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ABSTRACT

Arguments are given in favor of a hard asymptotic behavior of dynamically generated masses, its consequences for technicolor models are analyzed and a model is proposed, where effects of flavor changing neutral currents are highly supressed and pseudo Goldstone bosons get masses of O(30-90) Gev.

1. INTRODUCTION

Theories with dynamical symmetry breaking, particularly technicolor (TC) models[1], are more attractive than those involving the usual Higgs mechanism based on fundamental scalar fields.Their attraction resides in the introduction of new gauge symmetries leading to composite scalars in place of the fundamental ones,elim inating the proliferation of free parameters and implementing cal culability settled upon the knowledge of the dynamics of these theories.

In technicolor models the computation of current fermion masses is extremely dependent of the dynamically generated fermion mass asymptotic behavior. An example of how the dynamics is impor tant and modify the predictions of the theory has been evidenced by Holdom[2] and Georgi and Glashow[3], where the existence of a non-trivial ultraviolet (UV) fixed point in the TC theory leads to a fermion mass weakly dependent on the extended technicolor (ETC)[1] gauge boson mass, consequently this one can be made heavier and problems of flavor changing neutral currents (FCNC) eliminated.

The approach of references (2) and (3) makes use of what is called regular solution of the fermion self-energy[4], whose de- pendence on the fermionic anomalous dimension is strengthened due to the non-trivial UV fixed point. There is another solution for the dynamically generated mass, with much harder asymptotic behav ior than the regular one, and called irregular[4]. The choice be- tween those has been based on the fact that by means of the opera tor product expansion (OPE) we obtain only the regular solution[4,5].

Langacker[6] has criticized the naive OPE procedure, and has argued about the uniqueness of the irregular solution. Recently a very nice application of the Nambu-Jona-Lasinio non-perturbative method to QCD has led uniquely to the irregular solution[7], and similar conclusion was also obtained by means of an effective potential method[8]. The problem can also be attacked by its phenomenological side, asking if a sensible model of dynamical mass generation based on the irregular solution can be found[9]. Obviously, if this solution does not lead to a disastrous phenomenology, but, oppositely, to a successful one, we shall have an incentive to deepen our efforts towards a better knowledge of the dynamics of mass generation.

The aim of this paper is directed to an analysis of the phenomenological consequences of the irregular solution, and it is organized as follows: In the second section we give our arguments in favor of the irregular solution, and the relation between this one and the asymptotic behavior of the electromagnetic form factor of the pion is discussed. In the third section we discuss the determination of fermion masses in simple models, and in the next section a model is introduced. Section five contain an evaluation of technifermion and pseudo-Goldstone masses, and in section six we present our conclusions.

2. THE ASYMPTOTIC BEHAVIOR OF DYNAMICAL FERMION MASSES

The asymptotic behavior of dynamically generated fermion masses in non-Abelian gauge theories has two possible solutions[4] :

$$\Sigma_I(-p^2) \underset{-p^2\to\infty}{\sim} \frac{\mu^3}{p^2}\left(\ln\frac{-p^2}{\mu^2}\right)^{3c/16\pi^2 b}, \tag{1}$$

$$\Sigma_{II}(-p^2) \underset{-p^2\to\infty}{\sim} \mu\left(\ln\frac{-p^2}{\mu^2}\right)^{-3c/16\pi^2 b}, \tag{2}$$

called regular (Σ_I) and irregular (Σ_{II}), where μ is a dynamical mass,

$$c = \frac{1}{2}\left[C_2(R\psi)+C_2(R\chi)-C_2(R\phi)\right]$$

and

$$b = \frac{1}{16\pi^2}\left[\frac{11}{3}\;C_2(G)-\sum_{R_i}\frac{4}{3}\;T(R_i)\right]$$

In the above expressions we considered fermions ψ and χ in the representations $R\psi$ and $R\chi$, $R\phi$ is the condensate representation contained in $R\psi \times R\chi$, $C_2(R_i)$ their second Casimir operator, and b is the g^3 coefficient of the β function expansion.

