Ref.TH.3586-CERN

FERMION MASS GENERATION AND ELECTROWEAK
SYMMETRY BREAKING FROM COLOUR FORCES

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A B S T R A C T

The colour gauge group is extended to $SU(3) \otimes SU(3)$ and is subsequently broken to diagonal $SU(3)_c$. Under the diagonal $SU(3)_c$ the fundamental fermionic constituents of the larger strong group become ordinary quarks plus new quarks with exotic quantum numbers. Chiral symmetry breaking in the exotic quark sector may occur at much larger mass scales than ordinary chiral symmetry breaking, and could produce dynamical breaking of electroweak gauge symmetry and radiative masses for the light fermions.

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The standard electroweak model¹⁾ is in excellent agreement with existing experimental results. However, although it does predict and explain the interactions of the fermions involved, it does not predict their masses and mixing angles.

The breaking of $SU(2)_L \otimes U(1)$ invariance is caused by the Higgs fields, which have to develop a vev of 250 GeV in order to produce enough mass for the vector boson W . Fermions also acquire masses from the spontaneous symmetry breaking via Yukawa interactions with the same Higgs fields. Thus, even if the fermion masses are free parameters and can be put in by hand in the theory, one has to choose unnaturally small Yukawa coupling constants in order to fit their measured values.

An attractive alternative to the symmetry breaking and generation of masses by the Higgs mechanism is the idea that $SU(2)_L \otimes U(1)$ is broken dynamically²⁾ by a system of massless fermions, the technifermions, which in addition to $SU(3)_C \otimes SU(2)_L \otimes U(1)$, feel a new QCD-like interaction, the technicolour interaction, whose scale Λ_{TC} is of the order of 1 TeV. The ordinary fermions in this new picture obtain masses through their interactions with new heavy vector bosons, which are able to couple them to the technifermions, according to the figure³⁾. Thus, in the full technicolour scenario, the technicolour gauge group G_{TC} has to be enlarged to an extended technicolour group G_{ETC} ⁴⁾, under which the standard quarks and leptons and the technifermions transform non-trivially. For the current quark and lepton masses, one finds from the diagram in the figure

$$m_f \sim g_{ETC}^2 \frac{\mu_{TC}^3}{\mu_{ETC}^2} \quad (1)$$

where $\mu_{TC} \sim O(\Lambda_{TC})$ is fixed by the W gauge boson masses and μ_{ETC} is the mass of a heavy ETC gauge boson which couples fermions and technifermions. For $g_{ETC}^2 \sim O(1)$, $\mu_{TC} \sim 300$ GeV and $\mu_{ETC} \sim 2 - 60$ TeV, the known quark and lepton masses can be obtained. Unfortunately, no successful model exists as yet which realizes the above scenario. The main problem⁵⁾ is that in addition to the ETC gauge bosons, there exist heavy neutral horizontal ETC gauge bosons with similar masses which couple ordinary quarks of the same charge to each other, producing in this way flavour-changing neutral currents and predicting excessive rates for flavour-changing neutral current processes.

Here we are going to discuss another proposal^{6),7)} in the direction of dynamical symmetry breaking which is much more attractive than the technicolour scheme, since it requires only ordinary QCD with additional fermions in higher than triplet representations of the $SU(3)_C$ gauge group. Our version of this scenario

involves a larger $SU(3) \otimes SU(3)$ gauge symmetry which breaks down to ordinary colour $SU(3)$. The fundamental fermions, the haplons, of $SU(3) \otimes SU(3)$ produce new fermions which belong to different representations of the $SU(3)_C$ and have to be identified with the ordinary fermions and new echo-fermions^{*}). The most useful aspect of our scheme is that we are able to give masses to ordinary fermions from the echo-fermions, which was an unsolved problem in the original proposals. Moreover, the great difficulty which ruled out the ETC model, the violation of the bounds for flavour-changing neutral currents, does not exist at all here.

