СООБЩЕНИЯ
ОБЪЕДИНЕННОГО
ИНСТИТУТА
ЯДЕРНЫХ
ИССЛЕДОВАНИЙ

ДУБНА

E2 - 9329

S.V.Goloskokov, S.G.Mikhov, P.T.Morozov,
D.B.Stamenov

DEEP INELASTIC SCATTERING IN SPONTANEOUSLY BROKEN GAUGE MODELS

1975

E2 - 9329

S.V.Goloskokov, S.G.Mikhov, P.T.Morozov,
D.B.Stamenov

DEEP INELASTIC SCATTERING IN SPONTANEOUSLY BROKEN GAUGE MODELS

INTRODUCTION

The discovery of asymptotic freedom in non-abelian gauge theories[1] revived the hopes to describe the physical phenomena at high energies and high momentum transfers by means of quantum field theory. If the strong interactions are described by a field theory of this kind, it is possible to calculate some dynamical quantities using perturbation theory methods suitably improved with the help of renormalization group. These ideas have been used to analyze e^+e^- annihilation into hadrons[2] and deep inelastic lepton hadron scattering[3,4]. The asymptotic q^2 behaviour of the moments of the structure functions has been calculated and it has been shown that Bjorken scaling in deep inelastic lepton hadron processes is violated by powers of $\ln(-q^2/\mu^2)$. To obtain all these results a gauge invariant interaction of massless Yang-Mills field with coloured quarks as a model of strong interactions has been used. In order to construct realistic non-abelian gauge models of strong interactions one has to ascribe masses to Yang-Mills fields in a gauge invariant way (only in this case the model will be renormalizable). The usual mechanism to do this is the Higgs mechanism of spontaneous symmetry breaking[5].This mechanism is based on the introduction of new scalar fields and their quartic self-interaction. It is well known that a scalar field by itself, with $h\phi^4$ coupling is not asymptotically free in the main logarithmic approximation. Therefore it is the vector mesons only, that should provide a chance to preserve asymptotic freedom in spontaneously broken gauge theories. In paper[3] the wide class of gauge models including scalar fields was considered. Unfortunately the authors have not succeeded in finding any physically acceptable model which is asymptotically free with respect to both (Yang-Mills and quartic) coupling constants in main logarithmic approximation. Basically the problem is that asymptotic freedom requires large gauge groups and scalar representations of small dimensions. But if such representations are used, then only a part of the vector mesons acquires masses, but not all of them. Having in mind all this the authors of this work were forced to believe in some hypothetical "dynamical" mechanism of symmetry breaking.

But the account of the lower logarithmic terms is capable of changing the situation in a non-trivial way. In papers[6,7] the contribution of the lower logarithmic terms to the asymptotic behaviour of the invariant coupling constants (ICC) has been investigated in massive Yang-Mills theories based on the Higgs-Kibble mechanism. It was shown that in the second logarithmic approximation the invariant Yang-Mills coupling constant tends to zero in the region of large space-like momenta, while the quartic coupling constant of the scalar field tends to a finite value. The analysis of the higher order approximations has shown the property of having a finite asymptotic value of the quartic scalar selfinteraction to be sensitive upon the range of approximation. The next approximation not only can shift the asymptotic value of this coupling constant but even give rise to the so-called "zero-charge" situation /8/.The renormalization group analysis has pointed out that within the framework of pertirbation theory it is impossible to make any conclusions about asymptotic behaviour of the ICC in models that are not asymptotically free.

In the present paper the simplest model of the Yang-Mills theory with spontaneous symmetry breaking is considered. That is the SU(2) Kibble model. Within this model all the vector mesons acquire masses. It has two coupling constants: g the Yang-Mills coupling constant and h that of the quartic selfinteraction of the scalar fields. The aim of the paper is to analyze the deep inelastic lepton-hadron scattering under the assumption that the ICC $\bar{h}$ tends to a finite asymptotic value without spoiling the asymptotic freedom with respect to the ICC $\bar{g}$. Just the asymptotic q^2 behaviour of the moments of the structure functions is calculated. It is shown that Bjorken scaling is violated by powers of $\ln(-q^2/\mu^2)$.To obtain this result the renormalization group equations of Callan-Symanzik for the coefficient functions /10/ in the Wilson expansion of the product of two currents near the light-cone are used.