In the following we shall briefly digress about previous results leading exclusively to Eq. (2), and compute the electromagnetic form factor of the pion ($F_\pi(q^2)$), since it has been argued that (1) and (2) imply in different asymptotic behaviors for $F_\pi(q^2)$[10].

A very beautiful and simple determination of the Eq. (2) behavior was done by Chang and Chang[7]. Following Nambu and Jona-Lasinio the authors of ref. (7) look for consistent massive solutions of the QCD Lagrangian

$$\mathcal{L} = (\mathcal{L}_0 - M\bar{\psi}\psi) + (\mathcal{L}_{INT} + \delta M\bar{\psi}\psi), \tag{3}$$

with $\delta M = M$. Computing the high energy behavior in the broken phase we are naturally led to Eq. (2), which is not different from the behavior with bare masses, however the interesting point is that (2) comes from a consistency condition (analog to the gap equation of Nambu and Jona-Lasinio)proved at two-loop accuracy , which in a low momentum limit could be read roughly as $g^2c \approx$ cte., agreeing with results expected intuitively[11] and numerically[12] .

We have computed the effective potential for composite operators of Cornwall, Jackiw and Tomboulis and verified that the chiral breaking of non-Abelian gauge theories is dominated by the irregular solution and happens when[8]

$$g^2 \gtrsim 0.273\ (3c/4\pi^2 - b)^{-1} \quad , \tag{4}$$

for small b we obtain

$$g^2 c \gtrsim 3.6 \tag{5}$$

It is interesting to note that this value is very close to the one obtained by Kogut et al ($g^2c \gtrsim 4$)[12,F1]. There is also a constraint for the existence of chiral symmetry breaking in this scheme (weak coupling), which is[8]

$$(3c/8\pi^2 - b) > 0 \tag{6}$$

One should emphasize that the effective potential[8] has a vacuum expectation value (VEV) proportional to $1/q^2$, and it is needless to reinstate the arguments of Quinn and Gupta[14] of how ill defined is the OPE in that case, fact remembered elsewhere[6,7,9] as the one behind the failure of the OPE in obtaining Eq. (2).

It has been argued that solutions (1) and (2) give different asymptotic behavior to the electromagnetic form factor of the pion[10]. We point out that both solutions give $F_\pi(q^2) \to O(1/q^2)$ (the canonical power counting behavior) when $-q^2 \to \infty$.

To compute $F_\pi(q^2)$ we make use of the dynamical perturbation theory (DPT) proposed by Pagels and Stokar[15], stating that in the computation of a given diagram the amplitudes that do not vanish in all orders of perturbation theory are given by their free field values, and amplitudes vanishing as $e^{-4\pi/g^2(q^2)}$ are given by their lowest order approximation. In lowest order of DPT $F_\pi(q^2)$ is given by[15]

$$F_\pi(q^2) \underset{q^2 \to \infty}{\sim} \frac{3 \ln 2}{q^2 (2\pi)^2 f_\pi^2} \int_0^\infty dp^2 \, \frac{p^2 \, \Sigma^2(-p^2)}{p^2 + \Sigma_1^2(p^2)} \, , \qquad (7)$$

and this result is obtained as long as the normalization condition is obeyed, i.e. $F_\pi(0) = 1$, and for Σ_{IR} this does not imply any other constraint than the one already exhibited in Eq. (6). Using the correctly normalized behavior of (2)[16]

$$\Sigma_1(p^2) = \mu \left(1 + b g^2 \ln p^2/\mu^2 \right)^{-3c/16\pi^2 b} \, , \qquad (8)$$

and Eq.(7), we get

$$F_\pi(q^2) \underset{q^2 \to \infty}{\sim} \frac{0.05 \, \mu^4}{f_\pi^2 \, b g^2 (3c/8\pi^2 b - 1)} \, \frac{1}{q^2} \, . \qquad (9)$$

We see that within the DPT approximation $F_\pi(q^2)$ does go as $1/q^4$

asymptotically. We do not attempt to compare it to the experimental result due to the critical dependence on the not well known values of μ and g^2.

After the above analysis we believe to have sufficient reasons to justify our choice of the irregular solution, which will be assumed in the remaining sections without any further consideration.

