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Infrared asymptotics of gluon Green's functions in covariant gauge

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Abstract

It is considered whether the singular infrared behaviour $D(k) \sim M^2/k^4$ of the gluon propagator and the corresponding asymptotic behaviour of the gluon vertices are compatible with the Schwinger–Dyson equation in a covariant gauge. For the investigated asymptotic behaviours, the two-loop terms most singular in the infrared region are calculated. It is shown that there exists a distinguished covariant gauge in which one can find a self-consistent description of the lowest gluon and ghost Green's functions in the infrared region of QCD.

1. The infrared region has fundamental importance for the solution of basic problems in QCD. Weighty theoretical arguments suggest the singular infrared behaviour $D(k) \sim M^2/k^4$ of the gluon propagator, these arguments being based both on the traditional approach in quantum field theory [1, 2] as well as on Monte Carlo calculations in the lattice variant of QCD (see, for example, [3]). On the other hand, since this "stringlike" asymptotic behaviour corresponds to a linearly-rising potential [4, 5, 6], it finds experimental confirmation, particularly in the spectra of bound states of heavy quarks [7].

As regards the traditional approach [1, 2] based on asymptotic solution of the Schwinger–Dyson equations with allowance for the gauge identities, the conditions for realization of the asymptotic behaviour $1/k^4$ have in fact been investigated only in the ghostless axial gauge [1, 2, 8]. In this case, multiplication of the equation for the gluon propagator by the gauge vector η_μ makes it possible to separate the one-loop terms, thus greatly simplifying the analysis, which however is rendered incomplete, since the terms transverse with respect to η_μ are eliminated [9], and these, in general, may be more singular in the infrared region. In addition, in QCD with axial gauge there are certain difficulties in the quantization [10], which may lead to complications in the gluon propagator even in the perturbation theory.

It is therefore desirable to study the conditions for realization of the asymptotic behaviour $D(k) \sim M^2/k^4$ in a covariant gauge, in which, in particular, the problem of calculating the two-loop terms is much simpler than in the axial gauge. In this case, we must take into account the contribution of the Faddeev–Popov ghosts in both the equations and the gauge identities. However, there are grounds for believing that there exists a distinguished gauge in which the contribution of the ghosts is unimportant for determining the leading terms in the infrared asymptotics of the Green's functions. These arguments come from consideration of the Schwinger–Dyson equation for the ghost propagator $S(p)$:

$$S^{-1}(p) = {S^{(0)}}^{-1}(p) + \frac{g^2\, C_2}{(2\pi)^n \mathrm{i}} \int \mathrm{d}^n k\; D_{\mu\nu}(k)\; \Lambda^{(0)}_\mu(p-k, p; k)\; S(p-k)\; \Lambda_\nu(p, p-k; -k)\;, \quad (1)$$

where C_2 is a color factor (for the group $\mathrm{SU_c}(N)$ we have $C_2 = N$, $\Lambda^{(0)}_\mu$ and Λ_ν are the bare and the total ghost-ghost-gluon interaction vertices [11], and $D_{\mu\nu}(k)$ is the total gluon propagator. Here and below, we use dimensional regularization ($n = 4+2\varepsilon$, $\varepsilon \to 0$) when calculating the Feynman integrals. Following the basic assumptions, we consider the infrared behaviour of the gluon propagator

$$D_{\mu\nu}(k) = \frac{M^2}{(k^2)^2}\left(g_{\mu\nu} - d\,\frac{k_\mu k_\nu}{k^2}\right) \quad (2)$$

and show that when (2) is substituted Eq. (1) is satisfied by the free ghost propagator and free ghost vertex, $S = S^{(0)}$ and $\Lambda_\mu = \Lambda^{(0)}_\mu$, when the gauge parameter d takes the value (in n-dimensional space)

$$d = \frac{4}{5-n} = 4 + 8\varepsilon + 16\varepsilon^2 + \mathcal{O}(\varepsilon^3)\;. \quad (3)$$

The distinguished nature of this gauge has already been noted in the study of the behaviour of the quark propagator in QCD [12] and derives above all from the fact that the Fourier transform $\widetilde{D}_{\mu\nu}(x)$ of the propagator (2) under the condition (3) is transverse: $\widetilde{D}_{\mu\nu}(x)\, x_\nu = 0$. This property ensures cancelling of the principal infrared singularities [5] and, in particular, leads to vanishing of the loop integral in (1). In quantum electrodynamics the well-known Soloviev–Yennie gauge [13] has the same property of transversality in the x-space and, as a consequence of this, cancelling of the principal infrared singularities. The property of transversality in the x-space is also inherent in the propagator $< T(A_\mu(x) \times A_\nu(0)) >$ in the Fock–Schwinger gauge $x_\mu A_\mu(x) = 0$ [14], in which it is well known that there is no ghost contribution.

