Analyticity and Minimality of Nonperturbative Contributions in Perturbative Region for $\bar{\alpha}_s$

Aleksey I. Alekseev and Boris A. Arbuzov

Institute for High Energy Physics, Protvino, Moscow Region, 142284 Russia

Abstract

It is shown, that the possibility of a freezing of QCD running coupling constant at zero in the approach with "forced analyticity" can not be in accord with Schwinger-Dyson equation for gluon propagator. We propose to add to the analytic expression the well-known infrared singular term $1/q^2$ as well as pole term corresponding to "excited gluon". With this example we formulate the principle of minimality of nonperturbative contributions in perturbative (ultraviolet) region, which allows us to fix ambiguities in introduction of nonperturbative terms and maintain the finiteness of the gluon condensate. As a result we obtain estimates of the gluon condensate, which quite agree with existing data. The nonzero effective mass of the "excited gluon" leads also to some interesting qualitative consequences.

PACS number(s): 12.38.Aw, 12.38.Lg

---

1Electronic address: alekseev@mx.ihep.su
2Electronic address: arbuzov@mx.ihep.su
The discovery of the asymptotic freedom property \[1\] in non-Abelian gauge theories turned to be a decisive factor in the formation QCD as the strong interaction theory. The negative sign of QCD $\beta$-function $\beta(g^2) = \beta_0 g^4 + \ldots$, $\beta_0 = -b_0/(16\pi^2), b_0 = 11C_2/3 - 2N_f/3$ in the vicinity of zero provided the number of active quarks being not too large (for $SU_c(3)$ $N_f \leq 16$) gives coupling constant, which describes quarks and gluons interaction at large Euclidean $q^2$, i.e. at small distances,

$$\bar{g}^2(q^2/\mu^2, g) = \frac{g^2}{1 - \beta_0 g^2 \ln(q^2/\mu^2)},$$ (1)
tending towards zero. Therefore in the deep Euclidean region we are allowed to use perturbation theory. In expression (1), which takes into account the main logarithms, $\mu$ is a normalization point. An account of the next $g^2$ corrections does not change asymptotic behaviour (1) for $q^2 \to \infty$. By introducing dimensional constant $\Lambda^2 = \mu^2 \exp(-4\pi/(b_0 \alpha_s)), \alpha_s = g^2/4\pi$, we turn from explicitly renormalization invariant expression (1) to the following formula

$$\bar{\alpha}_s(q^2) = \frac{4\pi}{b_0 \ln(q^2/\Lambda^2)},$$ (2)

It is reasonable to estimate parameter $\Lambda$ in approximate expression (2) to be around of few hundreds MeV. With decreasing $q^2$ effective constant (4) increases, that may indicate a tendency of unlimited growth of the interaction at large distances, leading to a confinement of coloured objects. However, at $q^2 = \Lambda^2$ in expression (3) the pole is present, which is nonphysical at least due to failing of the perturbation theory, starting from which formula (2) has been obtained.

In recent work [2] a solution of the problem of ghost pole was proposed with imposing of a condition of analyticity in $q^2$. The idea of "forced analyticity" goes back to works [3, 4] of the late fifties, which were dedicated to the problem of Landau-Pomeranchuk pole [5] in QED. Using for $\bar{\alpha}_s(q^2)$ a spectral representation without subtractions, the following expression for the running coupling constant was obtained in paper [2]

$$\bar{\alpha}_s^{(1)}(q^2) = \frac{4\pi}{b_0} \left[ \frac{1}{\ln(q^2/\Lambda^2)} + \frac{\Lambda^2}{\Lambda^2 - q^2} \right].$$ (3)

This expression has the asymptotic freedom property and its analyticity in the infrared region is due to nonperturbative contributions. It does not contain any additional parameter and has finite limit at zero, $\bar{\alpha}_s^{(1)}(0) = 4\pi/b_0 \simeq$
1.40 (freezing of the coupling constant), which depends only on symmetry factors. This limit turns out to be stable with respect to higher orders corrections.

As it is noted in work [4], a procedure of summation of leading logarithmic terms is not defined uniquely. A partial fixation of this ambiguity in QED is realized by using of a method of summation of the perturbation theory series under the sign of the spectral integral of the Källen-Lehmann representation. Nevertheless after such summation a functional ambiguity remains, which from one side does not violate correct analytic properties of Green functions in a complex plane of a corresponding invariant variable and from the other side contains nonanalytic dependence on constant $g^2$. In work [6] while investigating the photon propagator in QED it was shown, that ambiguities in summation procedure of the diagram series can be removed provided one demands not only the validity of spectral representation, but also the fulfillment of equations of motion.