3. CURRENT FERMION MASSES

In this section we shall discuss the determination of current fermion masses in simple models, gradually obtaining hints to elaborate a realistic model.

The diagram that generates fermion masses is displayed in Fig. 1, where we assume that the gauge boson intermediates a non-Abelian interaction with a running coupling g_w, and has a mass M_W. The fermions χ interact strongly with a coupling g_s, inducing a dynamical mass μ_s. The computation of m_f in DPT (using Eq. (8) and the running coupling g_W) entails

$$m_f \approx \frac{3}{4} (c_W/\pi)\, \alpha_W\, \mu_s \left[1 + b g_s^2 \ln \frac{M_W^2}{\mu_s^2} \right]^{-3c/16\pi^2 b} , \qquad (10)$$

where c_W is the Casimir operator associated to the vertex $\chi\psi W$, $\alpha_W = g_W^2/4\pi$, b and c are related to the strong interaction, and to obtain (10) we assumed $g_s^2 c \sim 1$. The rather weak dependence on M_W will permit its increase to the point of complete elimination of FCNC problems[9].

It is worth mentioning that DPT gives a very concise rule to compute non-perturbative diagrams like Fig. 1, however it cannot be trivially applied to QED and scalars theories, because their amplitudes grow at large momenta. Therefore we shall compute $m_q(\mathrm{QED})$ using the following procedure: we assume α_{QED} constant until a certain grand unified theory (GUT) scale (which is a quite reasonable approximation), and from this scale to infinity we use the amplitudes of this unified theory.

To exemplify the above comment let us compute the electromagnetic contribution to the mass of a quark of charge Q, assuming the existence of a GUT at high energies, we obtain

$$m_q \simeq \frac{3\,Q^2}{4\,\pi}\,\frac{\alpha_{QED}\;\mu_{QCD}}{g^2 b\,(3c/16\pi^2 b - 1)} + \frac{3\,C_{GUT}}{4\,\pi}\,\mu_{QCD}\,\alpha_{GUT}\left[1 + b g^2 \ln \frac{M^2_{GUT}}{\mu^2_{QCD}}\right]^{-3c/16\pi^2 b} \qquad (11)$$

notice that. a) the QCD condensation is important when determining the current quark mass of the first generation ($m_q \sim O(1)\,\mathrm{Mev}$ for $\mu_{QCD} \sim 300\,\mathrm{Mev}$), b) for QCD with six quarks the first term in (11) is negative. This negative contribution to m_q had already been noticed before by Brodsky et al[17] in a different context, and is behind the good result obtained, in the more sophisticated approach of Chang and Li[18], for the up-down mass difference. The problem encountered by DPT for QED and scalar theories is also present in the renormalization group technique used in references (17)and (18), which is dominated by QED above the GUT scale, becoming trustless as we go up in the momentum scale, consequently the convenience of grand unification becomes translucent, freeing us from an apparent unnaturality of theories which are not asymptotically free.

As a simple unrealistic example we study the dynamical mass generation in the Georgi-Glashow model[19] (i.e. without the Higgs structure responsible for the breaking of $SU(2)_L \times U(1)$). Straightfowardly we see that quarks of same charge acquire a mass $m_q \sim O(1)$ Mev accordingly to Eq. (11) due to QCD condensation. The surprising fact is that QCD also induce lepton masses. Due to SU(5) interaction, from Eq. (10) with $\mu_s = \mu_{QCD} \sim 300$ Mev, $\alpha_W = \alpha_{GUT} \sim 1/40$, $M_W = M_{GUT} \sim 10^{15}$ Gev and $\alpha_s \sim 0.39$ (critical coupling constant obtained from Eq. (4)), we obtain $m_e \sim 0.18$ Mev, which within the present uncertainties is an excellent result. The fermion mass spectrum in this version of the SU(5) model would be given by

$$m_i = O(1-10) \text{ Mev}; \quad m_{\nu e} = 0 \quad , \tag{12}$$

where i signals the e, u and d fermion types, the neutrinos are massless, and we recall that masses determined at GUT scales should be corrected to the low energy values. The model is very reasonable concerning the first fermion generation, although is a failure for the higher families and weak boson masses.

