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HIGH-COLOUR AND MASS HIERARCHIES

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ABSTRACT

We present a model based on the gauge group $G = G_{HC} \times G_S \times SU(2)_L \times U(1)$, where the hypercolour gauge group G_{HC} is responsible for the dynamical breaking of the strong group G_S to $SU(3)_C$ of QCD. Chiral symmetry breaking of high colour representations produces dynamical breaking of the electroweak $SU(2)_L \times U(1)$ gauge group. Fermion masses and flavour mixing are dynamically generated from the condensations of high colour representations. A phenomenological analysis of the model is also presented.

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In this letter, we shall present a study of a fully dynamical scenario for the breaking of the electroweak gauge symmetry¹⁾ in which, contrary to the technicolour scheme²⁾, one only needs ordinary QCD forces with new quarks in high colour representations. According to this scenario³⁾⁻⁷⁾, chiral symmetry breaking with quarks in high colour representations of the strong group may occur at much larger mass scales than ordinary chiral symmetry breaking and could produce the dynamical breaking of the electroweak gauge symmetry. Such a scenario can be justified if one assumes a) that chiral symmetry breaking depends entirely on the strength of the quark-antiquark binding potential, which is based on the one-gluon exchange approximation and takes common value for all representations, i.e.,

$$C_2(R)\alpha_s = \text{const.} \quad (1)$$

and b) that perturbation theory holds in the region of chiral symmetry breaking.

Such a dynamical scheme is strongly supported by recent Monte Carlo simulations on the scales of chiral symmetry breaking⁸⁾ on a lattice. It was found that in an SU(2) and SU(3) colour model, the fermions in different colour representations do condense at exponentially separated mass scales. Other non-perturbative methods also support the validity of Eq. (1)⁹⁾.

To generate fermion masses, one has to enlarge the colour gauge group and subsequently to break it down to SU(3)_C⁴⁾⁻⁶⁾. Then, if the fermions are assigned to sufficiently high representations of the enlarged group, then under the SU(3)_C, they decompose to ordinary fermions plus new fermions with exotic colour quantum numbers. Therefore, when the new quarks condense, in addition to electroweak symmetry breaking, they also provide masses to fermions via the exchanges of the heavy sector bosons as in Fig. 1.

As we have described previously⁶⁾, if the strong group is SU(4)_S or SU(5)_S, then it is not possible to describe in this framework the variety of fermion masses and their mixings, in a realistic model. If the strong group is [SU(3)×SU(3)]_S, there was a hope for a construction of a realistic model. However, a detailed analysis of a three-family model based on the group [SU(3)×SU(3)]_S×SU(2)_L×U(1) showed that perturbation theory does not hold in a region in which we would like to obtain phenomenological predictions. Thus, we were led to consider an [(SU(3)×SU(3)×SU(3))_S×SU(2)_L×U(1)] model in which the perturbation theory is safe and a reasonable fermion mass spectrum emerges.

As we have mentioned in Ref. 6), if one stays with the naive "Casimir scaling" of Eq. (1), then even for small dimensionality representations of $SU(3)_C$, the predicted scale at which chiral symmetry breaking occurs exceeds the Planck scale. We assumed that the "Casimir scaling" does not work with full strength but there exists some "screening of the Casimir scaling". Such an effect has been seen in recent lattice calculations with an $SU(3)$ -colour gauge group⁸⁾.

We parametrized this effect, associating with every representation of $SU(3)_C$ a screening factor $f(R)$, and then Eq. (1) is replaced by

$$f^2(R) C_2(R) d_S(\mu_R) = \text{const.} \quad (2)$$

where μ_R is the scale where the chiral symmetry breaking will take place. These screening factors are calculated by fitting the fermion mass spectrum and mixing angles. These values of $f(R)$ are the predictions of our model.