RELATION BETWEEN THE STRUCTURE FUNCTIONS AND THE WILSON COEFFICIENT FUNCTIONS

The behaviour of the cross sections of the deep inelastic lepton hadron processes is governed by the behaviour of the structure functions $F_k^a(q^2,\omega)$ in the region: $-q^2 \to \infty$, $\nu \to \infty$ and $\omega = -\frac{2M\nu}{q^2}$ fixed ($k=1,2$ for electron hadron and $k=1,2,3$ for lepton hadron processes, the index a is related to the internal symmetry). These functions are defined by the matrix element of the commutator of two electromagnetic or weak hadron currents. For the electron hadron processes we have for instance:

$$\frac{1}{2\pi M}\int d^4x\, e^{iqx} \langle p|[J_\mu(x)\, J_\nu(0)]|p\rangle =$$

$$= (-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2})\, W_1(q^2,\nu) + \frac{1}{M^2}(p_\mu - \frac{p\cdot q}{q^2}q_\mu)(p_\nu - \frac{p\cdot q}{q^2}q_\nu)\, W_2(q^2,\nu), \quad (1)$$

where M is the hadron mass, p is the momentum of the hadron, q the momentum transfer to the hadrons, $\nu = \frac{p\cdot q}{M}$, J_μ is the electromagnetic hadron current, spin averaging is implicit, $MW_1 = F_1$, $\nu W_2 = F_2$.

As is well known/12/, the large q^2 and ν behaviour of the matrix element (1) is determined by the singularity of the product $J_\mu(x)\, J_\nu(0)$ on the light cone, $x^2=0$. The structure of this product is given by the Wilson operator product expansion at light-like distances which in this case takes the form:

$$J_\mu(x)J_\nu(0) =$$

$$=(g_{\mu\nu}\partial_\alpha\partial^\alpha - \partial_\mu\partial_\nu)\frac{1}{x^2 - i\epsilon x_0}\{\sum_{n=0}^{\infty}\sum_i G_i^{(n)}(x^2 - i\epsilon x_0; g_k) O^i_{\mu_1 \dots \mu_n}(0) x^{\mu_1}\dots x^{\mu_n} + \dots\}. \tag{2}$$

$$+(g_{\mu\nu}\partial^\alpha\partial^\beta + g_\mu^\alpha g_\nu^\beta \partial_\rho\partial^\rho - g_\nu^\beta\partial_\mu^\alpha - g_\mu^\beta\partial_\nu^\alpha) \times$$

$$\times\{\sum_{n=0}^{\infty}\sum_i E_i^{(n)}(x^2 - i\epsilon x_0; g_k) O^i_{\alpha\beta\mu_1\dots\mu_n}(0) x^{\mu_1}\dots x^{\mu_n} + \dots\},$$

where $O^i_{\mu_1 \dots \mu_n}$ are spin n finite local operators (traceless and symmetric) with smallest canonical twist (canonical twist = canonical dimension - spin $= n+2-n=2$). n is even because of charge-conjugation invariance. The label i denotes the various operators of equal twist which might occur in the expansion. $G_i^{(n)}(x\ )$ and $E_i^{(n)}(x\ )$ are c-number functions of x^2 and the coupling constants g_k of the theory. The three dots denote other operators of higher twist, which contribution to the structure functions is suppressed, to any finite order of perturbation theory by powers of q^2.