Let us look at the basis for the speculations^{6),7)} that a new very strong interaction like technicolour is unnecessary for dynamical breaking of $SU(2)_L \otimes U(1)$, and that instead QCD with new quarks in representations $R > 3$ can play the same rôle. The strength of the quark-antiquark binding potential based on the approximation of one gluon exchange is proportional to

$$C_2(R) \alpha_s(\mu) \quad (2)$$

where $C_2(R)$ is the quadratic $SU(3)_C$ Casimir invariant of the R representation¹⁰⁾. The quantity in Eq. (2) is a measure of the effective colour flux linking a quark-antiquark pair at mass scale μ . Thus, quarks belonging to representations $R > 3$ which have bigger Casimir invariants than triplets, carry more colour charge and hence interact more strongly than ordinary quarks. The weak coupling regime for ordinary triplet quarks may be a strong coupling regime for quarks in $R > 3$. If chiral symmetry breaking in the triplet sector occurs at $\alpha_s(\mu) = \alpha_s(\mu_3)$ and in the R -plet sector at $\alpha_s(\mu) = \alpha_s(\mu_R)$, then the dynamical mass scales of these two sectors are related by

$$\frac{\mu_R}{\mu_3} = \exp \left[\int_{\alpha_s(\mu_3)}^{\alpha_s(\mu_R)} \frac{d\alpha_s}{\beta(\alpha_s)} \right] \quad (3)$$

where β is the QCD β function. Thus there is an exponential enhancement of μ_R/μ_3 if $\alpha_s(\mu_R) \neq \alpha_s(\mu_3)$, and they are small enough for perturbation theory to be applicable. Of course, a better understanding of chiral symmetry breaking in QCD is needed before one can say that $\alpha_s(\mu_R) \neq \alpha_s(\mu_3)$; however, recently there has been an important hint that they are different from lattice results.

^{*}) Echo-quark phenomenology has been discussed in a different context in Ref. 8). Phenomenological studies of exotic quarks also exist in a completely different context⁹⁾. Phenomenological consequences of exotic quarks of the type of our echo-quarks have been examined in Ref. 7).

In a Monte Carlo study¹¹⁾ of the lattice $SU(2)$ colour model, it has been shown that different colour representations indeed condense at exponentially separated mass scales. Thus there exist some reasons to justify the speculations^{6),7)} that for quarks belonging to $R=6$, the chiral symmetry breaking occurs at $\mu \sim 1$ TeV. In the following, we assume that this is indeed the case.

Let us consider a model based on the gauge group $G = G_c \otimes SU(2)_L \otimes U(1)$, where the strong sector is $G_c = SU(3) \otimes SU(3)$ and the electroweak sector is the standard $SU(2)_L \otimes U(1)$. We impose a discrete symmetry interchanging the two $SU(3)$ subgroups of G_c . The elementary fermions are not identified with quarks. There are quark-haplons, belonging to $(3,3)$ of G_c , that have the following quark assignments under the $SU(2) \otimes U(1)$:

$$G = (SU(3) \otimes SU(3))_c \otimes SU(2)_L \otimes U(1)$$

$$\begin{pmatrix} H^u \\ H^d \end{pmatrix}_L = \left((3, 3) \quad ; \quad 2 \quad ; \quad -1/3 \right) \quad (4)$$

$$H_R^u = \left((3, 3) \quad ; \quad 1 \quad ; \quad -4/3 \right)$$

$$H_R^d = \left((3, 3) \quad ; \quad 1 \quad ; \quad 2/3 \right)$$

We shall discuss later the cancellation of anomalies in this model (in connection with the lepton-haplons).

Denoting the gauge bosons corresponding to the two $SU(3)$'s of G_c by A_μ^a and B_μ^a , respectively, and writing the quark-haplon field H_{ij} [with i and j referring to the two $SU(3)$'s] as a 3×3 matrix, we obtain the following strong gauge interaction for the quark-haplons:

$$g_s \text{Tr} \bar{H}^\dagger \gamma^\mu (A_\mu \lambda H + H \lambda^\dagger B_\mu) \quad (5)$$

