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INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

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INTERNATIONAL ATOMIC ENERGY AGENCY

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ABSTRACT

We discuss the effect of scalar partons arising in QCD if the colour symmetry is spontaneously broken. We make use of a previous result, which states that such scalars can be incorporated into the theory without disturbing asymptotic freedom.

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In unified gauge theories such as the $SU(4)_{flavour} \times SU(4)_{colour}$ model of Pati and Salam [1)] scalar fields play an important role in spontaneously breaking the colour gauge symmetry and providing masses to the gluons. One remarkable property of this spontaneous breaking mechanism is that for integer-charge quark (and gluon) models, the colour part of the photon in lepton induced deep inelastic processes essentially decouples from the flavour sector in the deep inelastic limit, so that quarks are effectively fractionally charged at short distances. This decoupling increases as $m^2(q^2)/q^2 \to 0$, where $m^2(q^2)$ is the effective mass of the gluons. It also takes place in the scalar sector, so that the integer-charge scalar particles will effectively also behave as fractionally charged at short distances, and gluons as if they were electrically neutral.*) In this sense, then, the theory behaves like conventional QCD, with fractionally charged quarks and neutral transverse gluons plus additional (fractionally charged) scalars. In the exact symmetry limit when $m \to 0$, the photon is pure flavour rather than flavour-colour mixture, and the theory is exact conventional QCD. [2)]

In this note we wish to estimate the departures from the standard (fractionally charged quark) QCD, particularly in σ_L/σ_T due to the presence of scalars for large q^2 . While the formalism presented applies quite generally, we are specifically interested in the "minimal" scalar Higgs multiplet made up of nine complex scalar fields, consisting of three colour triplets, with the same charge matrix as u, d and s quarks.**) Such a multiplet is necessary if, after spontaneous breaking, the colour symmetry is to remain a "good" global classification symmetry. In the exact theory, eight

*) The precise statement is that the electron-hadron interaction is of the long-range form $\left(q^{-2} J^{lep} J^{flavour}\right)$ plus terms of the type $\left(q^{-2} - [q^2 - m^2(q^2)]^{-1}\right) J^{lep} J^{colour}$. For low frequencies $\left(q^2 \to 0 \text{ but } m^2(q^2) \neq 0\right)$ hadronic electrical charge is given by $(J^{flavour} + J^{colour})$ and the quarks, scalars and gluons are integrally charged. For large q^2 , when $m^2(q^2)/q^2 \to 0$ as $q^2 \to \infty$, however, the second term and with it the colour part of charge becomes decoupled.

**) In Ref.1, four triplets corresponding to u, d, s and c quarks are introduced. However, one may expect the fourth triplet to be very heavy; in any case three triplets constitute the "minimal" set for the emergence of a good colour classification symmetry.

of the 18 real scalar fields will join with the transverse gluons to produce the longitudinal components of the now massive gluonic fields, while the remaining scalars remain as live Higgs fields. This is assuming that the mass parameter for these scalar particles has the requisite sign for the Higgs-Kibble phenomenon to occur. In order to preserve asymptotic freedom of the theory including scalars [3], we shall also assume that the bare coupling parameters λ_0 in the $\lambda_0\phi^4$ terms of the Higgs potential are zero. This means there is only an effective λ, induced by the interaction of the scalars and the gluons. In this case the recent discussion of Salam and Strathdee applies and the theory retains asymptotic freedom.[4]

In accordance with the colour decoupling theorem mentioned above [1], we shall ignore the colour contribution to the photon, so that the starting point of our discussion will be the following effective extended QCD Lagrangian, containing scalars without the $\lambda_0\phi^4$ coupling:

$$\mathcal{L} = -\bar{q}\, \not{D}\, q - \frac{1}{4} F^{a}_{\mu\nu} F^{\mu\nu}_{a} + \frac{1}{2} (\nabla_{\mu}\phi)^2 - m_0^2\phi^2 \quad ,$$

where

$$(D_\mu)_{ij} = (\not{\partial} + m)\, \delta_{ij} + g\, A^a_\mu (t_a)_{ij} \quad ,$$

$$(\nabla_\mu)_{ij} = \partial_\mu\, \delta_{ij} - g\, A^a_\mu (\theta_a)_{ij} \quad . \tag{1}$$