The large q^2 behaviour of the moments of the structure functions is determined by the behaviour of the Fourier transforms of the functions $G_i^{(n)}$, $E_i^{(n)}$ in this region. In fact

$$\int_1^\infty \frac{d\omega}{\omega^{n+2}} F_2(\omega, q^2) \underset{-q^2\to\infty}{=} \sum_i C^i_{n+2}\tilde{E}_i^{(n)}(q^2; q_k), \tag{3}$$

$$\int_1^\infty \frac{d\omega}{\omega^{n+1}}\left[\frac{1}{2}\omega F_2(\omega,q^2) - F_1(\omega, q^2)\right] \underset{-q^2\to\infty}{=} \sum_i C^i_n \tilde{G}_i^{(n)}(q^2; g_k),$$

where

$$\tilde{E}_i^{(n)}(q^2;g_k) = \frac{i}{2}(q^2)^{n+2}\left(-\frac{\partial}{\partial q^2}\right)^n \int d^4x e^{iqx} E_i^{(n)}(x^2 - i\epsilon; g_k),$$

$$\tilde{G}_i^{(n)}(q^2;g_k) = i(q^2)^{n+1}\left(-\frac{\partial}{\partial q^2}\right)^n \int d^4x e^{iqx} \frac{G_i^{(n)}(x^2 - i\epsilon; g_k)}{x^2 - i\epsilon}. \qquad (4)$$

In Eq. (3) C_n^i are unknown constants determined by the hadronic matrix elements of the operators $O^i_{\mu_1 \dots \mu_n}$

$$\langle p|O^i_{\mu_1 \dots \mu_n}(0)|p\rangle_{\text{spin average}} = (i)^n C^i_n p_{\mu_1} \dots p_{\mu_n} + \qquad (5)$$

$+$(term containing $g_{\mu_i \mu_j}$).

Note that only the term proportional to $p_{\mu_1} \dots p_{\mu_n}$ in eq. (3) yields leading terms in the Wilson expansion on the light-cone.

Bjorken scaling implies that all functions $\tilde{E}_i^{(n)}$ are constants independent of q^2. It is the case of free quark models only, where the latter can hold. In interaction theory the functions $\tilde{E}_i^{(n)}$, $\tilde{G}_i^{(n)}$ are non-trivial functions of q^2 and of the coupling constants of the theory (g_k). To determine the q^2 dependence of $\tilde{E}_i^{(n)}$, $\tilde{G}_i^{(n)}$ we shall use the renormalization group equations for the Wilson coefficients.

DEEP INELASTIC SCATTERINC IN THE KIBBLE MODEL

Let us now consider the Kibble model. Its original Lagrangian looks as follows:

$$\mathcal{L} = -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu}_a + |\partial_\mu \chi - ig\frac{\vec{\tau}}{2}\vec{B}_\mu \chi|^2 + \frac{m^2}{2}\bar{\chi}\chi - \lambda(\bar{\chi}\chi)^2 + \mathcal{L}_g, \quad (6)$$

where

$$F^a_{\mu\nu} = \partial_\mu B^a_\nu - \partial_\nu B^a_\mu + g\epsilon^{abc} B^b_\mu B^c_\nu$$

and χ is nonhermitian spinor, $\chi = \begin{pmatrix} \chi^o \\ \chi^+ \end{pmatrix}$. . $\mathcal{L}_g$ in Eq. (6) describes the interaction with the so-called Feynman-Faddeev-Popov ghosts/13/.

Due to spontaneous symmetry breaking the vacuum expectation value of the field χ is different from zero. The latter provides masses for the vector mesons. By means of a canonical transformation one can go over from the scalar fields χ to the scalar fields σ, $\vec{\phi}$ having zero vacuum expactation value

$$\chi^o = \xi + \frac{1}{\sqrt{2}}(\sigma + i\phi_3), \quad \chi^+ = \frac{1}{\sqrt{2}}(i\phi_1 - \phi_2), \qquad (7)$$

where ξ is a constant iso-spinor equal to the vacuum expectation value of χ.