One possible extension of the SU(5) model including a TC sector, which shall provide the heavier mass scale of weak bosons, is the Farhi-Susskind SU(7) model[20]. This theory has a fermion content belonging to the representations [7,2] + [7,4] + [7,6] (where by [N,m] we mean the SU(N) representation with m antisymmetric indices). We follow the assignments of ref.(20) and assume that SU(7) is broken directly to $SU(3)_c \times SU(2)_{TC} \times SU(2)_L \times U(1)$. We shall not bother ourselves about the origin of this breaking, and do not discard the possibility of fundamental scalar bosons (maybe associated to supersymmetry).

The $SU(2)_{TC}$ group will generate the condensates[20]

$$\langle \bar{U}\, U\rangle = \langle \bar{D}\, D\rangle = \langle \bar{E}\, E\rangle = \langle \bar{N}\, N\rangle \sim -\mu^3_{TC}$$

of techniquarks (U and D) and technileptons (E and N), and if we approximate Eq. (10) simply by

$$m_f \approx c^f_W\, \alpha_W\, \mu_s \tag{13}$$

where $c^f_W\, \alpha_W$ accounts for the charge of a given fermion f under the gauge interaction W, they produce a fermionic spectrum which can be written as

$$\begin{matrix} m_{\nu_e} & 0 \\ m_e(m_\mu) & c^e_7\, \alpha_7\, \mu_{TC} \\ m_u(m_c) & (c^u_1 \alpha_1 + c^u_7 \alpha_7)\mu_{TC} \end{matrix} \;+\; \begin{pmatrix} \text{QCD} \\ \text{contributions} \\ \mu_{TC} \to \mu_{QCD} \end{pmatrix} . \tag{14}$$

Notice that in (14) we should add contributions of QCD condensates, surpassed by far by those of TC. Contrarily to the first example, in this one all the fermions get very heavy masses.

At this point we have learned enough to establish some rules directed to a realistic model building. Basically we need one unified group containing the standard model and a TC group, this GUT will play the role of the ETC group, and, more fundamentally, we must have some symmetry preventing the light generations from getting masses in leading order when interacting with technifermions. In the next section we will present a model with these characteristics.

4. A MODEL

To build a theory containing the standard model ($SU(3)_C \times SU(2)_L \times U(1)$), a TC group (stronger than QCD) and three fermion generations is not an easy task, and within certain popular rules is even impossible[21]. Here we shall forsake unification of TC with three families for a semi-simple unification of a SU(N)theory with one family times a horizontal symmetry[22,23]. Our model is based on the group SU(9) with the following fermionic representations

$$5 \times \left[9,8\right] + 1 \times \left[9,2\right] \quad . \qquad (15)$$

This theory is anomaly free, asymptotically free, has one SU(5)family, and can be broken by a set of fundamental scalars to $SU(4)_{TC} \times SU(3)_C \times SU(2)_L \times U(1)$, as long as we respect constraints of proton lifetime and FCNC (possible consequences of a different pattern breaking will be discussed afterward). Notice that (15) belongs to a class of theories with the same properties, given by

$$\left[N,2\right] + A(N,2)\left[N,N-1\right] \quad , \qquad (16)$$

where $A(N,m)$ is the anomaly of the representation $\left[N,m\right]$.

The replication of families is due to a horizontal symmetry. We choose SU(3) as the family group, which may be local[22] or global[23], and will be broken by two scalar octets. If the symmetry is made local (although we shall keep the discussion in general grounds admitting both possibilities) a given group of fermions must be introduced to render the theory anomaly free, these fermions will get heavy masses in the $SU(3)_H$ breaking, assumed to happens at $M \sim M_{GUT}$. Since our main purpose is the determination of a realistic fermionic spectrum, we now turn to the question of fermionic assignment under $SU(3)_H$, which is a delicate one if we want to obtain the Fritzsch matrix[24]:

$$m_f = \begin{bmatrix} 0 & A_f & 0 \\ A^*_f & 0 & B_f \\ 0 & B^*_f & C_f \end{bmatrix} \quad , \qquad (17)$$

where $|C_f| >> |B_f| >> |A_f|$,

and f denotes the fermions of same charge (e, u and d type).This is a very good ansatz for the fermion mass matrix and its diagonalization leads naturally to the Kobayashi-Maskawa matrix.