In (1) the matrices t_0 and θ_a depend on the particular colour representation of the quarks and scalars.(The flavour indices, not indicated here, are treated in the usual way.) We shall specialize to the particular case of the "minimal" Higgs representation defined above later. For large q^2, when masses are irrelevant, the gluons A_μ 's are transverse and all scalars (including those which in the exact theory make up the longitudinal gluons in the conventional Higgs mechanism) contribute equally. The scalars have no Yukawa coupling to quarks.

Now, in the limit that m_0^2 can be neglected compared with other frequencies in the problem, the scalar part of the Lagrangian in (1) leads to much the same ultraviolet behaviour as the fermion sector, and by itself would behave like QCD with charged scalar partons. This means we can readily use the technique developed for vectors and fermions [5] to deduce the anomalous dimensions. Further, we can extend the recent results in QCD [6),7)], by which the evolution with Q^2 of the quark and transverse gluon distribution functions $q(x,Q^2)$ and $G(x,Q^2)$ can be calculated from coupled sets of integral equations, to estimate the scalar (Higgs) parton distribution $S(x,Q^2)$ generated at short distances.

The usual anomalous dimensions or critical exponents of QCD $({}_a\gamma^b_n)$ (where $a, b = F, V$) , involving only fermions (F) and vectors (V) are shown to order g^2 in Figs.1(a)-(d). If we include the scalar sector in (1), we have in addition the exponents shown in Figs. 1(e)-(g). The anomalous dimensions associated with the gluon propagator and the β function will also be modified by scalar loops (see Ref.3).

For fermions, the vertex corresponding to the dominant twist-two operator has the form [5]:

$$\langle O^{(F)}_{\mu_1\ldots\mu_n}\rangle = \frac{1}{2n}\sum_i k_{\mu_1}\cdots\gamma_{\mu_i}\cdots\gamma_{\mu_n} \quad . \tag{2}$$

This does not lead to gauge invariant diagrams, so to order g^2 one must add diagrams involving a vertex with two fermion lines and a vector line (see Fig.1(a)). This has the form

$$\frac{1}{24n(n-1)}\sum_{i\neq j} k_{\mu_1}\cdots g_{\mu_i\nu}\, p_{\mu_{i+1}}\cdots\gamma_{\mu_j}\cdots p_{\mu_n} \quad . \tag{3}$$

For scalars, the corresponding twist-two operator $\langle O^{(s)}_{\mu_1\ldots\mu_n}\rangle$ simply involves the replacement of γ_μ by k_μ in (2) and (3) and the appropriate modification of the statistical weight factors, so that for example

$$\langle O^{(S)}_{\mu_1\ldots\mu_n}\rangle = k_{\mu_1}\cdots k_{\mu_n} \quad . \tag{4}$$

In Ref.7 it was shown that in the appropriate non-covariant gauge, the gauge correcting additional diagrams disappear in order g^2 for fermions. It is simple to see that this also holds for scalars. Using essentially the same computation as in Refs. 5 and 7, we obtain the set of anomalous dimensions in Table I, where we also give the usual exponents in QCD.

In conventional QCD without the scalars the anomalous dimension matrix

$$\underset{\sim}{\gamma}_n = \begin{bmatrix} {}_F\gamma^F_n & {}_V\gamma^F_n \\ {}_F\gamma^V_n & {}_V\gamma^V_n \end{bmatrix} \tag{5}$$

becomes rapidly diagonal as n becomes large [see Ref.8 for a review of this and related questions] because the off-diagonal terms drop off rapidly like $({}_F\gamma^V_n)({}_V\gamma^F_n) \sim 1/n^2$, so that

$$(\gamma_n^{\text{singlet}}) = (\gamma_n^{\text{non-singlet}})\left[1 - O\left(\frac{1}{n^2 \log n}\right)\right] \tag{6}$$

since ${}_F\gamma_n^F = \gamma_n^{\text{non-singlet}} \sim \log n$ as $n \to \infty$. This simply says that we expect the quark distributions to dominate as $x \to 1$, remembering that only the fermions couple to $F_2(x,Q^2)$ in deep inelastic scattering.