Within this model a triplet of Coldstone bosons $\vec{\phi}$ arises, but it can be eliminated from the spectrum of physical states by a suitable choice of the gauge. For this particular model the dominant gauge invariant operators of twist two appearing in the operator product expansion (2) of two electromagnetic currents near the light cone will be:

$$O^B_{\mu_1 \dots \mu_n} = \frac{i^{n-2}}{2n!} \mathrm{Tr}\{ F_{\mu_1 \alpha} \nabla_{\mu_2} \dots \nabla_{\mu_{n-1}} F^{\alpha}_{\mu_n} + \qquad (8a)$$

+ permutations of vector indices }

- trace terms

$$O^{\chi}_{\mu_1 \dots \mu_n} = \frac{i^n}{n!} \{ \bar{\chi} \nabla_{\mu_1} \dots \nabla_{\mu_n} \chi + \text{permutations} \} \qquad (8b)$$

- trace terms,

where

$$\nabla_\mu F^a_{\nu\lambda} = (\delta^{ac}\partial_\mu + g\epsilon^{abc} B^b_\mu) F^c_{\nu\lambda}$$

and

$$\nabla_\mu \chi = (\partial_\mu - ig\frac{\tau^a}{2} B_\mu^a)\chi , \qquad n\text{-even.}$$

Note that there is a gauge in which $\vec{\phi} = 0$ and the operators $O^{\chi}_{\mu_1 \dots \mu_n}$ go over to

$$O^{\sigma}_{\mu_1 \dots \mu_n} = \frac{i^n}{2}\{\sigma\, \partial_{\mu_1} \dots \partial_{\mu_n} \sigma - \frac{ig}{2n!} S \sum_{i=1}^{n} \sigma\, \partial_{\mu_1} \dots B_{\mu_i} \dots \partial_{\mu_n} \sigma -$$

$$- \frac{g^2}{8n!} S \sum_{i=1}^{n} \sum_{\substack{j=1 \\ i \neq j}}^{n} \sigma\, \partial_{\mu_1} \dots B_{\mu_i} \dots B_{\mu_j} \dots \partial_{\mu_n} \sigma + O(g^3)\} .$$

In addition to the operators O^{B}, O^{χ} there are composite operators formed of Faddeev-Popov ghosts which may have the same twist and therefore are expected to mix with these operators. But in reference /3/ it is shown that there is a gauge in which no ghosts are present so that the mixing does not appear.

The Fourier transforms of the Wilson coefficient functions $E_i^{(n)}(-q^2/\mu^2, g, h)$ ($i = 1, 2$ for our model, g is the Yang-Mills coupling constant, $h = (\frac{3}{4\pi^2})\frac{\lambda}{4}$ the coupling constant of quartic self-interaction $-\frac{4\pi^2}{3} h(\sigma^2 + \vec{\phi}^2)^2$, μ^2 is the point of subtraction) satisfy the following renormalization group equations:

$$[(\frac{\partial}{\partial t} - \beta_g(g,h)\frac{\partial}{\partial g} - \beta_h(g,h)\frac{\partial}{\partial h})\delta_{ij} + (\gamma^{T})_{ij}^{(n+2)}]\bar{E}_j^{(n)}(-q^2/\mu^2, g, h) = 0$$

where (9)

$$\beta_g(g,h) = \frac{d\bar{g}(t,g,h)}{dt}\Big|_{t=0}, \quad \beta_h(g,h) = \frac{d\bar{h}(t,g,h)}{dt}\Big|_{t=0}, \qquad (10)$$

$$t = \frac{1}{2}\ln(-q^2/\mu^2) .$$

In Eq. (10) $\bar{g}$ and $\bar{h}$ are the asymptotical invariant coupling constants which we choose in the form:

$$\bar{g} = g\Gamma_{3B} d_B^{3/2}; \qquad \bar{g}(0) = g,$$
$$\bar{h} = h\Gamma_{4\sigma} d_\sigma^2 ; \qquad \bar{h}(0) = h. \tag{11}$$

