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AN ATTEMPT AT REALISTIC SUPERCOMPOSITENESS

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ABSTRACT

Composite models of quarks, leptons and Higgs bosons based on softly-broken supercolour are both predictive and capable of giving some of the desired mass hierarchies. In the SQCD examples we have explored, however, massive neutrinos ($m_\nu \sim m_e$) and light spin- $\frac{1}{2}$ leptoquarks ($m_{\lambda q} \lesssim 1$ GeV) can only be avoided at the price of introducing further elementary particles (spectators).

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1. - INTRODUCTION

It is possible that we have discovered in the usual quark-lepton families the ultimate constituents of matter and in the gauge bosons of $SU(3)_C \times SU(2)_L \times U(1)_Y$ the basis of their interactions.

It looks more likely, however, that, as has been the case before, new structures will emerge when scales smaller than those yet tested, say 10^{-16} cm or less, are explored. If this should turn out to be so, one would be left with the puzzle of explaining the smallness of the dimensionless quantity $m_{q,l} \cdot r_{q,l} < 10^{-6}$, together with the fact that the presently "observed" non-gauge elementary particles have spin $\frac{1}{2}$. It has been an exciting, though frustrating, adventure to look for confining gauge theories that predict such a spectrum.

In a recent paper¹⁾, two of us have argued that confining supersymmetric gauge theories (Supercolour or SC), endowed with soft SUSY-breaking terms, can be promising candidates for such preonic theories. A few examples supporting this contention were given.

In this paper we report on a (non-exhaustive) search for a phenomenologically consistent model of this type for one family of composite quarks, leptons and Higgs particles. Unfortunately, the best model we are able to come up with is still unrealistic; it is ruled out (unless ad hoc spectators are added) by its prediction of some new light fermions. We wish nonetheless to present it for the following reasons:

- a) It shows the predictive power of supersymmetric composite models²⁾.
- b) It exhibits the multiple protection mechanism for quark and lepton masses suggested in Ref. 3).
- c) It can remove from the light spectrum the supersymmetric partners of quarks and leptons.
- d) It exhibits dynamical $SU(2)_L \times U(1)_Y$ breaking with a technicolour⁴⁾ scale $G_F^{-\frac{1}{2}}$ which is related (but not identical) to the supercolour compositeness scale Λ_{SC} .

- e) It is able to provide small, calculable masses for quarks and leptons without invoking new ["extended technicolour"⁴⁾] interactions.
- f) It has automatic strong CP conservation at the price of an "invisible" axion⁵⁾.

The phenomenological problems of the model reside in the fact that its fermionic content is richer than that of the usual model. Furthermore, the mass protection mechanism, being too democratic, tends to give the neutrino a mass $O(m_e)$ and some leptoquark fermions ($Q = \pm 2/3$, $B = \pm 1/3$, $L = \mp 1$ objects) with masses $O(1 \text{ GeV})$. In order to avoid this problem one can only invoke "spectators", i.e., elementary SC singlet fields that are able to form large Dirac masses with the leptoquarks and restore an almost massless neutrino. These modifications of the model look rather contrived, especially for the spectator leptoquarks.

Although we have not been able to avoid these light states in a natural way, we do not think that there is a general no-go theorem in this respect. We hope that, by a more judicious choice of the gauge group and/or of the preon quantum numbers, this problem can be circumvented, possibly within a scheme incorporating families.

This brings us to the second disappointing feature of our construction: families are not predicted, although there are straightforward ways of incorporating them by adding a family label to some of the preons.

The situation with respect to horizontal symmetries may turn out to resemble that of flavour in QCD where confinement and hadron compositeness do not lead to much "economy": flavour has to be added at the constituent level.

In the following sections, we shall present our specific model with this outline:

- A: Definition of the model, through specifications of the gauge group and the preonic quantum numbers (Section 2).
- B. Discussion of the model at three successive levels of approximation:
 - i) Level I, i.e., in the absence of supersymmetry breaking and of "weak" gauging (Section 3);
 - ii) Level II, i.e., in the presence of SUSY breaking but still without weak gauging (Section 4);

iii) Level III, i.e., after weak gauging and R-symmetry breaking (Section 5).

C. Finally, we show how to free the model of its main phenomenological problems by adding spectators (Section 6) and give some conclusions (Section 7).

2. - DEFINITION OF THE MODEL

The model we consider is of the type originally proposed in Refs. 6). It is essentially the specific model of Ref. 3) supplemented with a specific set of soft (mass) terms through which SUSY breaking is introduced¹⁾.

The preon binding supercolour group is taken to be $SU(6)_{SC}$ to which a standard "weak" gauge group

$$G_W = SU(3)_c \times SU(2)_L \times U(1)_Y \quad (2.1)$$

is added. All gauge particles are elementary both in $SU(6)_{SC}$ and in G_W . Supersymmetry will be broken at tree level in G_W (through gluino, photino, etc. mass terms), but only radiatively in the $SU(6)_{SC}$ gauge sector. The preon quantum numbers are given in Table 1 and correspond to an anomaly-free theory with asymptotically free (and hopefully confining) supercolour.

The model is not completely specified until one gives the explicit form of the superpotential and of the soft SUSY-breaking terms. For pedagogical reasons, we postpone the definitions and discussion of these terms until we study the theory at level II.

For the time being we will just recall that, in the absence of mass (or any other superpotential) and of weak gauging, this model has a global symmetry

$$G_F^{cl.} = G_F \times U(1)_A^a = SU(6)_L \times SU(6)_R \times U(1)_V \times U(1)_X \times U(1)_A^a \quad (2.2)$$

in which the $U(1)_A^a$ factor is SC-anomalous, and the non-anomalous $U(1)_X$ is an appropriate combination of axial $U(1)$'s that does not commute with SUSY (R-symmetry). We recall [see, for example, Ref. 7)] that, in this particular case, the

R-charge is zero for all the scalar fields and equal and opposite for the gaugino and the matter fermions [our convention will have $q_x(\psi) = -q_x(\lambda) = 1$].

3. - LEVEL I - THE SUPERSYMMETRIC MASSLESS LIMIT

We shall first analyze our model in the limit in which supersymmetry is exact and weak gauge couplings are turned off. The idea is that such a limit should already give us an idea of the final spectrum: at least all the light degrees of freedom should be present in that crude approximation, while states of mass $O(\Lambda_{SC})$ can be dropped from further discussion.