When we add the scalars, instead of the 2×2 matrix (5) we now have a 3×3 matrix of anomalous dimensions $\underset{\sim}{\gamma}_n$, which also has rapidly vanishing off-diagonal elements as n becomes large, while ${}_S\gamma_n^S$ also grows like $\log n$ as $n \to \infty$. This simply means in the Bethe-Salpeter language (Refs.6 and 7) that the fermion and scalar iterations exponentiate independently. The structure functions can be written in terms of the eigenvalues $\{\gamma_n^{(i)}\}$ of $\underset{\sim}{\gamma}_n$, through the decomposition

$$F_2(x,Q^2) = \sum_a e_a^2 \int \frac{dn}{2\pi i} \frac{1}{x^n} c_a P^{(i)}_{n,ab} \left[\frac{\alpha(Q^2)}{\alpha(Q_0^2)}\right]^{\gamma_n^{(i)}} A_n^b \; . \tag{7}$$

with $F_1 = -\frac{1}{2x} F_2$ for fermions, $F_1 = 0$ for scalars and where $P^{(i)}_{n,ab} = U^{-1}_{ai} U_{ib}$ is the projection operator onto the i^{th} eigenvalue $[U^{-1} \underset{\sim}{\gamma}_n U = \text{diag}(\underset{\sim}{\gamma}_n)]$ and $A_n^b\, p_{\mu_1} \ldots p_{\mu_n} = \langle P | O^{(b)}_{\mu_1 \ldots \mu_n} | P \rangle$; (a, b run over F, V and S).

The above statement about the large n dependence of γ_n , means that as $x \to 1$ only the fermions and scalars contribute to the sum over a and b and $P^{[F,S]}_{n,ab} \to \delta_{ab}$. For low $Q^2 < Q_0^2$ we expect any primordial scalar distribution $S_a(x)$ to vanish faster than the quark distribution $q_a(x)$ as $x \to 1$. This simply corresponds to saying that only valence quarks dominate the nucleon wave function as $x \to 1$. In turn, this means $A_n^{(S)} << A_n^{(F)}$ for large n , so that we expect to see the effect of the scalars only in the lower moments and at smaller x .

To estimate the size of the A_n^S 's in Eq.(7), one may observe that the factorization relation $M_{\text{non-singlet}}(n,Q^2) = K\left[M_{\text{non-singlet}}(n',Q^2)\right]^{\gamma_{n'}/\gamma_n}$ for the moments $M(n,Q^2)$ of the non-singlet structure functions does not hold if A_n^S is large for any n . *This factorization relation is currently being used* (Ref.9) to study QCD through the ratio of moments of F_2 and F_3 . However, the

relevant n's are rather large, typically $n \geq 3$, for which one may expect A_n^S 's to be small in any case.

To study the effects of scalars on the smaller moments, consider the $n = 2$ moment, for which the appropriate operators in the short distance expansion are related to the energy and momentum tensor $\theta_{\mu\nu}$. This moment corresponds to the momentum sum rule and we discuss the effect of the scalars on the QCD result concerning the fraction of the total momentum in the nucleon carried by the quarks. This fraction tends, at high Q^2, to the value $r_q = c_1(t)\,[2c_2(t) + c_1(t)] \approx 0.36$ for $SU(3)_{\text{colour}}$ and three quark flavours u, d and s . Here r_q is given by

$$r_q = \sum_i \int_0^1 dx\, x\, q_i(x,Q^2) = \int_0^1 dx \left\{\frac{9}{2}\,[F_2^{ep} + F_2^{en}] - \frac{3}{4}\,[F_2^{\nu p} + F_2^{\nu n}]\right\} . \tag{8}$$