We are not yet able to assert definitely that the principal infrared singularities of the loop diagrams in the higher ghost Green's functions also vanish in the gauge (3) However, if we assume that the ghost contribution in the infrared region is unimportant when the distinguished gauge is used, then the asymptotic behaviour (2) of the gluon propagator must satisfy the condition of cancelling of the leading infrared singularities in the Schwinger–Dyson equation without allowance for the ghosts in either the equation itself or in the gauge identities. The main aim of this paper is to verify this necessary condition.

2. The Schwinger–Dyson equation for the gluon propagator is shown graphically in Fig. 1. The hatched circles are the total propagators and one-particle-irreducible vertices. The dashed line is the ghost propagator. In this equation, we need to know the infrared behaviours of the propagators and vertices. For the gluon propagator, we use (2), an the one-particle-irreducible vertices, which satisfy the necessary properties of symmetry and the requirement of absence of kinematic singularities [15], and also the assumption of absence of a contribution of ghosts to the Slavnov–Taylor gauge identities [16], can be conveniently represented by using in the infrared region [2] the effective Lagrangian

$$\mathcal{L}_{\text{eff}} = \frac{1}{4M^2}\, D^{ab}_\rho F^b_{\mu\nu}\, D^{ac}_\rho F^c_{\mu\nu} - \frac{\xi\, g}{6M^2}\, f^{abc}\, F^a_{\mu\nu}\, F^b_{\nu\rho}\, F^c_{\rho\mu}\, . \tag{4}$$

Here, D^{ab}_ρ and $F^a_{\mu\nu}$ are the covariant derivative and field tensor of the gluon field defined in the usual way [11], M occurs in the definition of the propagator (2), and ξ is a parameter that occurs in the three- and four-gluon vertices we employ and is not fixed by the gauge identities. It can be seen from (4) that the three- and four-gluon vertices are homogeneous functions of third and second degree, respectively, in the momenta they contain. Proceeding from this and bearing in mind (2), we show by a simple power counting that the two-loop terms (the first two diagrams on the right-hand side of the equation in Fig. 1) have in the general case infrared behaviour of order g^4M^4/q^2, where q is the external momentum, and the one-loop terms (the two following ones) are less singular — of order g^2M^2. The contribution of the ghost loops with free propagators and vertices (and it is this case that we investigate) and, a fortiori, the contribution of the quark loops are still less singular as $q^2 \to 0$. Thus, the first necessary condition for consistency of the asymptotics (2) and the vertices that follow from (4) with the Schwinger–Dyson equation is the vanishing of the coefficient of the maximal singularity g^4M^4/q^2. Note that when the problem was treated in the axial gauge [1, 2, 8] only the projection of the equation not containing the maximally singular two-loop terms was used.

Thus, the problem is to calculate the considered two-loop diagrams in which we have

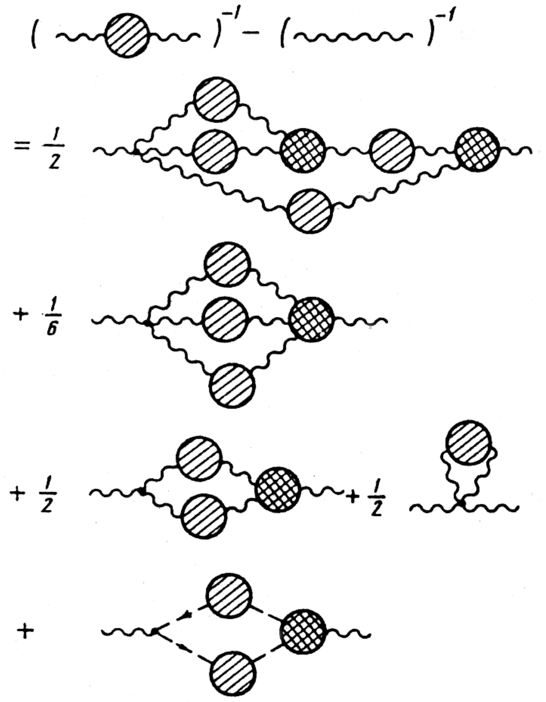

Figure 1: Schwinger–Dyson equation for gluon propagator

substituted the infrared behaviours of the propagator (2) and vertices for which the expressions follow from (4) [2]: the total three-gluon vertex