Of course, we can only argue in this way if G_W interactions are small at distances $O(\Lambda_{SC}^{-1})$ and if the SUSY-breaking mass parameters are chosen to be small with respect to Λ_{SC} .

At this level, our model reduces to the much-studied case of SQCD with an identical number of (super)colours and flavours ($N = N_f = 6$ in our case). Various arguments^{7),8)} and instanton-type calculations⁹⁾¹¹⁾ have led to an essentially complete understanding of the non-perturbative vacua of this theory and of its low-energy excitations. We claim that they are both given consistently by an effective Lagrangian generalizing slightly the one of Ref. 7) and containing the following $N^2 + 3$ chiral superfields:

$$\begin{aligned} T_{ij} &= \Phi_{i,\alpha} \tilde{\Phi}_j^\alpha & ; & \quad i, j = 1 \dots N \\ S &= g^2 / 32\pi^2 W^2 \\ X &= \det_{i,\alpha} \Phi_{i,\alpha} & ; & \quad \tilde{X} = \det_{i,\alpha} \tilde{\Phi}_i^\alpha \end{aligned} \quad (3.1)$$

Invariance requirements plus fulfilment of (anomalous and non-anomalous) Ward identities restrict L_{eff} to be of the form:

$$L_{eff} = L_D + L_F \quad (3.2)$$

where L_D are a set of invariant kinetic terms and the superpotential L_F is uniquely fixed to have the form

$$L_F = S [\log \det T / \Lambda^{2N} + f(Z)] \big|_{\theta^2} + \sum_i m_i T_{ii} \big|_{\theta^2} + h.c. \quad (3.3)$$

with f an arbitrary function of the chiral superfield

$$Z \equiv X \cdot \tilde{X} / \det T$$

In order to be general, we have added the most general superpotential, which in this case consists of supersymmetric masses m_i for our superfields $\Phi_i, \tilde{\Phi}_i$. The case $m_i \neq 0$ is the simplest one to describe. In this case, Konishi-type arguments⁸⁾ imply:

$$\begin{aligned} \langle T_{ij} \rangle &= \delta_{ij} m_i^{-1} \langle S \rangle \quad ; \quad i, j = 1 \dots N \\ \langle X \rangle &= \langle \tilde{X} \rangle = 0 \end{aligned} \quad (3.4)$$

Indeed, minimizing the potential that originates from (3.3), one finds the conditions

$$\begin{aligned} S \tilde{X} f' / \det T &= S X f' / \det T = 0 \\ S (T^{-1})_{ij} &= m_i \delta_{ij} \end{aligned} \quad (3.5)$$

which give (3.4) if $f' \neq 0$ and, in any case, are solved by (3.4). One also finds

$$\log \det T / \Lambda^{2N} + f(0) = 0 \quad (3.6)$$

which fixes

$$\det \langle T_{ij} \rangle = \Lambda^{2N} e^{-f(0)} \quad (3.7)$$

in agreement with instanton calculations^{10),11)}. The massless situation is more complicated. In this case, both our L_{eff} and anomaly arguments⁸⁾ require $\langle S \rangle = 0$. Besides, one just gets the further condition:

$$\log \det T / \Lambda^{2N} + f(Z) = 0 \quad (3.8)$$

which does not fix $\langle \det T \rangle$ and $\langle Z \rangle$. This appears to contradict the results of Ref. 11) if taken at face value. However, both the method of constrained instantons⁹⁾ and a recent extension of the ITEP approach¹²⁾ show that, in a Higgs phase with $\langle \phi_{i,\alpha} \rangle \sim \langle \tilde{\phi}_j^\alpha \rangle \gg \Lambda_{SC}$, there are still SUSY vacua with arbitrarily large values for $\langle \det T \rangle$. The calculations of Ref. 11) would then be justified only for the confining vacua at $\langle \phi_{i,\alpha} \rangle, \langle \tilde{\phi}_j^\alpha \rangle \ll \Lambda_{SC}$ and would be blind to the existence of Higgs-type vacua [the opposite being true for the calculations of Refs. 9)]. The amusing thing here is that our L_{eff} smoothly interpolates between the two pictures, since it possesses a flat direction in $\det T, X\tilde{X}$ space given by (3.8). To illustrate this point, take as an example

$$f(Z) = \log(1-Z) ; \quad L_F = S \log \left[\frac{\det T - X\tilde{X}}{\Lambda^{2N}} \right] - m_i \frac{T_{ii}}{\theta^2} + h.c. \quad (3.9)$$

In this case, the perturbative Higgs vacua

$$\langle \phi_{\alpha i} \rangle = \langle \tilde{\phi}_i^\alpha \rangle = v \delta_{\alpha i} ; \quad v \gg \Lambda \quad (3.10)$$

are recovered as

$$\langle \det T \rangle \gg \Lambda^{2N} ; \quad X\tilde{X} / \det T \rightarrow 1 \quad (3.11)$$

One can also verify that the SC singlets T_{ij} , X , $\tilde{X}$ contain, in the sense of complementarity^{*}), the massless superfields of the Higgs picture [including the Goldstone of the SC anomalous $U(1)_A^a$ symmetry].

In the confining picture we are adopting, S and one combination of $\det T$ and $X\tilde{X}$ pick up a mass $O(\Lambda_{SC})$ in agreement with R-invariance [recall that $q_x(S_\theta) = -q_x(T_\theta, (X\tilde{X})_\theta) = -1$], while the orthogonal combination and the remaining N^2 superfields remain massless. The resulting $N^2 + 1$ massless superfields are precisely what is needed³⁾ in order to satisfy all the appropriate 't Hooft consistency conditions¹⁴⁾. The actual vacuum stability group H depends, of

^{*}) For the use of complementarity in this context, see Ref. 13).

course, on which one of the many degenerate vacua (3.8) we choose. If, for instance, we take

$$\langle T_{ij} \rangle = c \delta_{ij} V_i^2 \quad ; \quad \langle X \rangle, \langle \tilde{X} \rangle \neq 0 \quad (3.12)$$

we break $U(1)_V$ and $SU(6)_R \times SU(6)_L$ down to a subgroup of $SU(6)_V$, which is $SU(6)_V$ itself if $V_1 = V$ and $U(1)^5$ if each V_1 is different. It can be seen, however, that in each case, the manifold of the vacua is the same, this being a consequence of the fact¹⁵⁾ that such a manifold represents the complexification of G_F and is invariant under a suitable complex extension of H . This is why the same set of degrees of freedom describes simultaneously all the inequivalent vacua which are present in the absence of mass and SUSY breaking.