The above result corresponds to a zero eigenvalue of the 2×2 anomalous dimension matrix γ_2 . When we add the scalars, γ_2 is replaced by the 3×3 matrix:

$$\underset{\sim}{\gamma}_2 = \frac{2}{3\pi}\begin{vmatrix} 2c_2(t) & 0 & -c_1(t) \\ 0 & \frac{3}{4}c_2(\theta) & -\frac{1}{4}c_1(\theta) \\ -2c_2(t) & -\frac{3}{4}c_2(\theta) & c_1(t) + \frac{1}{4}c_1(\theta) \end{vmatrix} \tag{9}$$

Specialize to the case of the three colour triplets of scalar fields, mentioned at the beginning of this note, which have the same representations as u, d and s quarks, so that $t^a_{ij} = \theta^a_{ij} = \lambda^a_{ij}$ (the Gell-Mann matrices for the 3* representation) and $c_1(t) = c_1(\theta) = 3/2$, $c_2(t) = c_2(\theta) = 4/3$. In this case γ_2 has the eigenvalues $\gamma^{(i)} = 0,\ 0.25,\ 0.92$ (compared with $\gamma^{(i)} = 0,\ 0.88$ for QCD without these scalars). Further, it is a simple matter to compute numerically the appropriate projection operators for these eigenvalues, from which one finds $r_q(\infty) = 0.29$, $r_s(\infty) = 0.19$ and $r_v(\infty) = 0.52$ (compared with $r_q(\infty) = 0.36$ and $r_v(\infty) = 0.64$ for conventional QCD). However, because of the small second eigenvalue $\gamma^{(2)} = 0.25$, which leads to corrections of the order $[\alpha(Q^2)/\alpha(Q_0^2)]^A$, with $A = 0.2$, these limits will only be reached at high Q^2 . Further, if the gluons and scalars are only generated at short distances, so that for $Q^2 < Q_0^2$, $r_q = 1$, $r_s = r_v = 0$, then it will be very difficult to see any difference from conventional QCD. On the other hand, if scalars are primordially present at lower Q^2, there should be measurable differences with QCD, depending on how large the primordial component is. For example, the sum rule (8) with the minimal set of scalars (with the average $\langle e_i^2\rangle_{scalars} = \langle e_i^2\rangle_{quarks} = 2/9$, will yield the value $r_q(\infty) + r_s(\infty) = 0.48$, which is appreciably different from the QCD value of $r_q(\infty) = 0.36$. However, this estimate pertains only to very high Q^2.

In any case, since $r_s(\infty)$ tends to a finite constant, the sum rule (8) with the replacement of F_2 by $F_2^{scalar} = F_2 - \frac{1}{2x}F_1$ can be used to determine its value in any particular model. In fact a useful sum rule to study will be:

$$\int_0^1 dx \left(F_2^{eN} - \frac{1}{2x}F_1^{eN}\right) = \sum_{i = \text{scalar flavour}} e'^2 \int_0^1 dx\, x\, S_i(x,Q^2) = \langle e_i^2\rangle_{scalars}\, r_s \quad . \tag{10}$$

This will then give an estimate of the normalization of S_i.

An even more useful conclusion is reached if one looks at the effect of the scalar partons on σ_L/σ_T, which will tend to a constant [10], as opposed to $1/\log Q^2$ for QCD without scalars. This constant can be estimated in the following way. Assuming that $G(x)$, $S(x)$ are, respectively, the primordial gluon and scalar distributions for $Q^2 < Q_0^2$, we can, to a reasonable approximation, estimate (Refs.6 and 7) the scalar distribution at large Q^2 from the formula:

$$S(x,Q^2) = S(x) + \frac{\alpha(Q^2)}{2\pi}\log Q^2/\Lambda^2 \int_x^1 \frac{dx'}{x'}\left[p_{SG}\left(\frac{x}{x'}\right) G(x') + p_{SS}\left(\frac{x}{x'}\right) S(x')\right] \quad , \tag{11}$$

where p_{SG} and p_{SS} are the Altarelli-Parisi kernels [6] for the scalar-gluon transitions. They satisfy

$$\int_0^1 dx\, x^{n-1}\, p_{SG}(x) = 2\pi\, {}_V\gamma_n^S \quad ,$$
$$\int_0^1 dx\, x^{n-1}\, p_{SS}(x) = 2\pi\, {}_S\gamma_n^S \quad . \tag{12}$$