Many attempts have been made to reproduce Eq. (17), and in some cases they call for fine-tuning of coupling constants[22], therefore we shall adopt an unconventional fermionic assignment in such a way that the leading order approximation to (17) will emerge solely for symmetry reasons. The standard left-handed fermions will transform under $SU(3)_H$ as triplets and the right-handed ones as antitriplets[25] (this does not imply a direct product of $SU(3)_H$ with SU(9)), consequently the allowed condensate("Higgs") representations are 3 and $\bar{6}$, and such condensates can be formed if technifermions also appear with the same assignments. Bellow we display the caracteristics of our "composite Higgs" system[F2]:

	CONDENSATE SCALES (singlets of $SU(4)_{TC} \times SU(3)_c$)	$SU(3)_H$ CONDENSATE REP.	π (π) DECAY CONSTANT
$SU(4)_{TC}$	$\langle \bar{\psi}_{TC} \psi_{TC} \rangle \sim -\mu^3_{TC} \sim -(300\ \text{Gev})^3$	$\bar{6}$	$F_\pi \sim 125$ Gev
$SU(3)_c$	$\langle \bar{\psi}_c \psi_c \rangle \sim -\mu^3_c \sim -(0.3\ \text{Gev})^3$	3	$f_\pi \sim 95$ Mev

If we consider only the TC composite Higgs, the sextet when written as a 3x3 matrix is not traceless, and we are naturally led to the following mass matrix[25]

$$m_f = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & c_f(\alpha_i \mu_{TC}) \end{pmatrix} \qquad (18)$$

Only the third generation has gotten mass. At this stage there is one mass splitting among the quarks b , t and the lepton τ . The τ receives mass at the SU(9) level, and from Eq.(10) with μ_{TC} given above, $M_W \sim 10^{15}$ Gev and $\alpha_W \sim 1/40$ we get $m_\tau \sim O(1-10)$ Gev . The b and t quarks shall be heavier than the τ lepton, and a renormalization factor of O(3) is usually expected as in the common scheme of SU(5) unification. Effects of different charges will also split b and τ masses (though a more complex mechanism of isospin breaking might be necessary[26]), and some complexity could advent of a different pattern of symmetry breaking. The value of α_W in Eq.(10) is important to determine m_f , suppose that instead of the breaking $SU(9) \rightarrow SU(4)_{TC} \times SU(3)_c \times SU(2)_L \times U(1)$ we have the one depicted in Fig. 2, then the diagram of Fig.3 will dominate the quark masses (due the larger coupling constant and also a slight increase coming from the term between brackets in (10)) .

When QCD is turned on the structure (18) cannot remain intact any longer, and the horizontal triplet will modify Eq.(18) to

$$m_f = \begin{pmatrix} 0 & c^f \alpha_i \mu_c & 0 \\ -c^f \alpha_i \mu_c & 0 & 0 \\ 0 & 0 & c^f \alpha_i \mu_{TC} \end{pmatrix} \qquad (19)$$

Our composite Higgs system, sextet (Φ_{TC}) and triplet (ϕ_c) of $SU(3)_H$, will interact with gauge and Higgs horizontal bosons at higher orders, and we can expect the complete breaking of the left over symmetry in (19) through radiative corrections, entailing the following structure of VEVs for the effective potential $V(\Phi_{TC},\phi_c)$:

$$V(\Phi_{TC}, \phi_c) \approx \begin{pmatrix} 0 & \mu_c & 0 \\ -\mu_c & 0 & \epsilon \\ 0 & -\epsilon & \mu_{TC} \end{pmatrix} . \qquad (20)$$

The above scheme has already been discussed by Wilczek and Zee[22] for fundamental scalars.Here the Higgs are condensates (at scales μ_c and μ_{TC}), and they align in the $SU(3)_H$ space generating one intermediate scale ϵ , perpendicular to the directions of $\langle \Phi_{TC} \rangle$ and $\langle \phi_c \rangle$. As a consequence (19) goes to the form (17), and generally we might find that all zeros in (19) are filled at higher orders in α_i .