In all of our previous attempts, because of a lack of a dynamical mechanism, the breaking of G_S to $SU(3)_C$ was done by introducing elementary Higgs fields. Here we shall employ another suggestion¹⁰⁾. Let us consider a gauge theory whose strong sector is governed by the gauge group $G = G_{HC} \times G_S$, where $G_{HC} = SU(N)_{HC}$ is a hypercolour strong gauge group, whose characteristic scale Λ_{HC} is much larger than the corresponding one Λ_{QCD} of QCD and $G_S = [SU(3) \times SU(3)]_S$. Let us introduce hyperfermions according to

$$G = SU(N)_{HC} \times (SU(3) \times SU(3))_S$$

$$HF_L \sim (N ; (\bar{3}, 1)) \quad (3)$$

$$HF_R \sim (N ; (1, \bar{3}))$$

Our theory has a chiral global symmetry

$$SU(3)_L \times SU(3)_R \quad (4)$$

consisting of the independent transformations of the left- and right-handed parts of the hyperfermions. This chiral symmetry is broken by the condensates

$\langle \overline{HF}_L \cdot HF_R \rangle \neq 0$ to the diagonal $SU(3)_V$. However, this breaking forces also the gauge $[SU(3) \times SU(3)]_S$ to break into the diagonal $SU(3)$ which is identified to be the colour gauge group¹¹⁾. This is because the introduction of the gauge group G_S can be considered as gauging the global symmetry (4). The axial parts of the $SU(3)_L \times SU(3)_R$ current can create, from the vacuum, hyperpions with decay constant $F_{H\pi}$. When these currents are identified with the corresponding currents of G_S , they still produce hyperpions but then they are coupled to the $G_S/SU(3)_{diag}$ gauge bosons which become massive.

Our fermion mass generation mechanism suggests that the horizontal splitting between different families should emerge simply from the fact that the various fermion flavours and the corresponding echo-fermions (new fermions with exotic colour) come from the decomposition under $SU(3)_C$ of haplons (elementary fermions) which belong to different representations under $[SU(3) \times SU(3)]_S$. On the other hand, the vertical mass splitting between fermions of the same $SU(2)_L$ -weak multiplet should be due to different representations under $[SU(3) \times SU(3)]_S$ of the right handed haplon multiplets which are $SU(2)_L$ singlets. One must be very careful at this point. Assigning right handed haplons to different representations, one is introducing strong $[SU(3) \times SU(3)]_S$ anomalies, since the model is not any more vector-like in this sector. Nevertheless, it is always possible to cancel such strong anomalies. The family mixing appears because some of the echo-quarks of one chirality happens to be attracted according to Eq. (2) with equal strength from echo-quarks of opposite chirality which belong to different families. Thus, they produce condensates with all of them, while the radiative mechanism of fermion mass generation produces also off-diagonal elements in addition to the diagonal ones which lead to fermion mixing.

Before we present our model, let us discuss the general strategy for the cancellation of the anomalies. There are no anomalies in the hypercolour sector because this sector is vector-like. Then we choose the representations under which the hyperfermions transform, in such a way that the two $SU(3)$'s in G are anomaly free. Then we add multiplets of the form $(\underline{1}, \underline{1}, R)$ under G_S in such a way that the anomalies of the third $SU(3)$ group cancel. In this way all strong anomalies cancel by construction. Moreover, the additional particles are the missing partners of those echo-fermions which cannot condense, since the corresponding opposite chirality echo-fermions belong to a different representation. The most important feature of these additional particles is that when they condense they disappear from the low energy fermionic spectrum because they belong to high $SU(3)_C$ representations. For the cancellation of the anomalies in

the electroweak $SU(2)_L \times U(1)$ sector the condition

$$\sum_{\text{all repr.}} Q_{e.l} = 0 \quad (5)$$

must be satisfied. One way to satisfy the above condition is to look for quark and lepton representations which have exactly the right multiplicity to meet Eq. (5). We know that in the standard model, there exists an easy solution to this problem. In the present context, to find such a solution is very difficult. But there is another way to cancel the electroweak anomalies. We add multiplets of the following type

$$\begin{aligned} & (SU(3) \times SU(3) \times SU(3))_S \times SU(2)_L \\ & ((1, 1, R), 2)_L \\ & ((1, 1, R), 1)_R \\ & ((1, 1, R), 1)_R \end{aligned} \quad (6)$$

where R is a real representation of $SU(3)$ in order of not having strong anomalies, and we fix the hypercharge in such way as to cancel the electroweak anomalies. These new multiplets, when they condense, disappear from the physical fermion spectrum, since they belong to high colour representations.