Next we demand that the G_c breaks down spontaneously to the "diagonal" $SU(3)_c$. Of course, at this stage of the theory we would also prefer to see a dynamical symmetry breaking rather than symmetry breaking by elementary Higgs scalars. However, one could postpone the question of dynamical symmetry breaking and introduce a Higgs complex scalar field Φ belonging to (8,8) and a Ω belonging to (3,3) of G_c . Examining the potential, one finds that Higgses can develop vevs in the directions

$$\begin{aligned}\langle \Phi \rangle_0 &= \langle \Phi_A \rangle_0 + \langle \Phi_B \rangle_0 \\ \langle \Omega \rangle_0 &= \langle \Omega_A \rangle_0 + \langle \Omega_B \rangle_0\end{aligned}\tag{6}$$

where A,B refer to the two $SU(3)$ s and

$$\begin{aligned}\langle \Phi_A \rangle_0 &= \langle \Phi_B \rangle_0 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & \alpha \\ 0 & \beta & 0 \end{pmatrix} \\ \langle \Omega_A \rangle_0 &= \langle \Omega_B \rangle_0 = \begin{pmatrix} \omega \\ 0 \\ 0 \end{pmatrix}\end{aligned}\tag{7}$$

With the above vevs, a linear combination $V_\mu^- \equiv (A_\mu - B_\mu)/\sqrt{2}$ forms a colour octet of massive vector bosons, while the orthogonal combination $V_\mu^+ \equiv (A_\mu + B_\mu)/\sqrt{2}$ emerges as the colour octet of massless gauge bosons of QCD. The heavy vector bosons and their masses are:

$$\begin{aligned}V_{1,2\mu}^- &= \frac{1}{2} (A_{1\mu} \pm i A_{2\mu}) - (A_{i\mu} \rightarrow B_{i\mu}); \mu_{V_{1,2}^-}^2 = \frac{g_s^2}{4} (\alpha^2 + \beta^2 + 2\omega^2) \\ V_{4,5\mu}^- &= \frac{1}{2} (A_{4\mu} \pm i A_{5\mu}) - (A_{i\mu} \rightarrow B_{i\mu}); \mu_{V_{4,5}^-}^2 = \mu_{V_{1,2}^-}^2 \\ V_{6,7\mu}^{1-} &= \left[\cos(\epsilon/2) A_{6\mu} - \sin(\epsilon/2) A_{7\mu} \right] - [A_{i\mu} \rightarrow B_{i\mu}] \\ V_{6,7\mu}^{2-} &= \left[\sin(\epsilon/2) A_{6\mu} + \cos(\epsilon/2) A_{7\mu} \right] - [A_{i\mu} \rightarrow B_{i\mu}] \\ V_{3,8\mu}^{1-} &= \frac{1}{2} (A_{3\mu} - \sqrt{3} A_{8\mu}) - (A_{i\mu} \rightarrow B_{i\mu}); \mu_{V_{3,8}^{1-}}^2 = g_s^2 \alpha^2\end{aligned}\tag{8}$$

$$V_{3,8\mu}^{2-} = \frac{1}{2} (\sqrt{3} A_{3\mu} + A_{8\mu}) - (A_{i\mu} \rightarrow B_{i\mu}); \mu_{V_{3,8}^{2-}} = \frac{2}{3} g_s^2 \omega^2 \quad (8) \text{ contd.}$$

where g_s is the coupling constant of G_c and $\epsilon = \arg(\alpha^* \beta)$.

The quark-haplons decompose into a colour $\bar{3}$ and a colour 6 under the "diagonal" $SU(3)_c$. Naturally we identify the $\bar{3}$ with ordinary quarks q (to be precise, with the antiquarks) and the 6 with new exotic quarks which we call echo-quarks Q . Thus, for example, to the usual colour triplet u quark corresponds a colour antisextet echo $\bar{u}$ quark of the same charge, and this holds for every quark flavour. The gauge interactions of quark-haplons given in Eq. (5) can now be expressed as interactions of the ordinary and echo-quarks with the new gluon fields V_μ^+ and V_μ^- . These interactions are:

$$g_s \text{Tr } \bar{Q} \gamma^\mu \lambda V_\mu^+ Q \quad (9)$$

$$g_s \bar{q} \gamma^\mu \lambda V_\mu^+ q \quad (10)$$

$$g_s (\text{Tr } \bar{Q} \gamma^\mu \lambda V_\mu^- q + \text{Tr } \bar{q} \gamma^\mu \lambda V_\mu^- Q) \quad (11)$$