Despite the difficulties to deal with the little known composite Higgs, the picture envisaged here is very interesting; it explains the order of magnitude of lighter and heavier generations, only one set of (effective) Higgs bosons is responsible for lepton and quark masses, FCNC problems are absent and the extra requisites are those (very common) of unification and horizontal symmetries.

5 TECHNIFERMION AND PSEUDO-GOLDSTONE BOSON MASSES

In earlier works[2,9] an analysis of pseudo-Goldstone (π^i) masses based in the asymptotic behavior of fermion self-energies was performed, however it seems that the pseudo-Goldstone bosons known in Nature (π's, K's,...) do not receive large mass contribution from short distances (as exemplified by the $\pi^+ - \pi^0$ mass diference[27]), moreover the quark masses play a significant role in that spectrum. Following this reasoning we verify that in the present scheme (and in the one of references (3)) the π^i masses will be dominated by the "tadpole terms"[27], i.e. the current technifermion masses cannot be neglected and technipions, etc..., will get masses in the same way as the QCD Goldstone bosons.

Technifermions acquire masses through the diagram of Fig. 1. To exemplify, notice that in the model discussed in section 3 ($SU(7)$[20]), the charged technifermions U, D and E obtain electroweak masses, and the neutral one, N, acquire mass interacting with the boson called b' in ref. (20) (all of them acquire mass at $SU(7)$ level). Similarly, for the model of previous section we obtain

$$m_{TC} \sim c^{TC}\, \alpha_{GUT}\, \mu_{TC} \sim O(1\text{-}10)\ \text{Gev} \qquad . \tag{21}$$

To compute π^i masses we use PCAC, defining their propagators by

$$\tilde{\Psi}^i(q^2) = \frac{2\, F^2_{\pi i}\, m^2_{\pi i}}{-q^2 + m^2_{\pi i}} \qquad , \tag{22}$$

where we neglected the continuum contribution. $\tilde{\Psi}^i(q^2)$ can also be written as the propagators of the covariant divergences of the hadronic currents (A^i_μ (x))

$$\tilde{\psi}^i(q^2) = i \int d^4 x \, e^{iqx} \langle 0|T D^\mu A^i_\mu (x) \, D^\nu A^i_\nu (x)|0\rangle \qquad (23)$$

The divergences are in respect to all the weaker interactions contributing to the hadronic current with the quantum number of π^i .
At q=0 (22) and (23) give the relation[27]

$$2 \, F^2_{\pi^i} m^2_{\pi^i} = \sum_i m^i_{TC} \langle \bar{\psi}^i_{TC} \psi^i_{TC} \rangle + \text{weak corrections} \quad . (24)$$

With the masses estimated in Eq. (21) the first term in the right-hand side of (24) dominates, leading to[F3]

$$m_{\pi^i} \sim O(30\text{-}90) \text{ Gev} \qquad (25)$$

Masses of that order will produce a very interesting phenomenology at the weak (W and Z) bosons scale.

6. CONCLUSIONS

We have argued, based in recent non-perturbative studies, that the dynamically generated fermionic self-energy has a hard behavior, i.e. is given by the usually called irregular solution. Its phenomenology has been analyzed and a technicolor model proposed.

We have built a SU(9) theory containing a $SU(4)_{TC}$ subgroup and a SU(5) Georgi and Glashow model with only one family. The family replication is introduced by a horizontal SU(3) symmetry, with a fermionic assignment leading to the Fritzsch matrix, generated uniquely by symmetry reasons. The QCD and TC theories play a major role in this scheme, their alignment under $SU(3)_H$ provide all fermionic spectrum, and the reason of its extremes (m_e to m_t) appear naturally. Technicolor breaks the $SU(2)_L \times U(1)$ symmetry, and fundamental scalars are responsible by the SU(9) and $SU(3)_H$ breaking at grand unification scale, consequently the problems of hierarchy and FCNC are absent.

The model has small sensibility to the physics of very high energies, and bears some similarity to the interesting line proposed in ref. (3), when not only TC but also QCD has a non-trivial UV fixed point.