4. - LEVEL II - ADDING SUSY BREAKING

We are now ready to add soft SUSY-breaking terms. We shall do this via chiral invariant bosonic mass terms with the hope that fermion masses will keep their chirality-based protection. It is at this point that the effect of the supercurrent anomaly comes into play. While the tree-level vacua all go to the zero value of ϕ and $\tilde{\phi}$ for non-vanishing (real) bosonic masses, the instanton-induced term (3.3) makes such a trivial vacuum disappear. Rather, the true ground state originates from playing the mass term against the effect of the anomaly.

In order to study how things go qualitatively, we have to represent the additional bosonic mass terms of the underlying theory:

$$\delta L = - \sum_i (\mu_i^2 \phi_i^* \phi_i + \tilde{\mu}_i^2 \tilde{\phi}_i^* \tilde{\phi}_i) \quad (4.1)$$

In Ref. 1), such terms were represented in L_{eff} by

$$\delta L_{\text{eff}} = - \sum_{ij} \bar{\mu}_{ij}^2 \pi_{ij}^* \pi_{ij} \quad (4.2)$$

where we have denoted by π_{ij} the lowest component of T_{ij} . Indeed, δL and δL_{eff} have the same transformation properties under G_F .

A more convincing way to take the perturbation (4.1) into account would be to introduce further (vector) superfields into L_{eff} . Following Ref. 16), these are:

$$R_{ij} \equiv \Phi_i^* e^V \Phi_j \quad ; \quad \tilde{R}_{ij} \equiv \tilde{\Phi}_i^* e^{\tilde{V}} \tilde{\Phi}_j \quad (4.3)$$

where $V, \tilde{V}$ are the supercolour gauge superfields in the $N, \bar{N}$ ($6, \bar{6}$) representations. The terms (4.1) can then be directly represented as

$$\delta L_{\text{eff}} = - \sum_i (\mu_i^2 R_{ii} + \tilde{\mu}_i^2 \tilde{R}_{ii}) \Big|_{\theta=\bar{\theta}=0} \quad (4.4)$$

Following the methods of Ref. 16), we have studied the consequences of adding (4.4) to our L_{eff} and we have checked that the results thus obtained agree with those derived from the simpler term (4.2), if R and $\tilde{R}$ carry only particles of mass $O(\Lambda_{\text{SC}})$, as expected on general grounds. In the following, we shall thus use the approach of Ref. 1) based on the SUSY-breaking term (4.2).

Since we want SUSY to be essentially unbroken at the Λ_{SC} scale, we shall take all the $\bar{\mu}$'s much smaller than Λ_{SC} . Also, invariance under G_W forces the conditions

$$\begin{aligned} \mu_1 &= \mu_2 = \mu_3 = \mu \\ \tilde{\mu}_1 &= \tilde{\mu}_2 = \tilde{\mu}_3 = \tilde{\mu} \\ \mu_5 &= \mu_6 \end{aligned} \quad (4.5)$$

while in general $\tilde{\mu}_5 \neq \tilde{\mu}_6$. Since the "diagonal" masses $\bar{\mu}_{ii}$ will play an important rôle in the following, we shall also adopt the simplified notation

$$\bar{\mu}_{ii} = \bar{\mu} \quad , \quad i=1,2,3 \quad ; \quad \bar{\mu}_{jj} = \bar{\mu}_j \quad , \quad j=4,5,6 \quad (4.6)$$

Adding the extra term (4.2) to the previous supersymmetric Lagrangian yields

$$L_{\text{eff}} = \frac{1}{\alpha \Lambda^4} (S^* S)_D + \frac{1}{\beta \Lambda^2} \text{Tr}(\tau^* \tau)_D + \\ + \left[S \log \det T / \Lambda^{12} \right]_{\theta^2} + \text{h.c.} - \frac{1}{\Lambda^2} \sum_{ij} \bar{\mu}_{ij}^2 \pi_{ij}^* \pi_{ij} \quad (4.7)$$

where we have taken canonical kinetic terms (up to a rescaling) and have omitted for simplicity the $X, \tilde{X}$ superfields. Neither one of these simplifications affects the main results of our subsequent discussion. The effective potential following from (4.7) is:

$$V_{\text{eff}} = \alpha \Lambda^4 \left| \log \det T / \Lambda^{12} \right|^2 + \beta \Lambda^2 \sum_{ij} |\pi_{ij}|^2 (\pi^{-1})_{ij}^2 + \\ + \Lambda^2 \sum_{ij} \bar{\mu}_{ij}^2 \pi_{ij}^* \pi_{ij} \quad (4.8)$$

where $\pi_\lambda \equiv S_{\theta=0}$. It is then easy to see that, for $\bar{\mu}^2 \ll \Lambda^2$, V_{eff} leads to a candidate diagonal vacuum where

$$\langle \pi_\lambda \rangle = 0 \quad ; \quad \langle \pi_{ij} \rangle = \delta_{ij} \pi_i \quad ; \quad \pi_i \bar{\mu}_i = \pi_j \bar{\mu}_j$$

$$\pi_1 = \pi_2 = \pi_3 = \left(\frac{\bar{\mu}_4 \bar{\mu}_5 \bar{\mu}_6}{\bar{\mu}^3} \right)^{1/6} \Lambda_{\text{sc}}^2, \\ \pi_4 = \left(\frac{\bar{\mu}^3 \bar{\mu}_5 \bar{\mu}_6}{\bar{\mu}_4^5} \right)^{1/6} \Lambda_{\text{sc}}^2, \\ \pi_5 = \left(\frac{\bar{\mu}^3 \bar{\mu}_4 \bar{\mu}_6}{\bar{\mu}_5^5} \right)^{1/6} \Lambda_{\text{sc}}^2, \quad \pi_6 = \left(\frac{\bar{\mu}^3 \bar{\mu}_4 \bar{\mu}_5}{\bar{\mu}_6^5} \right)^{1/6} \Lambda_{\text{sc}}^2 \quad (4.9)$$

In a range of values for the $\mu_i, \tilde{\mu}_i$ parameters, this can be shown to be the true ground state [see also Ref. 1)]. The vacuum (4.9) possesses some interesting features:

- a) It respects the R-symmetry ($\langle \pi_\lambda \rangle = 0$). This is of course important for fermionic mass protection.