Here Γ_{3B} and $\Gamma_{4\sigma}$ are the normalized symmetric vertex functions, d_B and d_σ are dimensionless Creen functions. In Eq. (9) γ is the matrix of anomalous dimensions of the operators O^B and O^σ (γ^T is the transposed γ)

$$\gamma^{(n)}_{ij} = 2\delta_{ij}\gamma_j + 2b^{(n)}_{ij}, \tag{12}$$

where

$$\gamma_i(g,h) = \frac{\partial \ell n\, d_i(L,g,h)}{\partial L}\Big|_{L=0} (L = 2t), \tag{13}$$

is the anomalous dimensions of fields B, σ ($i = 1$ stands for field B and $i = 2$ for field σ) and b_{ij} are related to the relevant matrix elements of $O^i_{\mu_1\cdots\mu_n}$ by

$$b^{(n)}_{ij} p_{\mu_1}\cdots p_{\mu_n} + (g_{\mu_i\mu_j}) = \frac{\partial \Gamma_{\Phi_i\Phi_i O^j_{\mu_1\cdots\mu_n}}(L,g,h)}{\partial L}\Big|_{L=0}. \tag{14}$$

In Eq. (14) $\Gamma_{\Phi_i\Phi_i O_j}$ are

$$\Gamma_{\Phi_i\Phi_i O^j_{\mu_1\cdots\mu_n}} = \langle 0|T(\Phi_i(-p) O^j_{\mu_1\cdots\mu_n}(0)\Phi_i(p))|0\rangle_A. \tag{15}$$

Here we use notation $\Phi_1 = B$, $\Phi_2 = \sigma$, $O^1 = O^B$, $O^2 = O^\sigma$.

Note that Eqs. (9) are also satisfied by the functions $\tilde{G}^{(n)}_i$. That is why in the following we consider the functions $\tilde{E}^{(n)}_i$ only, having in mind that the same results hold for $\tilde{G}^{(n)}_i$.

The solution of Eq. (9) is

$$\tilde{E}_i^{(n)}(-q^2/\mu^2, g, h) = \{T\exp[-\int_0^t \gamma^{T(n+2)}(\bar{g},\bar{h})\,dx]\}_{ij}\,\tilde{E}_j(1,\bar{g},\bar{h}) , \quad (16)$$

where T implies that the exponential is to be t-ordered. The large q^2 behaviour of $\tilde{E}_i^{(n)}$ is governed by the large t behaviour of ICC $\bar{g}$ and $\bar{h}$, which in turn is determined by the ultraviolet stable fixed points of the following system of renormalization group equations

$$\frac{d\bar{g}^2(t,g,h)}{dt} = 2\bar{g}\beta_g(\bar{g},\bar{h}) ,$$
$$\frac{d\bar{h}(t,g,h)}{dt} = \beta_h(\bar{g},\bar{h}) . \quad (17)$$

The functions β_g and β_h are defined by the conditions (10) and can be obtained from the perturbation calculations as the series in $\bar{g}$ and $\bar{h}$:

$$2\bar{g}\beta_g(\bar{g},\bar{h}) = -a(\bar{g}^2)^2 + b(\bar{g}^2)^2\bar{h}^2 + c(\bar{g}^2)^3 + \ldots ,$$
$$\beta_h(\bar{g},\bar{h}) = \alpha\bar{h}^2 - \beta\bar{h}^3 - \gamma\bar{g}^2\bar{h} + \delta(\bar{g}^2)^2 + \ldots . \quad (18)$$

It is well known from the proofs of the renormalization of spontaneously broken Higgs theories /14/ that their ultraviolet behaviour is that of the underlying symmetry theory. In other words the symmetry breaking is a "generalized mass terms" and does not affect the asymptotic (Euclidean) behaviour of the theory. That is why all perturbation calculations can be performed as if all masses and dimensional coupling constants are zero.