The difficulty of our approach is the usual one, i.e. the lack of a better knowledge of the strongly interacting Higgs system does not permit a precise determination of the mass spectrum, actually the long discussion[4-9] of which one is the correct behavior of the self-energy reflects another aspect of this problem, however under a series of hypothesis the phenomenology of this model can be deeply investigated, and we hope to return to this study in the future.

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FOOTNOTES

1. A reasonable agreement has also been obtained for QED between the result of ref. (8) ($e_c^2/4\pi \gtrsim 0.25$) and the numerical one of Bartholomew et al[13].

2. Notice that the condensates transformations under $SU(3)_H$ have been assumed arbitrarily, other cases, both triplets or sextets, are possible, however a conclusive answer would only come out with a complete analysis of the $SU(4)_{TC}$ and $SU(3)_C$ alignment under the perturbation of $SU(3)_H$, which shall be discussed elsewhere. Whatever the case, the composite Higgs structure is sufficient to generate the complete breaking of the global SU(3) symmetry of the fermionic mass matrix.

3. We consider Eq.(25) as a lower bound for pseudo-Goldstone masses, usually charged bosons get heavier masses given by the weak corrections of Eq.(24).

FIGURE CAPTIONS

Fig. 1- Diagram that generates current fermion masses.

Fig. 2- Altenative breaking of SU(9), which could be responsible for an increasing of quark masses in relation to the leptonic ones.

Fig. 3- Diagram connecting quarks to techniquarks through an interaction stronger than SU(9).

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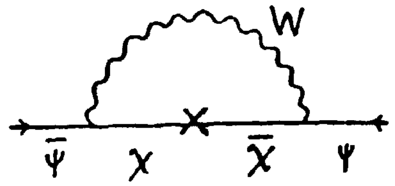

FIGURE 1

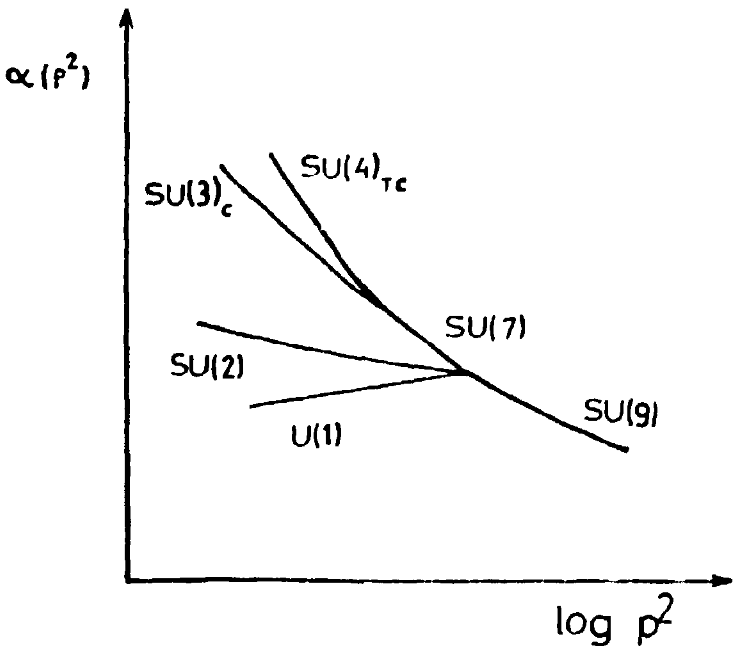

FIGURE 2

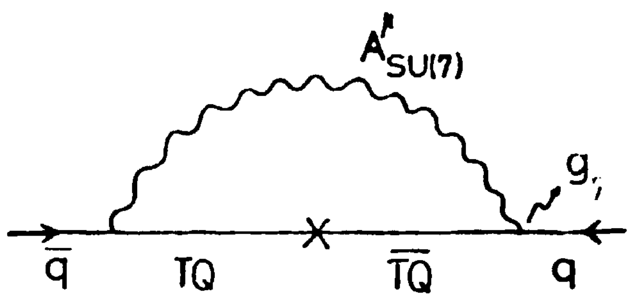

FIGURE 3