- b) It preserves ordinary colour as a consequence of the vector-like nature of the preons under $SU(3)_C$.
- c) It breaks $SU(2)_L \times U(1)_Y$ (through $\langle \pi_5 \rangle, \langle \pi_6 \rangle$) as a consequence of the chiral nature of the preons under the electroweak group.

Since the Fermi scale (~ 100 GeV) is related directly to π_5 and π_6 and we wish it to be (much?) smaller than Λ_{SC} , we have to play on the mass parameters $\bar{\mu}$ and make $\pi_{5,6} \ll \Lambda^2$. This is possible due to the original huge vacuum degeneracy of the massless theory. Small masses can induce a widely asymmetric vacuum, as is evident from Eqs. (4.9). We are clearly led to the choice^{*}

$$\bar{\mu}^2 \sim \bar{\mu}_4^2 < (\ll) \bar{\mu}_{5,6}^2 \quad (4.10)$$

which yields

$$\begin{aligned} \pi_1 = \pi_2 = \pi_3, \pi_4 &> (\gg) \Lambda_{SC}^2 \\ \pi_5, \pi_6 &< (\ll) \Lambda_{SC}^2 \end{aligned} \quad (4.11)$$

Proceeding as in Ref. 1), we study now the light spectrum:

Fermions: They all remain massless except for the Dirac fermion, which had already become massive at Level I. This result is hardly surprising; since the R-symmetry has not been broken, 37 massless fermions are still necessary in order to fulfil the $U(1)_X^3$ anomaly constraint.

Bosons: Only the Goldstone bosons remain massless, while the others pick up , masses $O(\bar{\mu})$. For $\mu \neq \mu_4 \neq \mu_5 = \mu_6, \tilde{\mu} \neq \tilde{\mu}_4 \neq \tilde{\mu}_5, \tilde{\mu}_6$, the expected Goldstone bosons are:

- i) A colour octet from $SU(3)_L \times SU(3)_R \supset SU(3)_V \equiv SU(3)_C$.
- ii) Three "technipions" from $SU(2)_L \times U(1)_Y \supset U(1)_Q$.
- iii) Two more pseudoscalars associated with the spontaneous breaking of the two non-anomalous axial symmetries $U(1)_A, U(1)_{A'}$, of Table 2.

^{*}) Insisting on a huge ratio $\bar{\mu}_{5,6}/\bar{\mu}$ could turn out to be unnatural due to the coupling of different masses through higher orders in α_{SC} .

- iv) A combination of $\text{Im } X$, $\text{Im } \tilde{X}$ for the breaking of $U(1)_V$.

The second column of Table 3 summarizes the mass spectrum at Level II.

5. - LEVEL III - "WEAK" GAUGING AND R-SYMMETRY BREAKING

We finally gauge the subgroup of G_F :

$$G_W = SU(3)_C \times SU(2)_L \times U(1)_Y$$

and, at the same time, break the R-symmetry by introducing explicit masses for the gauginos of G_W , in particular for the gluino $\tilde{g}$.

Here we encounter the problem that the hypergluino mass m_λ is no longer protected by the R-symmetry (unlike that of the fermionic preons) and is generated by radiative corrections such as that of Fig. 1. Setting $m_\lambda = 0$ or $m_\lambda \ll \alpha_c m_{\tilde{g}}$ amounts thus to an unnatural fine tuning.

In conclusion, weak gauging and gluino mass bring in the following modifications of L_{eff} :

1. A direct (super)gluino mass term providing a correction¹⁾

$$\int L_{\text{eff}} = -m_\lambda \Pi_\lambda + \text{h.c.} ; m_\lambda \sim \alpha_c m_{\tilde{g}} \ll \mu \quad (5.1)$$

2. Weak gauge interactions for our composite fields. Neglecting Wess-Zumino type terms¹⁷⁾ (which should be irrelevant for our present purposes) we follow Ref. 1) and make the replacements

$$(S^* S)_D \rightarrow (S^* S)_D ; (T_{ij}^* T_{ij})_D \rightarrow T_{ij}^* [e^V]_{ij}^{\kappa\ell} T_{\kappa\ell} \quad (5.2)$$

where V includes all the gauge fields of G_W and is, in general, in a reducible representation of it. The modification (5.2) brings in additional contributions to the scalar potential in the form of (non-negative) D-terms.

We have to see whether these modify the previously-found vacuum (4.9). The colour D-term vanishes in the vacuum (4.9) and hence tends to stabilize it.

The $SU(2)_L \times U(1)_Y$ D-terms would tend to favour a vacuum in which

$$\langle \pi_{55} \rangle = \langle \pi_{66} \rangle \quad (5.3)$$

For $\bar{\mu}_5 = \bar{\mu}_6$, Eq. (5.3) is obviously satisfied by the vacuum (4.9).

We shall see, however, that in order to distinguish charge 2/3 from charge 1/3 quarks, it is important to allow $\bar{\mu}_5 \neq \bar{\mu}_6$. In this situation, there is a competition between the $SU(2)_L$ D-term:

$$g_w^2 (\pi_{55}^2 - \pi_{66}^2)^2 \quad (5.4)$$

and the mass term

$$\Lambda_{sc}^{-2} (\bar{\mu}_5^2 \pi_{55}^2 + \bar{\mu}_6^2 \pi_{66}^2) \quad (5.5)$$

The ratio of the two terms is of the order of

$$\frac{V_D}{V_{mass}} \approx g_w^2 \Lambda_{sc}^{-2} (\pi_{55}^2 + \pi_{66}^2) \frac{(\pi_{55}^2 - \pi_{66}^2)^2}{\bar{\mu}_{5,6}^2 \pi_{55}^2 (\pi_{55}^2 + \pi_{66}^2)} \lesssim \frac{m_w^2 (\bar{\mu}_5^2 - \bar{\mu}_6^2)^2}{\bar{\mu}_{5,6}^2 (\bar{\mu}_5^2 + \bar{\mu}_6^2)} \quad (5.6)$$

where we have anticipated the expression for the W,Z mass. Since $\bar{\mu}_{5,6}$ are the large masses, we can assume $V_D/V_{mass} \ll 1$, in which case the level II vacuum is not appreciably modified.