Like in reference [7] we assume that ICC $\bar{h}(t)$ tends to a finite value H for large space-like momenta

$$\bar{h}(t) \underset{t \to \infty}{\to} H. \tag{19}$$

Then it follows from Eqs. (17,18) that the point $\bar{g}^2 = 0, \bar{h} = H$ will be an unltaviolet stable fixed point for sufficiently small H. Equations (17,18) imply that large t -behaviour of $\bar{g}^2$ is:

$$\bar{g}^2(t) \underset{t \to \infty}{\to} \frac{1}{\lambda(H^2)\, t}. \tag{20}$$

where

$$\lambda(H^2) = a - bH^2 \tag{21}$$

with account of the lowest in H^2 perturbation theory order. Here a and b are the corresponding coefficients in Eq. (18). The coefficient a has been calculated in papers [3,7] and for the present model is $a = \frac{43}{6 \cdot 8\pi^2}$. The anomalous dimension matrix $\gamma^{(n)}(\bar{g},\bar{h})$ for large t takes the form

$$\gamma^{(n)}(\bar{g},\bar{h}) = \begin{pmatrix} \gamma^{(n)}_{BB}(H^2)\,\bar{g}^2 + O(\bar{g}^4) & \gamma^{(n)}_{B\sigma}(H^2)\,\bar{g}^2 + O(\bar{g}^4) \\ \gamma^{(n)}_{\sigma B}(H^2)\,\bar{g}^2 + O(\bar{g}^4) & \gamma^{(n)}_{\sigma\sigma}(H^2) + O(\bar{g}^2) \end{pmatrix} \tag{22}$$

Then the solution of the system of equations (16) for the functions $\tilde{E}_i^{(n)}$ in the limit of large t, taking into account the explicit form (22) of $\gamma^{(n)}$, is

$$\tilde{E}_B^{(n)}(-q^2/\mu^2, g, h) \underset{-q^2\to\infty}{\to} [A_1 \tilde{E}_B^{(n)}(1;0,H) + A_2 \tilde{E}_\sigma^{(n)}(1;0,H)] \left(\frac{1}{\ln(-q^2/\mu^2)}\right)^{\frac{\gamma_{BB}^{(n+2)}(H^2)}{\lambda(H^2)}} \quad (23)$$

$$\tilde{E}_\sigma^{(n)}(-q^2/\mu^2, g, h) \underset{-q^2\to\infty}{\to} [A_1 \tilde{E}_B^{(n)}(1;0,H) + A_2 \tilde{E}_\sigma^{(n)}(1;0,H)] \left(\frac{1}{\ln(-q^2/\mu^2)}\right)^{1+\frac{\gamma_{BB}^{(n+2)}(H^2)}{\lambda(H^2)}},$$

where A_1, A_2 are unknown constants, $n = 2, 4 \ldots$. Note that in the case $n = 0$ (which account the contribution of the energy-momentum operator in the Wilson expansion) the solution of the Eq. (16) is

$$\tilde{E}_B^{(0)} \underset{-q^2\to\infty}{\to} \text{const}, \qquad \tilde{E}_\sigma^{(0)} \underset{-q^2\to\infty}{\to} \text{const}.$$

In Eq. (23) $\gamma_{BB}^{(n)}(H^2)$ with account of the lowest in H^2 perturbation theory order is

$$\gamma_{BB}^{(n)}(H^2) = \hat{\gamma}_{BB}^{(n)} - dH^2. \quad (24)$$

$\hat{\gamma}_{BB}^{(n)}$ have been calculated in papers [3,4] For the present case the result is as follows

$$\hat{\gamma}_{BB}^{(n)} = \frac{1}{4\pi^2}\left[1 - \frac{4}{n(n+1)} - \frac{4}{(n+1)(n+2)} + 4\sum_{\ell=2}^{n} \frac{1}{\ell}\right]. \quad (25)$$

The constant d in Eq. (24) is the coefficient for the term $g^2 h^2 \ln(-p^2/\mu^2)$ in the perturbation expansion of the dimensionless Green function d_B.