In the present note we consider a problem of consistency of the constant behaviour of the effective charge in the infrared region with Schwinger-Dyson (SD) equation for a gluon propagator. Further we include into consideration nonperturbative terms, the singular in the infrared region term $\sim 1/q^2$ in particular, the necessity of the renormalization invariance being taken into account. Then we discuss possibilities of an adjustment of demands of confinement, asymptotic freedom, analyticity, an accordance with the perturbation theory and correspondence with estimates of the gluon condensate value.

To study the problem of a possibility of a constant behaviour of the running constant in the infrared region let us consider the integral SD equation for the gluon propagator in ghost-free axial gauge [7] $A_\mu^a \eta_\mu = 0$, $\eta_\mu$ — gauge vector, $\eta^2 \neq 0$. In this gauge the effective charge is directly connected with the gluon propagator and Slavnov-Taylor identities [8] have the most simple form. The important preference of the axial gauge consist in the possibility to exclude the term from SD equation, which contains the full four-gluon vertex by means of contraction of the equation with tensor $\eta_\mu \eta_\nu / \eta^2$.

In what follows we shall work in the Euclidean momentum space, where smallness of the momentum squared is immediately connected with smallness of its components. The equation to be considered has the form:

$$[D_{\mu\nu}^{-1}(p) - D^{(0)}_{\mu\nu}(p)] \frac{\eta_\mu \eta_\nu}{\eta^2} = \Pi_{\mu\nu}(p) \frac{\eta_\mu \eta_\nu}{\eta^2}.$$
\[\Pi_{\mu\nu}(p) = -C_2 g^2 \mu^{4-n} \int d^n k \Gamma^{(0)}_{3\mu\lambda\rho}(p, -k, k - p) D_{\lambda\sigma}(k) \times \Gamma_{3\sigma\delta\nu}(k, p - k, -p),\]  

where \(\Pi_{\mu\nu}(p)\) is the one-loop part of the polarization operator, \(D_{\mu\nu}(p)\) is the propagator, \(\Gamma^{(0)}_{\sigma\delta\nu}(k, p - k, -p)\) is the one-particle irreducible three-gluon vertex function, \(\Gamma^{(0)}_{\mu\lambda\rho}(p, -k, k - p)\) is the free three-gluon vertex.

We suppose the approximation \(D_{\mu\nu}(p) = Z(p^2) D^{(0)}_{\mu\nu}(p)\) to be appropriate to study the infrared region. Let us divide the momentum integration domain in expression (4) in two parts: \(k^2 < \lambda^2\) and \(k^2 > \lambda^2\), where \(\lambda\) is sufficiently small, but finite. Then domain \(k^2 > \lambda^2\) in the case of absence of kinematic singularities in three-gluon vertex gives a contribution, which is regular in \(p^2\) for \(p^2 \to 0\), and in domain \(k^2 < \lambda^2\) full Green functions can be approximated by free ones up to constant factors according to an assumption of running constant be frozen at zero. Then one can write

\[\Pi_{\mu\nu}(p) \frac{\eta_{\mu} \eta_{\nu}}{\eta^2} = -C_2 g^2 \mu^{4-n} Z(0) \int_0^\lambda d^n k \Gamma^{(0)}_{3\mu\lambda\rho}(p, -k, k - p) \times \Gamma_{3\sigma\delta\nu}(k, p - k, -p) \eta_{\mu} \eta_{\nu} / \eta^2 + Q(p^2; y, \lambda, n).\]  