The extra term (5.1) induces the non-vanishing VEV for π_λ

$$\langle \pi_\lambda \rangle = \frac{m_\lambda}{\Lambda^2} \left(\sum_i \langle |\pi_{ij}^{-1}|^2 \rangle \right)^{-1} \approx \frac{m_\lambda}{\Lambda^2} \cdot \min \langle \pi_{ii} \rangle^2 \quad (5.7)$$

as can be seen from the extremum condition.

Consequently, the $S \log \det T$ term generates the fermionic mass term

$$L_{mass}^{tree} = \langle \pi_\lambda \rangle \prod_{k \neq i,j} \frac{\langle \pi_{kk} \rangle}{\langle \det \pi \rangle} \chi_{ij} \chi_{ji} \approx m_\lambda \frac{\min \langle \pi_{kk} \rangle^2}{\langle \pi_{ii} \rangle \langle \pi_{jj} \rangle} \frac{\chi_{ij} \chi_{ji}}{\Lambda^2} \quad (5.8)$$

where χ_{ij} is the θ -component of T_{ij} . One also gets extra masses for the spin-zero bound states, but these are small compared with the level II masses, since $m_\lambda \ll \bar{\mu}$. The true Goldstone bosons are not affected by these terms.

The other effects of weak gauging are of two types: the usual Higgs mechanism and the radiative contributions to fermionic and bosonic masses. Through the Higgs mechanism, the W and Z bosons and their SUSY partners, $\tilde{W}$ and $\tilde{Z}$, become massive. The longitudinal components of W and Z are provided by the would-be Goldstone bosons in the π_{ab} sector ($a, b = 5, 6$) which play the role of technipions. We obtain

$$M_W^2 = \frac{1}{2} g_2^2 (\pi_5^2 + \pi_6^2) / \Lambda^2 = \cos^2 \theta_W M_Z^2; \quad \sin^2 \theta_W = \frac{g_1^2}{g_1^2 + g_2^2} \quad (5.9)$$

Analogously, the $\tilde{W}$ and $\tilde{Z}$ get a mass by eating up the fermionic partners of the technipions. In the limit $m_\lambda = 0$ and $\pi_5 = \pi_6$ the two eigenvalues of the $\tilde{W}$ mass matrix coincide. In the case of a mass splitting mainly due to $m_\lambda \neq 0$, we get

$$M_{\tilde{W}_{1,2}} = M_W \pm \alpha_c m_{\tilde{g}} \quad (5.10)$$

whereas in the situation where $\pi_5 \neq \pi_6$ is the major source of this splitting, we obtain

$$M_{\tilde{W}_{1,2}} = M_W \pm \frac{\bar{\mu}_5^2 - \bar{\mu}_6^2}{\bar{\mu}_5^2 + \bar{\mu}_6^2} M_W \quad (5.11)$$

Hence, in any case, one of two mass eigenvectors is lighter than the W boson with relevant phenomenological consequences. Analogous results apply to the $\tilde{Z}$ case.

Finally, one gets¹⁾ fermionic masses from radiative corrections such as that of Fig. 2^{*)}. Keeping $SU(2)_L \times U(1)_Y$ gaugino masses of $O(m_{\tilde{g}})$ or less, their

^{*)} Presumably these diagrams have to be dressed with form factors¹⁸⁾. Since they already converge, however, the final outcome will differ at most by a factor $O(1)$ from the answer given below.

effect will be negligible compared to that of $m_{\tilde{g}}$ or m_{λ} . In the approximation $m_{\tilde{g}} < \mu_1$, one finds

$$\begin{aligned} m_{\nu}^{(rad)} &= 0 ; \quad m_{e^{\pm}}^{(rad)} = O\left(\frac{\alpha}{3\pi} m_{\tilde{g}} \bar{\mu}_4/\bar{\mu}_6\right) \\ m_{u,d}^{(rad)} &= \frac{\alpha_c}{\pi} C_F m_{\tilde{g}} \bar{\mu}/\bar{\mu}_{5,6} \\ m_{eq}^{(rad)} &= \frac{\alpha_c}{\pi} C_F m_{\tilde{g}} \log\left(\frac{\bar{\mu}_4+\bar{\mu}}{\bar{\mu}_4-\bar{\mu}}\right) \\ m_{oct.}^{(rad)} &= \frac{\alpha_c}{\pi} C_A m_{\tilde{g}} \log(\bar{\mu}/m_{\tilde{g}}) \end{aligned} \quad (5.12)$$

where $C_F = (N^2-1)/2N$, $C_A = N$ are the SC Casimirs. The identification of the ordinary fermions with appropriate composites is, of course, dictated by the quantum numbers and is shown in Table 1. Combining tree level [Eq. (5.8)] and radiative [Eq. (5.12)] masses, we get the final result [$m_{\lambda} = O(\alpha_c m_{\tilde{g}})$]

$$\begin{aligned} m_{e^{\pm}} &= \frac{\alpha_c}{\pi} C_F m_{\tilde{g}} a \left(\frac{\bar{\mu}_4 \bar{\mu}_6}{\sum \bar{\mu}_j^2} \right) + O(\alpha m_{\tilde{g}}) \\ m_{\nu} &= \frac{\alpha_c}{\pi} C_F m_{\tilde{g}} a \left(\frac{\bar{\mu}_4 \bar{\mu}_5}{\sum \bar{\mu}_j^2} \right) \\ m_u &= \frac{\alpha_c}{\pi} C_F m_{\tilde{g}} \left[a \frac{\bar{\mu} \bar{\mu}_5}{\sum \bar{\mu}_j^2} + \bar{\mu}/\bar{\mu}_5 \right] \\ m_d &= \frac{\alpha_c}{\pi} C_F m_{\tilde{g}} \left[a \frac{\bar{\mu} \bar{\mu}_6}{\sum \bar{\mu}_j^2} + \bar{\mu}/\bar{\mu}_6 \right] \\ m_{eq} &= m_{eq}^{(rad)} ; \quad m_{oct.} = m_{oct.}^{(rad)} \quad (\text{see 5.12}) \end{aligned} \quad (5.13)$$

where a is a number $O(1)$ related to the ratio $m_{\lambda}/\alpha_c m_{\tilde{g}}$.