$$
\begin{aligned}
\Gamma^{abc}_{\lambda\mu\nu}(p,q,r) &= f^{abc}\,\Gamma_{\lambda\mu\nu}(p,q,r)\ , \\
\Gamma_{\lambda\mu\nu}(p,q,r) &= \frac{\mathrm{i}g}{M^2}\Big\{ g_{\lambda\mu}[p_\nu(p^2-(pq)) - q_\nu(q^2-(pq))] \\
&+ g_{\mu\nu}[q_\lambda(q^2-(qr)) - r_\lambda(r^2-(qr))] + g_{\nu\lambda}[r_\mu(r^2-(rp)) - p_\mu(p^2-(rp))] \\
&+ q_\lambda p_\mu p_\nu - q_\lambda p_\mu q_\nu + q_\lambda r_\mu q_\nu - r_\lambda r_\mu q_\nu + r_\lambda r_\mu p_\nu - r_\lambda p_\mu p_\nu \\
&+ \xi[g_{\lambda\mu}(p_\nu(qr) - q_\nu(pr)) + g_{\mu\nu}(q_\lambda(pr) - r_\lambda(pq)) + g_{\nu\lambda}(r_\mu(pq) - p_\nu(rq)) \\
&+ r_\lambda p_\mu q_\nu - q_\lambda r_\mu p_\nu]\Big\}\ , \qquad (5)
\end{aligned}
$$

and the total four-gluon vertex

$$
\begin{aligned}
\Gamma^{abcd}_{\mu\nu\rho\lambda}(p,q,r,k) &= f^{abe}f^{cde}\,\Pi_{\mu\nu\rho\lambda}(p,q,r,k) + f^{ace}f^{bde}\,\Pi_{\mu\rho\nu\lambda}(p,r,q,k) \\
&+ f^{ade}f^{cbe}\,\Pi_{\mu\lambda\rho\nu}(p,k,r,q)\ , \\
\Pi_{\mu\nu\rho\lambda}(p,q,r,k) &= \frac{g^2}{M^2}\Big\{ g_{\mu\rho}g_{\nu\lambda}[2(pr) + 2(qk) + (pk) + (qr)] \\
&- g_{\mu\lambda}g_{\nu\rho}[2(qr) + 2(pk) + (pr) + (qk)] \\
&- g_{\mu\rho}[p_\nu p_\lambda + r_\nu r_\lambda - r_\nu(p_\lambda + q_\lambda) - (r_\nu + k_\nu)p_\lambda + k_\nu q_\lambda] \\
&+ g_{\mu\lambda}[p_\nu p_\rho + k_\nu k_\rho - k_\nu(p_\rho + q_\rho) - (r_\nu + k_\nu)p_\rho + r_\nu q_\rho]
\end{aligned}
$$

$$
\begin{aligned}
&-g_{\nu\lambda}[q_\mu q_\rho + k_\mu k_\rho - k_\mu(p_\rho + q_\rho) - (r_\mu + k_\mu)q_\rho + r_\mu p_\rho] \\
&+g_{\nu\rho}[q_\mu q_\lambda + r_\mu r_\lambda - r_\mu(p_\lambda + q_\lambda) - (r_\mu + k_\mu)q_\lambda + k_\mu p_\lambda] \\
&-\xi\Big[(g_{\mu\sigma}g_{\nu\lambda} - g_{\mu\lambda}g_{\nu\rho})\,((pq) + (kr)) \\
&-g_{\mu\rho}\,(k_\nu r_\lambda + p_\nu q_\lambda) + g_{\mu\lambda}\,(r_\nu k_\rho + p_\nu q_\rho) \\
&-g_{\nu\lambda}\,(r_\mu k_\rho + q_\mu p_\rho) + g_{\nu\rho}\,(k_\mu r_\lambda + q_\mu p_\lambda) \\
&+g_{\mu\nu}\,(p_\rho q_\lambda - q_\rho p_\lambda) + g_{\rho\lambda}\,(r_\mu k_\nu - k_\mu r_\nu)\Big]\Big\}
\end{aligned}
\tag{6}
$$

Using the symmetry properties, we can significantly simplify the expressions for the two-loop terms. Then after summation over the colour indices for the group $\mathrm{SU_c}(N)$, we obtain the following expression for the sum of the two-loop terms of the right-hand side of the equation in Fig. 1:

$$
\begin{aligned}
&-\delta^{ab}\,\frac{3N^2}{4}\,\frac{1}{(2\pi)^{2n}}\,\left(U_{\mu\nu}(q) + W_{\mu\nu}(q)\right)\,, \\
U_{\mu\nu}(q) &= \int\!\!\int \mathrm{d}^n p\,\mathrm{d}^n t\;\Pi^{(0)}_{\mu\sigma\lambda\rho}\;\Gamma_{\lambda'\rho'\tau'}(p, t-p, -t)\;\Gamma_{\tau\sigma'\nu}(t, q-t, -q) \\
&\qquad \times D_{\lambda\lambda'}(p)\;D_{\rho\rho'}(t-p)\;D_{\sigma\sigma'}(q-t)\;D_{\tau\tau'}(t)\,, \\
W_{\mu\nu}(q) &= -\int\!\!\int \mathrm{d}^n p\,\mathrm{d}^n t\;\Pi^{(0)}_{\mu\sigma\lambda\rho}\;\Pi_{\lambda'\rho'\sigma'\nu}(p, t-p, q-t, -q) \\
&\qquad \times D_{\lambda\lambda'}(p)\;D_{\rho\rho'}(t-p)\;D_{\sigma\sigma'}(q-t)\,,
\end{aligned}
\tag{7}
$$

where

$$
\Pi^{(0)}_{\mu\sigma\lambda\rho} = g^2(g_{\mu\lambda}g_{\sigma\rho} - g_{\mu\rho}g_{\sigma\lambda})\,.
$$