Let us now construct a model according to the above anomaly cancellation discussion. The model is based on the gauge group $G = SU(6)_{HC} \times [SU(3) \times SU(3) \times SU(3)]_S \times SU(2)_L \times U(1)$. We assign the hyperfermions and elementary haplons to the following representations under G :

$$\begin{aligned} HF_{1L} & \sim (21 ; (\bar{3}, 1, 1) ; 1 ; 0) \\ HF_{1R} & \sim (21 ; (1, \bar{3}, 1) ; 1 ; 0) \\ HF_{2L} & \sim (6 ; (1, \bar{3}, 1) ; 1 ; 0) \\ HF_{2R} & \sim (6 ; (1, 1, \bar{3}) ; 1 ; 0) \\ \begin{pmatrix} H^{u,c} \\ H^{d,s} \end{pmatrix}_L & \sim (1 ; (3, 3, \bar{3}) ; 2 ; 1/3) \end{aligned} \quad (7)$$

$$\begin{aligned}
 H_R^u (H_R^d) &\sim (1 ; (\bar{3}, \bar{3}, 1) ; 1 ; 4/3 (-2/3)) \\
 H_R^c &\sim (1 ; (1, 6, \bar{3}) ; 1 ; 4/3) \\
 H_R^s &\sim (1 ; (1, \bar{3}, \bar{3}) ; 1 ; -2/3) \\
 \begin{pmatrix} H^{\nu e} \\ H^e \end{pmatrix}_L &\sim (1 ; (\bar{3}, 3, 1) ; 2 ; -1) \\
 H_R^e &\sim (1 ; (\bar{3}, 3, 1) ; 1 ; -2) \\
 \begin{pmatrix} H^{\nu \mu} \\ H^\mu \end{pmatrix}_L &\sim (1 ; (1, \bar{3}, 3) ; 2 ; -1) \\
 H_R^\mu &\sim (1 ; (1, \bar{3}, 3) ; 1 ; -2) \\
 E_L^1 &\sim (1 ; (1, \bar{6}, 1) ; 1 ; -2/3) \\
 E_R^2 &\sim (1 ; (1, 15, 1) ; 1 ; -2/3)
 \end{aligned} \tag{7}$$

cont.

where, $HF_{1,2}$ refer to the hyperfermions, H^f are the haplons which after the decomposition under $SU(3)_C$ give the fermions, while $E^{1,2}$ are new multiplets required for the cancellation of the strong anomalies.

According to our previous discussion, when the condensations $\langle HF_{1L} \cdot \overline{HF}_{1R} \rangle$ and $\langle HF_{2L} \cdot \overline{HF}_{2R} \rangle$ form, the strong group G_S breaks to

$$G_S = SU(3) \times SU(3) \times SU(3) \xrightarrow{\text{diag.}} SU(3) \times SU(3) \longrightarrow SU(3)_C \tag{8}$$

and if A_μ, B_μ, C_μ are the gauge bosons of the three $SU(3)$'s gauge groups, then after diagonalization of the gauge boson mass matrix, the physical massive gauge bosons are the

$$V_{1\mu} = \frac{1}{\sqrt{2}} (A_\mu - B_\mu) \quad , \quad V_{2\mu} = -\frac{1}{\sqrt{6}} (A_\mu + B_\mu) + \sqrt{\frac{2}{3}} C_\mu$$

with mass-squared m_1^2, m_2^2 respectively, while the orthogonal linear combination of the previous ones, $-1/\sqrt{3}(A_\mu + B_\mu + C_\mu)$ remain massless and we identify them with the massless gluons of QCD.