Notice that $m_u = m_d$ unless $\bar{\mu}_5 \neq \bar{\mu}_6$. Concerning the bosons, we still had, at level II, a massless octet coming from the breaking of $SU(3)_R \times SU(3)_L$ to $SU(3)_C$. Since this symmetry is explicitly broken by colour interactions, we expect these bosons to acquire masses $O(\alpha_c)$. Indeed, the loop diagram of Fig. 3, after inserting the octet fermion mass obtained beforehand [Eq. (5.13)] gives

$$m(\pi_{oct}) \simeq m(\chi_{oct.}) \quad (5.14)$$

In other words, the two bosons of each member of the octet get masses $O(\mu)$ and $O(\alpha_c m_g^C \log \bar{\mu}/m_g)$.

In the diagonal sector $(\pi_{ii}, X, \tilde{X})$, the only remaining light bosons consist of

- a) The true Goldstone bosons of $U(1)_V$ and $U(1)_A$ (see Table 2);
- b) The Goldstone of the $U(1)_A$, symmetry which, because of a colour anomaly, behaves precisely as an axion. Since the scale of breaking of $U(1)_A$, is not the same as that of $SU(2)_L \times U(1)_Y$ breaking, the axion is of the "invisible"⁵⁾ type with mass and couplings depressed with respect to the ordinary axion by a factor $O(\pi_{5,6}/\pi, \pi_4) \approx \bar{\mu}/\bar{\mu}_{5,6} \ll 1$. Furthermore its bare coupling to 2γ is depressed by the presence of light composite fermions which essentially saturate the relevant non-Abelian anomalies^{*)}.

As far as the true (massless) Goldstone boson $\text{Im}(\pi_0)$ (Table 3) is concerned, one may worry about its couplings to the weak gauge bosons. One finds, however, that the only relevant coupling is through the Z_0 which connects $\text{Im} \pi_0$ to $\text{Re} \pi_0$ and to other massive states. Also the other true Goldstone boson $\text{Im}(X-\tilde{X})$ is found to decouple from all our light states, in particular from quarks and leptons.

Turning to the fermions, we find that the partner of π_0 , i.e., χ_0 of Table 3, can present a vertex $\bar{\chi}_0 \gamma_\mu \chi_0 Z_\mu$ if $\pi_5 \neq \pi_6$. This could simulate some extra neutrinos contributing to the Z_0 total width.

On the other hand, the results (5.13) for the fermions exhibit two conspicuous problems:

- i) The neutrino gets a Dirac mass of $O(m_e)$ from the tree-level mechanism due to $m_\lambda \neq 0$ [Eqs. (5.8)].
- ii) The spin- $\frac{1}{2}$ leptoquarks (which have charge $\pm 2/3$) receive [Eqs. (5.13)] masses $O(\alpha_c m_g^C)$ which could be dangerously small. In order to see this point, let us try to keep the leptoquarks heavier than 10 GeV. It follows that

*) This question has been studied meanwhile in some detail by E. Guadagnini, K. Konishi and M. Mintchev (private communication).

$m_{\tilde{g}} > 10^2 \text{ GeV}$, but also, from Eq. (5.1), that $\bar{\mu}, \bar{\mu}_4 > 10^2 \text{ GeV}$. We also want

$$m_{\tilde{q}}/m_{\tilde{e}q} \approx \bar{\mu}/\bar{\mu}_5 \approx 10^{-3} \Rightarrow \bar{\mu}_5 > 10^2 \text{ TeV} \quad (5.15)$$

and, on the other hand, we have from (5.9) and (4.9),

$$\begin{aligned} \Lambda_{SC} &= \frac{M_{WU}}{g_2} \left(\frac{\Lambda_{SC}^2}{\pi_{5,6}} \right)^{1/2} \approx (300 \text{ GeV}) \cdot (\bar{\mu}_{5,6}/\bar{\mu})^{2/3} \\ &\approx 30 \text{ TeV} < \bar{\mu}_{5,6} \end{aligned} \quad (5.16)$$

We are thus led to a situation where $\bar{\mu}_{5,6} \gg \Lambda_{SC}$, which is the opposite of our initial assumption (SUSY breaking small compared to Λ_{SC}).

Suppose in fact that we require

$$\bar{\mu}_{5,6}/\Lambda_{SC} \lesssim 10^{-1} \quad (5.17)$$

we then get from (4.9)

$$\bar{\mu}^2 < (30 \text{ GeV})^3 / \bar{\mu}_5 \quad (5.18)$$

Since we want $\bar{\mu} < \bar{\mu}_5$, we must have $\bar{\mu}_5 > 30 \text{ GeV}$ (i.e., $\Lambda_{SC} > 300 \text{ GeV} \approx G_F^{-1/2}$), but the higher we take $\bar{\mu}_5$, the smaller $\bar{\mu}$ [Eq. (5.18)] and $m_{\tilde{\chi}q}$ [Eqs. (5.13)] will be. We are squeezed into a compromise, e.g.,

$$\Lambda_{SC} \gtrsim 1 \text{ TeV} \Rightarrow \bar{\mu}_{5,6} \gtrsim 100 \text{ GeV} \Rightarrow \bar{\mu} < 16 \text{ GeV} \quad (5.19)$$

which gives (for $m_{\tilde{g}} \approx \bar{\mu}, \bar{\mu}_4$):

$$m_{\tilde{e}q} \approx 1.5 \text{ GeV} ; m_e, m_q \approx 250 \text{ MeV} \quad (5.20)$$

The quark and lepton masses could be fit for the second generation, but the leptoquark mass is obviously too small.