Here \(y = (p\eta)^2 / p^2 \eta^2\) is the gauge parameter. The integration in formula (4) can be extended up to all the domain of momentum, that results in a change of the regular in \(p^2\) contribution \(Q\). Thus one has

\[\Pi_{\mu\nu}(p) \frac{\eta_{\mu} \eta_{\nu}}{\eta^2} = Z(0) \Pi^{(1)}_{\mu\nu}(p) \frac{\eta_{\mu} \eta_{\nu}}{\eta^2} + Q(p^2; y, n),\]  

where \(\Pi^{(1)}_{\mu\nu}(p)\) is the one-loop perturbation theory contribution to the polarization operator. This contribution is calculated in paper [4] and has rather complicated structure. Let us present the expression for the leading terms of convolution (6) at \(y \to 0\). We have

\[\Pi^{(1)}_{\mu\nu}(p) \frac{\eta_{\mu} \eta_{\nu}}{\eta^2} = C p^2 \left[ -\frac{22}{3\epsilon} - \frac{22}{3} \left( \gamma - 2 + \ln \frac{p^2}{4\pi \mu^2} \right) - \frac{70}{9} + \frac{40}{3} y \ln y + O(y, y^2 \ln y) \right].\]
Here $C = g^2 C_2 / 32 \pi^2$, $\gamma$ is the Euler constant. From expression (7) we see, that singularity at $y = 0$ is smooth and the limit at $y = 0$ does exist. Term $\sim 1/\epsilon$ ($n = 4 + 2\epsilon$) as well as constant ones could be absorbed into function $Q$ while the logarithm of the momentum squared necessarily persists. The equation for function $Z(p^2)$ takes the form

$$Z^{-1}(p^2) = 1 + Z(0) \frac{g^2 C_2}{16\pi^2} \frac{11}{3} \ln p^2 + Q(p^2; n).$$

(8)

We see, that behaviour $Z(p^2) \simeq Z(0) \neq 0$ for $p^2 \to 0$ does not agree with the SD equation.

This conclusion stimulate us to look for the possibilities different from the assumption on the finiteness of the coupling constant at zero. Recently the possibility of the soft singular power infrared behaviour of the gluon propagator was discussed [10], $D(q) \sim (q^2)^{-\beta}$, $q^2 \to 0$, where $\beta$ is a small positive non-integer number. In Ref. [11] the consistency of such behaviour with Eq. (4) was studied. A characteristic equation for the exponent $\beta$ was obtained and this equation was shown not to have solutions in the region $0 < \beta < 1$. The authors of Ref. [12] also come to the conclusion on the inconsistency of the soft singular infrared behaviour of the gluon propagator. The case of possible interference of power terms was studied in Ref. [13] and it was shown that in a rather wide interval $-1 < \beta < 3$ of the non-integer values of the exponent the characteristic equation has no solutions. At present the more singular, in comparison with free case, infrared behaviour of the form $D(q) \simeq M^2/(q^2)^2$, $q^2 \to 0$ seems to be the most justified [14, 15, 16]. The physical consequences of such enhancement of zero modes are discussed in the reviews [17, 18]. Bearing in mind the remarks stated above let us consider the following expression for the running coupling:

$$\tilde{\alpha}_s(q^2) = \frac{4\pi}{b_0} \left[ \frac{1}{\ln q^2/\Lambda^2} + \frac{\Lambda^2}{\Lambda^2 - q^2} + c \frac{\Lambda^2}{q^2} \right].$$

(9)

Let us represent this expression in explicitly renormalization invariant form. It can be done without solving the differential renormalization group equations. In this order we write $\tilde{\alpha}_s(q^2) = \tilde{g}^2(q^2/\mu^2, \ g^2)/4\pi$ and use the normalization condition $\tilde{g}^2(1, g^2) = g^2$. Then we obtain the equation for wanted dependence of the parameter $\Lambda^2$ on $g^2$ and $\mu^2$:

$$g^2/4\pi = \frac{4\pi}{b_0} \left[ \frac{1}{\ln \mu^2/\Lambda^2} + \frac{\Lambda^2}{\Lambda^2 - \mu^2} + c \frac{\Lambda^2}{\mu^2} \right].$$
From dimensional reasons $\Lambda^2 = \mu^2 \exp\{-\varphi(x)\}$, where $x = b_0 g^2 / 16\pi^2 = b_0 \alpha_s / 4\pi$, and for function $\varphi(x)$ we obtain the equation:

$$x = \frac{1}{\varphi(x)} + \frac{1}{1 - e^{\varphi(x)}} + ce^{-\varphi(x)}.$$ 

The solution of this equation at $c > 0$ is monotone decreasing function $\varphi(x)$, which has the behaviour $\varphi(x) \approx 1/x$ at $x \to 0$ and $\varphi(x) \approx -\ln(x/c)$ at $x \to +\infty$. The relation obtained ensures the renormalization invariance of $\bar{\alpha}_s(q^2)$. At low $g^2$ we obtain $\Lambda^2 = \mu^2 \exp\{-4\pi/(b_0 \alpha_s)\}$, which indicates the essentially nonperturbative character of both last terms of the Eq. (9) and this terms are absent in the perturbation theory. With given value of the QCD scale parameter $\Lambda$ the parameter $c$ can be fixed by the string tension $\kappa$ or the Regge slope $\alpha' = 1/(2\pi\kappa)$ assuming the linear confinement $V(r) \approx \kappa r = a^2 r$ at $r \to \infty$. We define the potential $V(r)$ of static $q\bar{q}$ interaction [19, 20] by means of three-dimensional Fourier transform of $\bar{\alpha}_s(q^2)/q^2$ with the contributions of only one dressed gluon exchange taken into account. This gives the following relation

$$c\Lambda^2 = \left(3b_0/8\pi\right)a^2 = \left(b_0/16\pi^2\right)g^2 M^2. \quad (10)$$

At large $q^2$ from Eq. (11) one obtains

$$\bar{\alpha}_s(q^2) = \frac{4\pi}{b_0} \left[ \frac{1}{\ln q^2/\Lambda^2} + \left( c - 1 \right) \frac{\Lambda^2}{q^2} - \frac{\Lambda^4}{(q^2)^2} + O((q^2)^{-3}) \right]. \quad (11)$$

From Eq. (11) it is seen that in the ultraviolet region the nonperturbative contributions decrease more rapid then all renormalization group improved perturbation theory corrections. The value $c = 1$ corresponds to maximal suppression of nonperturbative contributions in the ultraviolet region. Accepting this condition one obtains the connection of the QCD scale parameter $\Lambda$ and string tension $\kappa = a^2$ of the form $\Lambda^2 = 3b_0\kappa/8\pi$. Taking $a \simeq 0.42 \text{GeV}$ one obtains for $\Lambda$ reasonable estimation, $\Lambda \simeq 0.434 \text{GeV}$ ($b_0 = 9$ in the case of 3 light flavours).

Considering the nonperturbative contributions the following arguments can be expressed. One knows QCD to be renormalizable in the perturbation theory and, as usually, the renormalization procedure can be developed to remove the divergences in all orders. However, what about the nonperturbative contributions? If they bring in the additional divergences then the
problem of renormalization turns out to be unsolved. The situation when nonperturbative contributions do not violate the perturbative renormalization properties seems to be more attractive. It take place if the nonperturbative contributions decrease at momentum infinity sufficiently fast and do not introduce the divergences in observables. So, it is natural to demand their fastest of possible decrease at large momenta. An application of the principle of minimalty of nonperturbative contributions in the ultraviolet region will be shown further with taking as an example the important physical quantity, namely, the gluon condensate, \( K = \langle \alpha_s / \pi : G_{\mu \nu} G_{\mu \nu}^a \rangle \). According to the definition (see e.g., [17]) up to the quadratic approximation in the gluon fields one has after Wick rotation

\[
K = \frac{48}{\pi} \int \frac{d^4 k}{(2\pi)^4} \left( \bar{\alpha}_s(k^2) - \bar{\alpha}_s^{\text{pert}}(k^2) \right) = \frac{3}{\pi^3} \int_0^\infty \bar{\alpha}_s^{\text{np}}(y) y dy ,
\] (12)

where \( \bar{\alpha}_s^{\text{np}} \) is nonperturbative part of the running coupling constant. In our case the two last terms of Eq. (9) should be taken. By substituting this terms in Eq. (12) one can see the logarithmic divergences of the integral at infinity and at finite point \( k^2 = \Lambda^2 \).

The acceptance of the cancellation mechanism for the nonphysical perturbation theory singularities (2) by the nonperturbative contributions leads to the necessity of supplementary definition of the integral (12) near point \( k^2 = \Lambda^2 \). This problem can be reformulated as a problem of dividing of perturbative and nonperturbative contributions in \( \bar{\alpha}_s \) resulting in introduction of some parameter \( k_0 = 1 \div 2 \text{ GeV} \). This provides absence of the pole at \( k^2 = \Lambda^2 \) in both perturbative and nonperturbative parts. The divergence of the integral (12) at infinity stimulate the further modification of the running coupling constant. Going over from Eq. (3) to Eq. (9) the isolated singularity was introduced. In this case the singularity corresponding to the unitary cut was not changed and in accordance with the approach of Refs. [3, 4, 2] is determined by perturbation theory. Following to this logic let us consider the expression for \( \bar{\alpha}_s \) with one more isolated singularity in the time-like region. The tachion singularity in the space-like region, of cause, is prohibited.