We can easily identify the ordinary quarks and leptons and the echo-fermions, if we decompose the fermionic haplons of Eq. (7) under the diagonal $SU(3)_C$

$$\begin{aligned}
 SU(3) \times SU(3) \times SU(3) &\longrightarrow SU(3) \times SU(3) \longrightarrow SU(3)_c \\
 (3, 3, \bar{3}) &\longrightarrow (\bar{3}, \bar{3}) + (6, \bar{3}) \longrightarrow 3 + \bar{6} + 3 + 15 \\
 (\bar{3}, \bar{3}, 1) &\longrightarrow (3, 1) + (\bar{6}, 1) \longrightarrow 3 + \bar{6} \\
 (1, 6, \bar{3}) &\longrightarrow (6, \bar{3}) \longrightarrow 3 + 15 \\
 (\bar{3}, 3, 1) &\longrightarrow (1, 1) + (8, 1) \longrightarrow 1 + 8
 \end{aligned} \tag{9}$$

We assume that the condensation $\langle 15_L \cdot 15_R \rangle$ of the highest echo-quark representation is responsible for the dynamical breaking of $SU(2)_L \times U(1)$. Then its condensation scale should be 0(500) GeV and the associated screening factor is given in Table 1.

The ordinary fermions obtain radiative masses from the echo-fermions via the exchange of the heavy vector bosons $V_{1\mu}, V_{2\mu}$ (Fig. 1). According to the assignments of haplons in Eq. (7) and the decompositions (9), the ordinary quarks obtain the following radiative (Dirac) mass contributions

$$\mathcal{M}^u = \begin{pmatrix} C(6) \frac{\mu_6^3}{m_1^2} & C(15) \frac{\mu_{15}^3}{m_1^2} \\ C(6) \frac{\mu_6^3}{m_1^2} & C(15) \frac{\mu_{15}^3}{m_2^2} \end{pmatrix} + h.c. \tag{10}$$

$$\mathcal{M}^d = \begin{pmatrix} C(6) \frac{\mu_6^3}{m_1^2} & C(6) \frac{\mu_6^3}{m_1^2} \\ C(6) \frac{\mu_6^3}{m_1^2} & C(6) \frac{\mu_6^3}{m_2^2} \end{pmatrix} + h.c. \tag{11}$$

where

$$C(R) = \frac{9d_s}{4\pi} C_2(R) \ln \left(\frac{m_{1(2)}^2}{(3C_2(R))^{2/3} \mu_R^2} \right) \quad (12)$$

with $C_2(R)$ the quadratic Casimir of the R representation, α_s the strong coupling constant and μ_R the scale where the chiral symmetry breaking occurs in the R sector.

These mass matrices after diagonalization give

$$\begin{aligned} m_\mu &\simeq 1 \text{ MeV} \\ m_d &\simeq 1 \text{ MeV} \\ m_c &\simeq 1.5 \text{ GeV} \\ m_s &\simeq 175 \text{ MeV} \end{aligned} \quad (13)$$

where we have taken

$$\begin{aligned} m_1 &= 270 \text{ TeV} \\ m_2 &= 14 \text{ TeV} \\ \mu_6 &= 260 \text{ GeV} \\ \mu_{15} &= 500 \text{ GeV} \\ d_s &= 0.10 \end{aligned} \quad (14)$$

The Cabibbo angle is given by $\tan \theta_c \sim m_d/m_s$. For the leptons we have

$$\mathcal{U}^E = \begin{pmatrix} C(8) \frac{\mu_9^3}{m_1^2} & C(8) \frac{\mu_9^3}{m_1^2} \\ C(8) \frac{\mu_8^3}{m_1^2} & C(8) \frac{\mu_8^3}{m_2^2} \end{pmatrix} + h.c. \quad (15)$$

from where we get

$$m_e \simeq 0.5 \text{ MeV} \quad (16)$$

$$m_\mu \simeq 105 \text{ MeV}$$

with

$$\mu_g = 220 \text{ GeV} \quad (17)$$

The fitting of the condensation scales μ_R leads to the screening factors listed in Table 1.