6. - AN UNAPPEALING REMEDY: SPECTATORS

We now present the only remedy we have been able to find for the problems i) and ii) mentioned above. This consists of adding spectators, i.e., SC-singlet elementary fields, to our list of preons. Let us look at the neutrino sector in some detail. The singlet N which is needed is neutral with respect to the entire gauge group. It is easy to see that, if we want the R symmetry, as well as B and L , not to be broken by the interactions of N with Φ and $\tilde{\Phi}$, we can give charges $-2, 0$ and $+1$ (under R, B and L) to N and add the gauge invariant interaction

$$\int L = \int d^4\theta N^* N + \int d^2\theta g_N \Phi_\alpha^4 \tilde{\Phi}_5^* N + h.c. \quad (6.1)$$

to the fundamental Lagrangian. Writing this new term in the effective theory as

$$\int L_{eff} = (N^* N)_D + g_N T_{45} \cdot N / \theta^2 + h.c. \quad (6.2)$$

we find easily that the new vacuum is consistent with the old vev's, with the addition of $\langle N \rangle = 0$.

For the masses nothing new happens in the bosonic sector, while in the fermionic sector one gets a 3×3 neutrino mass matrix of the form

$$\begin{matrix} \nu \\ N \\ \nu^c \end{matrix} \begin{pmatrix} 0 & 0 & 0(m_e) \\ 0 & 0 & g_N \Lambda \\ 0(m_e) & g_N \Lambda & 0 \end{pmatrix} \quad (6.3)$$

Thus the combination $\nu_{Phys} = \nu - (0(m_e)/g_N \Lambda) N$ remains strictly massless (as long as L number is conserved). We have just rediscovered in our context one of the known mechanisms¹⁹⁾ for eliminating the right-handed neutrino and its dangerous Dirac mass with the ordinary neutrino.

The situation with leptoquarks is similar, except that now the spectators

carry colour and hypercharge. We need the set of spectators [under $SU(6)_{SC} \times G_W$]:

$$M_i = (1, 3, 1, 4/3) ; \quad \tilde{M}_i = (1, \bar{3}, 1, -4/3) \quad (6.4)$$

which evidently do not change (gauge) anomaly cancellations. Again one can use symmetries to restrict the possible interactions of $M_i, \tilde{M}_i$ with $\Phi, \tilde{\Phi}$, and obtain large Dirac masses between the composite leptoquarks and the spectators (which have been chosen ad hoc for that purpose).

Of course, this way of proceeding is far from satisfactory. A possible alternative to the $M_i, \tilde{M}_i$ spectators is to start from a gauge group G_W in which $SU(3)_C$ is part of a Pati-Salam $SU(4)$. If we could think of a mechanism that spontaneously breaks $SU(4)_{PS}$ to $SU(3)_C$, the usual supersymmetric Higgs mechanism would be able to eat up spin-zero and spin- $\frac{1}{2}$ leptoquarks and give them masses $O(\Lambda_{PS})$. Unfortunately, with the order parameter at our disposal, we cannot find a way to achieve such a breaking.

7. - CONCLUSIONS

In spite of several encouraging features, like the prediction of light composite fermions, the models we have been able to construct are ruled out by lack of sufficient hierarchy in the fermionic mass spectrum.

It seems to us that, in order to avoid the above problems, one must start with SUSY gauge theories in which the preon-quantum numbers lead to a radically different global symmetry structure. One possibility would be to enlarge the flavour group (and possibly G_W itself) by adding a family label to some of the preons while keeping the size of the SC group unaltered. It is known^{7),20)} that in this case the vev's of the relevant order parameters are controlled by the masses in such a way that they are typically much smaller than Λ_{SC} . This could avoid the unwanted constraint $\bar{\mu}_{5,6} > \Lambda_{SC}$ which we encountered in order to get heavy enough leptoquarks. More radical alternatives, like that of a set of preons which are chiral under SC, should also be explored.

Unfortunately, we are short of hints as to which precise direction to follow in the search for modifications. At the same time, the techniques for analyzing a given model in sufficient detail are still primitive and require many long and

tedious calculations. We hope, however, to have convinced the reader that supercompositeness is not a game in handwaving and parameter adjustment. Its predictivity is what doomed our attempt, but it can be invaluable in providing a unique, workable scheme, if it exists. The stakes are well worth the effort, we believe.

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<div style="display: inline-block; transform: rotate(-45deg); transform-origin: center;"> Transfor. under L.H. S.field </div>	SU(6) _{SC}	SU(3) _C	SU(2) _L	U(1) _Y	Q = T ₃ +Y/2	
$\Phi_{\alpha}^{1,2,3}$	6	$\bar{3}$	1	-1/3	-1/6	
Φ_{α}^4	6	1	1	1	1/2	
$\Phi_{\alpha}^{5,6}$	6	1	2	0	$\pm 1/2$	
$\Phi_{1,2,3}^{\alpha}$	$\bar{6}$	3	1	1/3	1/6	
$\Phi_{4,5}^{\alpha} \quad 2 \times$	$\bar{6}$	1	1	-1	-1/2	
Φ_6^{α}	$\bar{6}$	1	1	1	1/2	
$\begin{pmatrix} u \\ d \end{pmatrix} = \Phi_{\alpha}^{5,6} \Phi_{1,2,3}^{\alpha}$	1	3	2	1/3	$\begin{matrix} 2/3 \\ -1/3 \end{matrix}$	$\left. \begin{matrix} \\ \\ \\ \\ \\ \\ \end{matrix} \right\} \text{1 family (with } \nu_R)$
$\begin{pmatrix} \nu \\ e \end{pmatrix} = \Phi_{\alpha}^{5,6} \Phi_4^{\alpha}$	1	1	2	-1	$\begin{matrix} 0 \\ -1 \end{matrix}$	
$\bar{u} = \Phi_{\alpha}^{1,2,3} \Phi_5^{\alpha}$	1	$\bar{3}$	1	-4/3	-2/3	
$\bar{d} = \Phi_{\alpha}^{1,2,3} \Phi_6^{\alpha}$	1	$\bar{3}$	1	2/3	1/3	
$\nu^c = \Phi_{\alpha}^4 \Phi_5^{\alpha}$	1	1	1	0	0	
$e^+ = \Phi_{\alpha}^4 \Phi_6^{\alpha}$	1	1	1	2	+1	
$\bar{q} = \Phi_{\alpha}^4 \Phi_{1,2,3}^{\alpha}$	1	3	1	4/3	2/3	
$\bar{l} = \Phi_{\alpha}^{1,2,3} \Phi_4^{\alpha}$	1	$\bar{3}$	1	-4/3	-2/3	B=-1/3, L=1
Oct = $\Phi_{\alpha}^{1,2,3} \Phi_{1,2,3}^{\alpha}$	1	8	1	0	0	B = L = 0

Table 1: The SU(6)_{SC}, SU(3)_C, SU(2)_L, U(1)_Y and U(1)_{em} quantum numbers of the preonic superfields ($\Phi, \bar{\Phi}$) and of the composite superfields ($\Phi\bar{\Phi}$).