So, the large q^2 behaviour of the moments of the structure function $F_2(\omega, q^2)$ observes a logarithmic deviation from Bjorken scaling

$$\int_1^\infty \frac{d\omega}{\omega^n} F_2(\omega, q^2) \underset{-q^2 \to \infty}{\to} \mathrm{const}\left(\frac{1}{\ell n(-q^2/\mu^2)}\right)^{\frac{\gamma^{(n)}_{BB}(H^2)}{\lambda(H^2)}} + O\left(\frac{1}{\ell n(-q^2/\mu^2)}\right). \qquad (26)$$

The result for the structure function $F_L(\omega, q^2)$ is just the same. The same results hold for neutrino hadron processes.

SUMMARY

Deep inelastic scattering in the Kibble model (the simplest spontaneously gauge model in which all the vector mesons acquire masses) is analysed. This model is not asymptotically free. Using the hypothesis that $\bar{h}(t) \underset{t \to \infty}{\to} H$ (Eq. (19)) the asymptotic q^2 behaviour of the moments of the structure functions is calculated. It is shown that Bjorken scaling is violated by powers of $\ell n(-q^2/\mu^2)$ (as in the case of asymptotically free models). The powers of the logarithms are functions of the unknown parameter H, which by assumption is small. So they can be calculated by perturbation theory methods. The present experimental situation shows that the deviation of Bjorken scaling is small so the hypothesis of small H seems to be plausible. We note that the model under consideration is not realistic, since the deep inelastic processes here are $\ell + \sigma \to \ell + X$, where σ is a hypothetic scalar particle. This is the reason why we did not

deal with the coefficients b and d in formulae (21) and (24) respectively. It should be interesting to consider models involving usual hadrons. We intend to perform these calculations which can give us the possibility to determine the constant μ from the present experimental data. We hope that it will be one and the same for all experiments of the kind. In that case such models can pretend to describe strong interaction.

We want to express our gratitude to D.V.Shirkov and A.V.Efremov for the useful discussions and remarks. We thank also D.V.Shirkov for attracting our attenuation to this problem. Thanks are also due to A.A.Vladimirov and D.I.Kazakov for some discussions.

REFERENCES

1. D.Gross, F.Wilczek. Phys.Rev.Lett., 30, 1343 (1973); H.D.Politzer. Phys.Rev.Lett., 30, 1346 (1973).
2. T.Appelquist, H.Ceorgi. Phys.Rev., D8, 4000 (1973); A.Zee. Phys.Rev., D8, 4038 (1973).
3. D.Cross, F.Wilczek. Phys.Rev., D8, 3633 (1973); Phys.Rev., D9, 980 (1974).
4. H.Ceorgi, H.D.Politzer. Phys.Rev., D9, 416 (1974).
5. P.W.Higgs. Phys.Rev.Lett., 13, 508 (1964); Phys.Rev., 145, 1156 (1966).
6. V.V.Belokurov et al. Phys.Lett., 47B, 359 (1973).
7. V.V.Belokurov et al. JINR preprint, E2-7562 and TMF, 19, 149 (1974).

8. L.D.Landau, I.Ya.Pomeranchuk. Doklady AN SSSR, 102, 489 (1955); E.S.Fradkin. JETP 28, 750 (1955).
9. T.W.B.Kibble. Phys.Rev., 155, 1554 (1967).
10. K.Symanzik. Comm.Math.Phys., 23, 49 (1971); C.C.Callan. Phys.Rev., D5, 3202 (1972); N.Christ, B.Hasslacher, A.Mueller. Phys.Rev., D6, 3543 (1972).
11. K.Wilson. Phys.Rev., 179, 1499 (1969).
12. H.Leutwyler, J.Stern. Nucl.Phys., B20, 77 (1970); R.Brandt. Phys.Rev., D1, 2808 (1970).
13. R.P.Feynman. Acta Phys. Polon., 26, 697 (1963); L.D.Faddeev, V.N.Popov. Phys.Lett., 25B, 30 (1967); B. de Witt. Phys.Rev., 162, 1195, 1293 (1967).
14. B.W.Lee, J.Zinn-Justin. Phys.Rev., D5, 3121, 3137, 3155 (1972); A.A.Slavnov. Particles and Nuclei, vol.5, part 3 (1974).

Received by Publishing Department
on November 21, 1975.