We turn now to the study of some phenomenological aspects of the model, namely the spectrum of pseudo-Goldstone bosons and the problems of flavour changing neutral currents. We analyze first the quark sector. According to the fermion assignment [Eq (7)], the quark fields under $SU(3)_C$ are:

$$\begin{aligned} \psi_{(3,3,\bar{3})_L}^i &= \psi_{3_L}^{i,1} + \psi_{\bar{6}_L}^{i,1} + \psi_{3_L}^{i,2} + \psi_{15_L}^{i,1} \\ \psi_{(\bar{3},\bar{3},1)_R}^1 &= \psi_{3_R}^{1,1} + \psi_{\bar{6}_R}^{1,1} \\ \psi_{(\bar{3},\bar{3},1)_R}^2 &= \psi_{3_R}^{2,1} + \psi_{\bar{6}_R}^{2,1} \\ \psi_{(1,6,\bar{3})_R}^1 &= \psi_{3_R}^{1,2} + \psi_{15_R}^{1,1} \\ \psi_{(1,\bar{3},\bar{3})_R}^2 &= \psi_{3_R}^{2,2} + \psi_{\bar{6}_R}^{2,3} \\ \psi_{(1,15,1)_R}^2 &= \psi_{15_R}^{2,1} \\ \psi_{(1,\bar{6},1)_L}^1 &= \psi_{\bar{6}_L}^{1,3} \end{aligned} \quad (18)$$

where the first index is the isospin index, while the second index counts the number of the fields. Then, the global symmetries in this sector (neglecting the triplets) are:

$$G_{gl} = SU(3)_L^{\psi_{\bar{6}L}} \times SU(3)_R^{\psi_{\bar{6}R}} \times U(1)_L^{\psi_{\bar{6}L}} \times U(1)_R^{\psi_{\bar{6}R}} \\ \times SU(2)_L^{\psi_{15L}} \times SU(2)_R^{\psi_{15R}} \times U(1)_L^{\psi_{15L}} \times U(1)_R^{\psi_{15R}} \quad (19)$$

The condensates $\langle \bar{\psi}_{\bar{6}L} \times \psi_{\bar{6}R} \rangle$ and $\langle \bar{\psi}_{15L} \times \psi_{15R} \rangle$ break the global symmetries G_{gl} to G'_{gl}

$$G'_{gl} = SU(3)_{L+R}^{\psi_{\bar{6}}} \times U(1)_{L+R}^{\psi_{\bar{6}}} \\ \times SU(2)_{L+R}^{\psi_{15}} \times U(1)_{L+R}^{\psi_{15}} \quad (20)$$

The Goldstone bosons in G_{gl}/G'_{gl} are 12. A certain linear combination of them are the three Goldstone bosons which are eaten by the $W^\pm$, Z^0 of the electroweak interactions. One would expect that some of the remaining Goldstone bosons become massive pseudo-Goldstone bosons as the result of the explicit breaking of G_{gl} by the broken $SU(3) \times SU(3)$ gauge interactions. In this model, this not possible because the broken generators of $SU(3) \times SU(3)$ do not provide couplings between $\psi_{\bar{6}}$ and ψ_{15} . The only source which can give mass to the charged Goldstone bosons in the first order of perturbation theory is the weak and the electromagnetic interactions. According to Ref. 12), the charged Goldstone bosons get a mass

$$m_{\pi_6^\pm} \simeq \sqrt{\alpha_{em}} \frac{4}{\sqrt{3}} \mu_6 \simeq 51 \text{ GeV} \quad (21)$$

while the other neutral Goldstone bosons remain massless. The $\pi_6^\pm$ pseudo-Goldstone bosons should be seen, in principle, in e^+e^- machines while the π_6^0 , π_6^0 true Goldstone bosons behave like unwanted axions.