The interactions given in Eqs. (7)-(9) have this form before the symmetry breaking of $SU(3) \otimes SU(3)$ determines which combinations of the gauge fields become physical. After symmetry breaking, the massive gauge fields are given in Eq. (8) and the massless gluons V_μ^+ are just orthogonal to them. The corresponding currents coupled to heavy and massless gluons are accordingly redefined and Eqs. (9)-(11) take the form:

$$g_s \text{Tr} \left[\left(\frac{1}{\sqrt{2}} V_{1,2\mu}^+ J_{1,2}^\mu(Q, Q) + \frac{1}{\sqrt{2}} V_{4,5\mu}^+ J_{4,5}^\mu(Q, Q) \right) + h.c. + V_{6,7\mu}^{1+} J_{6,7}^\mu(Q, Q) + V_{6,7\mu}^{2+} J_{6,7}^\mu(Q, Q) + \frac{1}{2} V_{3,8\mu}^{1+} J_{3,8}^\mu(Q, Q) + \frac{1}{\sqrt{3}} V_{3,8\mu}^{2+} J_{3,8}^\mu(Q, Q) \right] \quad (9')$$

$$g_s \left[\left(\frac{1}{\sqrt{2}} V_{1,2}^{1+} J_{V_{1,2}}^\mu(q, q) + \frac{1}{\sqrt{2}} V_{4,5}^{1+} J_{V_{4,5}}^\mu(q, q) \right) + h.c. + V_{6,7}^{1+} J_{V_{6,7}}^\mu(q, q) \right. \\ \left. + V_{6,7}^{2+} J_{V_{6,7}}^\mu(q, q) + \frac{1}{2} V_{3,8}^{1+} J_{V_{3,8}}^\mu(q, q) + \frac{1}{\sqrt{3}} V_{3,8}^{2+} J_{V_{3,8}}^\mu(q, q) \right] \quad (10')$$

$$g_s \text{Tr} \left[\left(\frac{1}{\sqrt{2}} V_{1,2}^{1-} J_{V_{1,2}}^\mu(Q, q) + \frac{1}{\sqrt{2}} V_{4,5}^{1-} J_{V_{4,5}}^\mu(Q, q) \right) + h.c. + V_{6,7}^{1-} J_{V_{6,7}}^\mu(Q, q) \right. \\ \left. + V_{6,7}^{2-} J_{V_{6,7}}^\mu(Q, q) + \frac{1}{2} V_{3,8}^{1-} J_{V_{3,8}}^\mu(Q, q) + \frac{1}{\sqrt{3}} V_{3,8}^{2-} J_{V_{3,8}}^\mu(Q, q) \right] \\ + (Q \leftrightarrow q) \quad (11')$$

where

$$J_a^\mu(Q(q), q(q)) = \bar{Q}(\bar{q}) \gamma^\mu \lambda_a Q(q)$$

$$J_a^\mu(Q(q), q(Q)) = \bar{Q}(\bar{q}) \gamma^\mu \lambda_a q(Q)$$

and

$$\begin{aligned} J_{V_{1,2}}^\mu &= J_1^\mu + i J_2^\mu & J_{V_{6,7}}^\mu &= \cos(\epsilon/2) J_6^\mu - \sin(\epsilon/2) J_7^\mu \\ J_{V_{4,5}}^\mu &= J_4^\mu + i J_5^\mu & J_{V_{6,7}}^\mu &= \sin(\epsilon/2) J_6^\mu + \cos(\epsilon/2) J_7^\mu \end{aligned}$$

$$J_{V_{3,8}}^\mu = -J_3^\mu + \sqrt{3} J_8^\mu$$

$$J_{V_{3,8}}^\mu = \frac{3}{2} J_3^\mu + \frac{1}{\sqrt{3}} J_8^\mu$$

So far the haplons given in Eq. (4) have produced two triplets of quarks and two sextets of echo-quarks. When echo-quarks are massless, there exists a chiral $SU(2)_L \otimes SU(2)_R$ symmetry which is expected to break down to the diagonal $SU(2)$