In particular, $p_{SG}(x) = 2c_1(\theta)x(1-x)$, so that if we assume for simplicity $S(x) = 0$ and $G(x) = c(1-x)^d/x$, it is a simple matter to show that

$$S(x,Q^2) \xrightarrow[x \to 0]{} \frac{\alpha(Q^2)\log Q^2/\Lambda^2}{2\pi} 2c_1(\theta)\frac{1}{6}\frac{c}{x} \quad , \tag{13}$$

while for $x \to 1$, $S(x,Q^2)/G(x,Q^2) \sim (1-x)$. The longitudinal structure function is given by $F_L(x,Q^2) = \sum_i e_i^2 S_i(x,Q^2)$, while in QCD without scalars, $F_L = 0$, except for higher order corrections. These can be written in the form (see Ref.8):

$$F_L(x,Q^2) = \alpha(Q^2)c_2(t)\, x^2 \int_x^1 \frac{dx'}{x'^3} F_2(x',Q^2) \quad . \tag{14}$$

From (14) we learn (assuming $F_2(x,Q^2) \to$ constant as $x \to 0$) that

$$\frac{F_L(x,Q^2)}{F_2(x,Q^2)} \xrightarrow[x \to 0]{} \frac{1}{2}\alpha(Q^2)c_2(t) \quad . \tag{15}$$

With the above simplifications, the presence of charged scalars leads to the result

$$\frac{F_L(x,Q^2)}{G(x,Q^2)} \xrightarrow[x \to 0]{} \sum_{i = \text{scalars}} e_i^2 \frac{c_1(\theta)/3}{[\frac{11}{3}c_2(G) - \frac{4}{3}c_1(t) - \frac{1}{6}c_1(\theta)]} \quad . \tag{16}$$

In the specific model of Ref.1, the ratio (16) is 0.08, where we have used the fact that in the minimal model $\sum e_i^2 = 2 \times \frac{2}{3}$ and at current values of Q^2, $G(x,Q^2)/F_2(x,Q^2) = 3.5$ (see Ref.11 for fits to the data, where, however, only conventional QCD is taken into account). Thus $\sigma_L/\sigma_T = F_L/F_2 \simeq 0.28$ as $x \to 0$. This should be compared with the QCD prediction using (15), with $\alpha(Q^2) = 0.1$, which gives a value $\sigma_L/\sigma_T \sim 0.1$ <u>at small x</u>. Since we have set the primordial scalar distribution $S(x) = 0$, the above estimate is minimal in the sense that it generated at short distances (i.e. large Q^2). It is clear from our observations that a very detailed measurement of σ_L/σ_T as a function of Q^2 and x, will be valuable in deciding whether the scalar sector is present or not, if only as a dynamical effect.

TABLE I

Anomalous dimensions ${}_a\gamma_n^b$ with a, b = F, S, V.

	FERMIONS		SCALARS
${}_F\gamma_n^F$	$\frac{c_2(t)}{\pi}\left[1-\frac{2}{n(n+1)}+4\sum_{j=2}^{n}\frac{1}{j}\right]$	${}_S\gamma_n^S$	$\frac{c_2(\theta)}{2\pi}\left[-1+4\sum_{j=2}^{n}\frac{1}{j}\right]$
${}_F\gamma_n^V$	$-\frac{c_2(t)}{\pi}\left[\frac{n^2+n+2}{(n-1)n(n+1)}\right]$	${}_S\gamma_n^V$	$-\frac{c_2(\theta)}{\pi}\frac{1}{(n-1)n}$
${}_V\gamma_n^F$	$-\frac{2c_1(t)}{\pi}\left[\frac{n^2+n+2}{n(n+1)(n+2)}\right]$	${}_V\gamma_n^S$	$-\frac{c_1(\theta)}{\pi}\frac{1}{(n+1)(n+2)}$
${}_V\gamma_n^V$	$\frac{c_2(G)}{\pi}\left[\frac{1}{3}-\frac{4}{n(n-1)}-\frac{4}{(n+1)(n+2)}+4\sum_{2}^{n}\frac{1}{j}\right]+\frac{2}{3\pi}\left[c_1(t)+\frac{1}{4}c_1(\theta)\right]$		

$c_2(G)\delta_{ab}=\sum_{cd} f_{acd}\, f_{bcd}$, $c_1(M)\delta_{ab}=\mathrm{Tr}(M^aM^b)$

$c_2(M)\delta_{ij}=\sum_a (M^aM^a)_{ij}$, f_{abc} refers to the structure constants of the gauge group and M^a_{ij} is a particular representation matrix.