Phenomenologically, the big advantage of the high colour scenario and therefore of the present model is the lack of tree order flavour changing neutral currents due to gauge boson exchange which was the main problem that ruled out the technicolour scenario. However, flavour changing neutral currents are induced by high order exchanges. The most dangerous is the contribution to the kaon mass difference due to the box diagram¹³⁾. From the kaon mass difference, we easily obtain a lower bound of $m_2 > 1.3 \text{ TeV}$ for the mass of the lighter vector boson. Thus, according to Eq. (14) the model satisfies this bound.

In the lepton sector we have

$$\begin{aligned}
 \psi_{(\bar{3}, 3, 1)_L}^i &= \psi_{1L}^{i,1} + \psi_{8L}^{i,1} \\
 \psi_{(\bar{3}, 3, 1)_R}^1 &= \psi_{1R}^{1,1} + \psi_{8R}^{2,1} \\
 \psi_{(1, \bar{3}, 3)_L}^i &= \psi_{1L}^{i,2} + \psi_{8L}^{i,2} \\
 \psi_{(1, \bar{3}, 3)_R}^1 &= \psi_{1R}^{1,2} + \psi_{8R}^{2,2}
 \end{aligned} \tag{22}$$

Because 8 is a real representation to find the global symmetries in this sector, we simply count how many 8's we have. Therefore, the global symmetry¹⁴⁾ is

$$G_{gl.} = SU(6) \tag{23}$$

The condensations $\langle \bar{\psi}_{8L} \psi_{8R} \rangle$ break this symmetry to

$$G'_{gl.} = SO(6) \tag{24}$$

and we have 20 massless Goldstone bosons. These are colour-singlet objects and the charged ones will get a mass from the electroweak interactions

$$m_{\pi_8^\pm} \simeq \sqrt{d_{em}} \frac{4}{\sqrt{3}} \mu_8 \simeq 43.5 \text{ GeV} \tag{25}$$

while the neutral remain massless and they behave like Majorons¹⁵⁾.

In conclusion, we discussed a model for dynamical symmetry breaking of the electroweak gauge group, with new quarks in high QCD representations. Our model is based on the group $G_{HC} \times G_S \times SU(2)_L \times U(1)$. In this model the hypercolour forces are responsible for the dynamical breaking of G_S to $SU(3)_C$ under which the elementary fermions (haplons) decompose to ordinary fermions and new ones with exotic quantum numbers. The condensation of the new fermions leads to the electroweak symmetry breaking while it generates fermion masses and mixings. It gives definite predictions on the various condensation scales which could be tested in future Monte Carlo calculations. Phenomenologically our model has certain virtues, since it avoids problems with flavour changing neutral currents, but on the other hand, in its particle spectrum there are unwanted axions and Majorons. The advantage of the present model compared to our previous model building attempts is that it is more economical in the introduction of high colour representations. This, in turn, means that we increase the "plateau" where we can use perturbation theory, up to TeV region where we do our calculations.

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$f^2(6)$	$f^2(8)$	$f^2(15)$
0.46	0.51	0.23

Table 1: the screening factors of the 6, 8 and 15 colour representations

NOTE ADDED

After this paper was typed, it came to our knowledge a new paper by J. Kogut, J. Polonyi, H.W. Wyld and D.K. Sinclair on "Hierarchical mass scales in lattice gauge theories with dynamical, light fermions", (Illinois preprint ILL-TH-85-8), in which the disparity in mass scales that was observed in their early simulations of the quenched approximation to lattice gauge theory⁸⁾, appears to survive the inclusion of fermion loops. Not only that, but they found a very large number for the ratio of the temperature of the $\ell = \frac{1}{2}$ and $\ell = 1$ colour representations of SU(2), namely $T_1/T_{\frac{1}{2}} = 175 \pm 50$.

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FIGURE CAPTION

Fig. 1: Diagram producing a radiative mass to an ordinary fermion f . A_μ denotes the massive gauge bosons of $G_S/SU(3)_C$. $Q(R)$ is a R-plet of exotic colour fermion and $\Sigma_R = \langle Q(R)_L \bar{Q}(R)_R \rangle$.