The principle of minimality of nonperturbative contributions in ultraviolet region then leads to the following unique expression for the running coupling constant

\[
\bar{\alpha}_s(q^2) = \frac{4\pi}{b_0} \left( \frac{1}{\ln(q^2/\Lambda^2)} + \frac{\Lambda^2}{\Lambda^2 - q^2} + \frac{c\Lambda^2}{q^2} + \frac{(1-c)\Lambda^2}{q^2 + m^2_\varphi} \right),
\] (13)
with fixed residue and mass parameter \( m_g \),

\[
m_g^2 = \Lambda^2 / (c - 1),
\]

for the newly introduced term. The expression (13) can be represented in explicitly renormalization invariant form in a similar way to the expression (14). Nonperturbative contributions in Eq. (13) decrease at infinity as \( 1/q^6 \), the integral in Eq. (12) converges and we can obtain

\[
K = \frac{4}{3\pi^2} \Lambda^4 \{ \ln(c - 1) + k_0^2/\Lambda^2 + \ln(k_0^2/\Lambda^2 - 1) \}.
\]

Phenomenology gives the positive value of the gluon condensate \( K \) in the interval \((0.32 \, GeV)^4 - (0.38 \, GeV)^4\) [21, 22]. As an example we take values \( k_0 = 1.2 \div 1.3 \, GeV \). If one regards the string tension parameter to be given, then from Eqs. (14), (17) and (19) one has the dependencies of all the values under consideration on the parameter \( c \), which are presented in the Table I.

Note that values \( c = 1.063, \Lambda = 422 \, MeV, m_g = 1.682 \, GeV, \) \( k_0 = 1.265 \, GeV \) corresponds to the conventional value of the gluon condensate \( K = (0.33 \, GeV)^4 \). Certainly these results should be considered as tentative, but nevertheless they seems encouraging.

It is seen from Eq. (13) that the pole singularities are situated at two points \( q^2 = 0 \) and \( q^2 = -m_g^2 \). It corresponds to two effective gluon masses, 0 and \( m_g \). Therefore the physical meaning of the parameter \( m_g \) is not the constituent gluon mass but rather the mass of the exited state of the gluon. It is essential that the residue at \( m_g^2 \) is very small, so the states with the exited gluons should be quite narrow in contrast to the spectrum of the coupled massless gluons.

The qualitative picture of the glueball states corresponding to the running coupling constant (13) with \( m_g \approx 1.7 \, GeV \) could be the following:

1) The states \( gg \) — continuous spectrum and very wide resonances are probable;

2) The states \( gg' \) — resonances with probable mass interval 1500 – 1800 MeV and with width suppression factor \((1 - c)\);

3) The narrow states \( gg' \) — resonances with possible masses 3000 – 3600 MeV and with width suppression factor \((1 - c)^2\).

Note that in the region 2) there are the glueball candidates. The region 3) is insufficiently investigated, some indications in favour of the narrow states are showing up (see e.g., [23]).
Table 1: Parameters of the running coupling constant (13) and gluon condensate as functions of parameter $c$.