	U(1) symmetries and light bosons						
Bos. Fields	$U(1)_V$	B	L	$U(1)_Y$	$U(1)_X$	$U(1)_A$	$U(1)_{A'}$
$\phi_{1,2,3}$	1	1/6	1/6	-1/3	0	0	1
ϕ_4	1	-1/6	-5/6	1	0	-2	-3
$\phi_{5,6}$	1	-1/6	1/6	0	0	1	0
$\tilde{\phi}_{1,2,3}$	-1	-1/6	-1/6	1/3	0	0	1
$\tilde{\phi}_4$	-1	1/6	5/6	-1	0	-2	-3
$\tilde{\phi}_5$	-1	1/6	-1/6	-1	0	1	0
$\tilde{\phi}_6$	-1	1/6	-1/6	1	0	1	0
Status	Broken. Spon.	Unbroken	Unbroken	Gauged, broken	Unbroken	Broken Spon.	Colour an. broken
Light Bosons	X, $\tilde{X}$	---	---	Eaten	---	Im π	Axion

Table 2: Global and local U(1) quantum numbers of the preonic superfields. The line denoted by "status" describes the fate of each U(1) symmetry of the theory. Notice that B and L are exact symmetries; the same occurs for $U(1)_X$ as long as we do not introduce some explicit breaking, such as a mass for gluinos.

Fermions	Level II Mass		Level III Mass	
Leptons	0		$\alpha_c m_{\tilde{g}}(\tilde{\mu}_4/\tilde{\mu}_{5,6})$	
Quarks	0		$\alpha_c m_{\tilde{g}}(\tilde{\mu}/\tilde{\mu}_{5,6})$	
Spin- $\frac{1}{2}$ leptoquarks	0		$\alpha_c m_{\tilde{g}}(\tilde{\mu}\tilde{\mu}_4/\tilde{\mu}^2+\tilde{\mu}_4^2)$	
Octet	0		$\alpha_c(C_A/C_F)m_{\tilde{g}} \log(\tilde{\mu}/m_{\tilde{g}})$	
Higgsinos	0		eaten by $\tilde{W}, \tilde{Z}$	
χ_λ, χ_0	Λ_{SC}		Λ_{SC}	
χ_0', χ_0''	0		$\alpha_c m_{\tilde{g}}, \alpha_c m_{\tilde{g}}(\tilde{\mu}/\tilde{\mu}_{5,6})^2$	
$\tilde{X}_F, \tilde{X}_F$	0		$\alpha_c m_{\tilde{g}} f(\tilde{\mu}_1/\tilde{\mu}_j)$	
Gluino	-		$m_{\tilde{g}}$ (input)	
$\tilde{W}, \tilde{Z}$	-		$m_{W,Z} \pm \max(\alpha_c m_{\tilde{g}}, (\tilde{\mu}_5^2 - \tilde{\mu}_6^2)/\tilde{\mu}_{5,6}^2 m_{W,Z})$	
Bosons	Level II (mass) ²		Level III (mass) ²	
Sleptons	$\tilde{\mu}_{5,6}^2$		$\tilde{\mu}_{5,6}^2$	
Squarks	$\tilde{\mu}_{5,6}^2$		$\tilde{\mu}_{5,6}^2$	
Scalar leptoquarks	$(\tilde{\mu} \pm \tilde{\mu}_4)^2$		$(\tilde{\mu} \pm \tilde{\mu}_4)^2$	
	scalar	pseudoscalar	scalar	pseudoscalar
Octet	$\tilde{\mu}^2$	0	$\tilde{\mu}^2$	$(\alpha_c m_{\tilde{g}} \log \tilde{\mu}/m_{\tilde{g}})^2$
Higgses	$\tilde{\mu}_{5,6}^2$	0	$\tilde{\mu}_{5,6}^2 + 0(m_W^2)$	eaten by W, Z
π_λ, π_0	Λ_{SC}^2	Λ_{SC}^2	Λ_{SC}^2	Λ_{SC}^2
π_0', π_0''	$\tilde{\mu}^2, \tilde{\mu}_{5,6}^2$	0, 0	$\tilde{\mu}^2, \tilde{\mu}_{5,6}^2$	$\lesssim \Lambda_{QCD}^4/\Lambda_{SC}^2; 0$
$X + \tilde{X}$	$\tilde{\mu}_1^2$	$\tilde{\mu}_1^2$	$\tilde{\mu}_1^2$	$\tilde{\mu}_1^2$
$X - \tilde{X}$	$\tilde{\mu}_1^2$	0	$\tilde{\mu}_1^2$	0
W, Z	-	-	$g_2^2 \Lambda_{SC}^2 \left(\frac{\tilde{\mu}\tilde{\mu}_4^{1/3}(\tilde{\mu}_5^2 + \tilde{\mu}_6^2)}{(\tilde{\mu}_5\tilde{\mu}_6)^{5/3}} \right)$	$m_W/\cos \theta_W$

Table 3: Order of magnitude of the fermionic and bosonic masses at the levels II and III of the theory. See Sections 4 and 5 of the text for details.

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FIGURE CAPTIONS

- Fig. 1 : Example of radiative corrections inducing the supergluino mass m_λ once the R-symmetry is broken by the gluino mass $m_g^\sim$.
- Fig. 2 : Radiative corrections giving rise to fermionic masses. These diagrams give the contributions in Eqs. (5.12).
- Fig. 3 : Radiative mass of the tree-level massless octet which comes from the breaking of $SU(3)_L \times SU(3)_R$ to $SU(3)_C$.

Fig. 1

Fig. 2

Fig. 3