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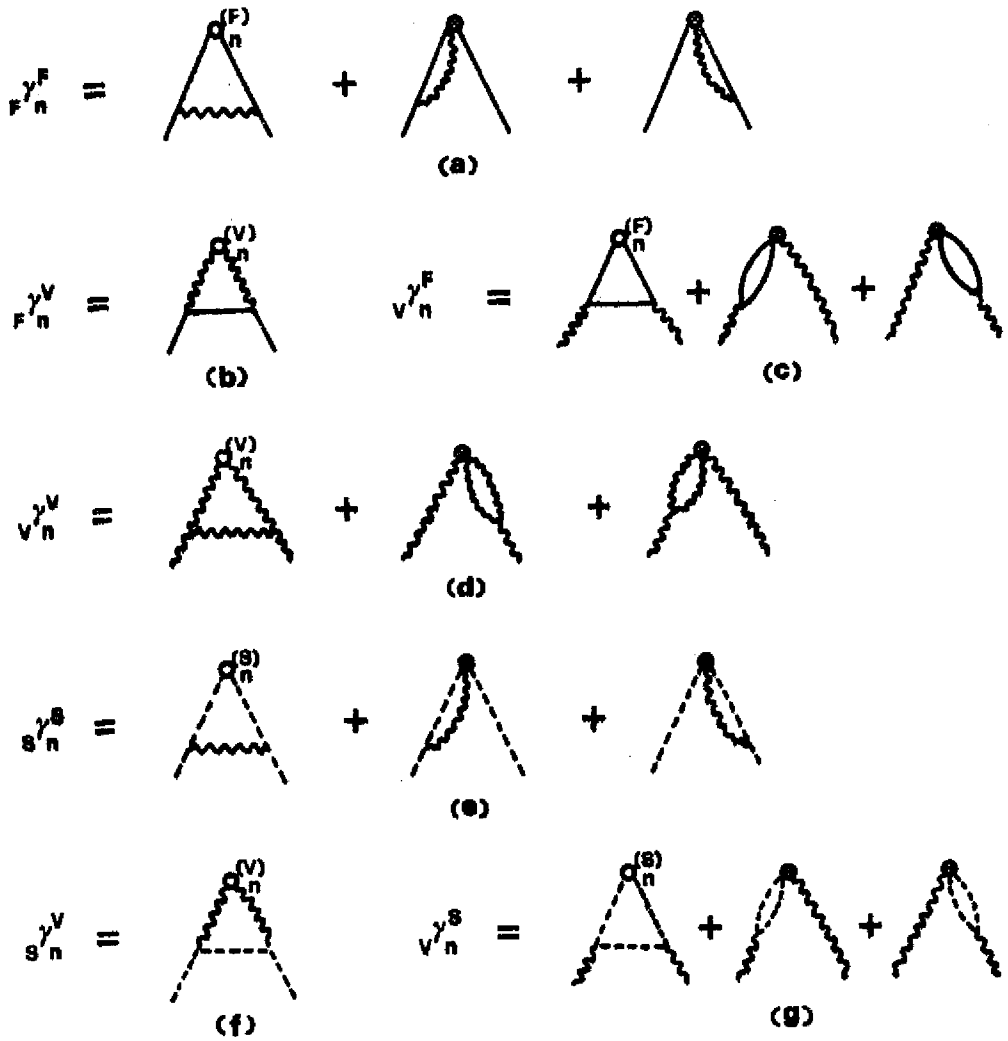

Fig. 1

Solid lines are quarks, wavy lines are gluons and dotted lines are scalars.

(a)-(d) graphs for anomalous dimensions of QCD;

(e)-(g) show the additional contribution for scalars.

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