| $c$  | $\Lambda$, GeV | $m_g$, GeV | $K^{1/4}$, GeV $k_0 = 1.2$ GeV | $K^{1/4}$, GeV $k_0 = 1.25$ GeV | $K^{1/4}$, GeV $k_0 = 1.3$ GeV |
|------|----------------|------------|--------------------------------|--------------------------------|--------------------------------|
| 1.01 | 0.433          | 4.332      | 0.298                          | 0.309                          | 0.318                          |
| 1.02 | 0.431          | 3.048      | 0.307                          | 0.317                          | 0.326                          |
| 1.03 | 0.429          | 2.476      | 0.312                          | 0.321                          | 0.330                          |
| 1.04 | 0.427          | 2.134      | 0.315                          | 0.324                          | 0.332                          |
| 1.05 | 0.425          | 1.900      | 0.317                          | 0.326                          | 0.334                          |
| 1.06 | 0.423          | 1.726      | 0.319                          | 0.327                          | 0.335                          |
| 1.07 | 0.421          | 1.591      | 0.320                          | 0.328                          | 0.336                          |
| 1.08 | 0.419          | 1.481      | 0.321                          | 0.329                          | 0.337                          |
| 1.10 | 0.415          | 1.313      | 0.322                          | 0.330                          | 0.337                          |
| 1.12 | 0.411          | 1.187      | 0.323                          | 0.330                          | 0.337                          |
| 1.16 | 0.404          | 1.010      | 0.323                          | 0.330                          | 0.337                          |
| 1.20 | 0.397          | 0.889      | 0.322                          | 0.329                          | 0.336                          |
| 1.24 | 0.391          | 0.798      | 0.321                          | 0.328                          | 0.335                          |
| 1.30 | 0.382          | 0.697      | 0.319                          | 0.326                          | 0.332                          |
We would like to thank Yu.F. Pirogov and V.E. Rochev for interesting discussion. A.I.A. is grateful also to C.D. Roberts, J.M. Namyslovski, and J.P. Vary for stimulating discussions.

References

[1] D.J. Gross and F. Wilczek, Phys. Rev. Lett. 30, 1343 (1973); H.D. Politzer, Phys. Rev. Lett. 30, 1346 (1973).

[2] D.V. Shirkov, I.L. Solovtsov, JINR Rapid Comm. 2[76], 5 (1996).

[3] P.J. Redmond, Phys. Rev. 112, 1404 (1958).

[4] N.N. Bogolubov, A.A. Logunov, D.V. Shirkov, Sov. Phys. JETP 37, 805 (1959).

[5] L.D. Landau, I.Ya. Pomeranchuk, Doklady Acad. Nauk SSSR 102, 489 (1955).

[6] B.A. Arbuzov, Doklady Acad. Nauk SSSR 128, 1149 (1959).

[7] W. Kummer, Acta Phys. Austr. 41, 315 (1975).

[8] A.A. Slavnov, Teor. Mat. Fiz. 10, 153 (1972); J.C. Taylor, Nucl. Phys. B33, 436 (1971).

[9] A.I. Alekseev, ”Generalized prescription for unphysical axial gauge singularities”, ICTP Report No IC/91/359, 1991 (unpublished).

[10] J.R. Cudell and D.A. Ross, Nucl. Phys. B359, 247 (1991); G.R. Cudell, A.J. Gentles, D.A. Ross, Nucl. Phys. B440, 521 (1995).

[11] A.I. Alekseev, Phys. Lett. B334, 325 (1995).

[12] K. Büttner and M.R. Pennington, Phys. Rev. D52, 5220 (1995).

[13] A.I. Alekseev, Teor. Mat. Fiz. 106, 250 (1996).
[14] H. Pagels, Phys. Rev. D15, 2991 (1977); C. Nash and R.L. Stuller, Proc. Roy. Irish Acad. 78A, 217 (1978); S. Mandelstam, Phys. Rev. D20, 3223 (1979); N. Brown and M.R. Pennington, Phys. Rev. D38, 2266 (1988).

[15] M. Baker, J.S. Ball and F. Zachariasen, Nucl. Phys. B186, 531, 560 (1981).

[16] A.I. Alekseev, Yad. Fiz. 33, 516 (1981); A.I. Alekseev, V.F. Edneral, Yad. Fiz. 45, 1105 (1987).

[17] B.A. Arbuzov, Phys. Element. Part. Atom. Nucl. 19, 5 (1988).

[18] C.D. Roberts and A.G. Williams, Prog. Part. Nucl. Phys. 33, 477 (1994).

[19] N.N. Bogolubov and D.V. Shirkov, An Introduction to the Theory of Quantized Fields (Wiley-Interscience, New York, 1980).

[20] W. Buchmüller and S.-H.H. Tye, Phys. Rev. D24, 132 (1981).

[21] A.I. Vainshtein, V.I. Zakharov, V.A. Novikov, M.A. Shifman, Phys. Element. Part. Atom. Nucl. 13, 542 (1982).

[22] W. Greiner, A. Schäfer, Quantum Chromodynamics (Springer - Verlag, Berlin, 1994).

[23] A.N. Aliev et al. (EXCHARM collaboration), Yad. Fiz. 56-10, 100 (1993).