when the echo-quarks form condensates at a scale $\mu_6 \sim 0(1 \text{ TeV})$. The three resulting Goldstone bosons from the spontaneous breakdown of the global chiral symmetry become longitudinal components of the $W^\pm$ and Z^0 which become massive with masses

$$\begin{aligned}\mu_W &= g F/2 \\ \mu_Z &= (g^2 + g'^2)^{1/2} F/2\end{aligned}\quad (12)$$

with $F = 246 \text{ GeV}$

where g, g' are the coupling constants of the groups $SU(2)_L$ and $U(1)$ respectively, and F is the analogue of the pion decay constant. The interesting feature of our model is that echo-quarks, being massive after the chiral symmetry breaking, couple to the massless ordinary quarks according to Eq. (11') via the heavy vector bosons V_μ^- . Thus the light quarks obtain a current quark mass according to the mechanism of the figure, where in our picture the techniquarks become echo-quarks and the extended technicolour vector bosons become V_μ^- . The value of the current quark mass is still given by Eq. (1) *). On the other hand, the great problem which ruled out extended technicolour does not appear at all here.

Until now, we have limited ourselves to haplon-quarks, quarks and echo-quarks of the first generation. One can repeat the game for the rest of the families. However, in that case one has to abandon asymptotic freedom, since the QCD β function given by

$$\beta = -\frac{1}{2\pi} \left(11 - \frac{2}{3} n_3 - \frac{10}{3} n_6 - 4 n_8 \dots \right) \alpha_s^2 + O(\alpha_s^3) \quad (13)$$

suggests that for six ordinary quarks ($n_3 = 6$) and for more than two echo-quarks ($n_6 > 2$) asymptotic freedom is lost. Of course, echo-quarks decouple at lower energies and perturbative QCD is safe.

Next, one could extend the mechanism to leptons. In that case, one could put the lepton-haplons in the $(3, \bar{3})$ representation of $SU(3) \otimes SU(3)$ which,

*) The proposal that light fermion masses are produced radiatively is quite old. The authors in Ref. 12) were trying to explain in this way the electron muon mass ratio. Similar attempts have been made for quarks¹³⁾ and the first realistic example was given in Ref. 14). In horizontal models, mass matrix elements have been introduced radiatively¹⁵⁾. The same proposal has been discussed in the framework of GUTs¹⁶⁾ as well as in the framework of subquark models¹⁷⁾.

in turn, is decomposed in $1+8$ under $SU(3)_C$. Thus, one can identify the singlets with the ordinary leptons and the octets with echo-leptons. The cancellation of the anomalies now requires quarks (ordinary and echo) to be three times more than leptons, since the usual factor 3 in the contribution of the quark anomalies which in the standard model balance the contribution of lepton anomalies, does not enter here.

In conclusion, we have presented a model where the strong sector is enlarged to $SU(3) \times SU(3)$, which is subsequently broken to $SU(3)_L$. The elementary fermionic constituents, the haplons, are decomposed under $SU(3)_L$ to ordinary quarks and to echo-quarks. Echo-quarks, since they belong to higher representations than ordinary quarks, are assumed to break chirality in the scale $O(1 \text{ TeV})$ and in this way they are able to give masses to $SU(2)_L \otimes U(1)$ vector bosons and current quark and lepton masses.

ACKNOWLEDGEMENTS

It is a pleasure to thank J. Iliopoulos and D. Sutherland for discussions, and J. Bell, J.-M. Gérard, D. Nanopoulos and J. Prentki for reading the manuscript and for giving their comments.

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FIGURE CAPTION

Fig. 1 : The mechanism for light fermion mass generation through the exchange of a heavy ETC gauge boson E_μ (wavy line), and the technifermion-antitechnifermion condensation $\bullet$. In the present scheme, ETC gauge bosons are replaced by V_μ^- and technifermions by echo-fermions.