Here, $U_{\mu\nu}$ and $W_{\mu\nu}$ have two tensor structures, $g_{\mu\nu}$ and $q_\mu q_\nu$, and we calculate the independent contractions of these tensors:

$$
\begin{aligned}
A\cdot U^{(1)} &= g_{\mu\nu}U_{\mu\nu}(q), \qquad A\cdot U^{(2)} = \frac{q_\mu q_\nu}{q^2}U_{\mu\nu}(q), \\
A\cdot W^{(1)} &= g_{\mu\nu}W_{\mu\nu}(q), \qquad A\cdot W^{(2)} = \frac{q_\mu q_\nu}{q^2}W_{\mu\nu}(q),
\end{aligned}
\tag{8}
$$

where for convenience we have introduced the factor

$$
A = 2\mathrm{i}^{2-4\varepsilon}\,\pi^{4+2\varepsilon}\,\frac{\Gamma^3(1+\varepsilon)\,\Gamma(1-2\varepsilon)}{\Gamma(1+3\varepsilon)}\,g^4 M^4 (q^2)^{-1+2\varepsilon}\,.
$$

Substitution of (2). (5) and (6) in (7) leads to a need to make a large number of calculations, which were done using the REDUCE system for analytic calculations [17] at the Institute for Nuclear Physics of Moscow State University. At the same time, the two-loop integrals encountered in the present work can be calculated by the formulae given in [18]. To calculate the specific integrals in these expressions, it is necessary to make a Laurent expansion with respect to ε.

We calculated the contractions (8) for arbitrary values of the gauge parameter d and the parameter ξ, which characterizes the contribution of the transverse part of the three-gluon vertex. Taking into account the singular and finite terms of the expansion in

$\varepsilon = (n-4)/2$, we obtain

$$
\begin{aligned}
U^{(1)} &= \left[-\frac{195-15\xi}{2\varepsilon^2}+\frac{271+86\xi-6\xi^2}{4\varepsilon}+\frac{81-124\xi-\xi^2}{4}\right] \\
&+d\left[\frac{546-27\xi}{8\varepsilon^2}-\frac{1666+63\xi}{16\varepsilon}+\frac{1314+597\xi}{32}\right] \\
&+d^2\left[-\frac{507-12\xi}{32\varepsilon^2}+\frac{2219-8\xi}{64\varepsilon}-\frac{3601+128\xi}{128}\right]+d^3\left[-\frac{39}{32\varepsilon^2}-\frac{103}{32\varepsilon}+\frac{25}{8}\right] , \\
U^{(2)} &= \left[-\frac{135}{8\varepsilon^2}+\frac{159+27\xi}{8\varepsilon}-\frac{868+69\xi}{16}\right]+d\left[\frac{363}{32\varepsilon^2}-\frac{1133+48\xi}{64\varepsilon}+\frac{4987+128\xi}{128}\right] \\
&+d^2\left[-\frac{81}{32\varepsilon^2}+\frac{9}{2\varepsilon}-\frac{141}{16}\right]+d^3\left[\frac{3}{16\varepsilon^2}-\frac{11}{32\varepsilon}+\frac{5}{8}\right] , \\
W^{(1)} &= \left[-\frac{63-6\xi}{\varepsilon^2}+\frac{3+4\xi}{\varepsilon}+48-6\xi\right]+d\left[\frac{93-6\xi}{2\varepsilon^2}-\frac{149+2\xi}{4\varepsilon}-\frac{28-11\xi}{2}\right] \\
&+d^2\left[-\frac{183-6\xi}{16\varepsilon^2}+\frac{137-\xi}{8\varepsilon}-\frac{33+4\xi}{4}\right]+d^3\left[\frac{15}{16\varepsilon^2}-\frac{2}{\varepsilon}+\frac{29}{16}\right] , \\
W^{(2)} &= \left[-\frac{15}{\varepsilon^2}+\frac{37+6\xi}{2\varepsilon}-\frac{91+8\xi}{2}\right]+d\left[\frac{21}{2\varepsilon^2}-\frac{133+6\xi}{8\varepsilon}+\frac{279+8\xi}{8}\right] \\
&+d^2\left[-\frac{39}{16\varepsilon^2}+\frac{35}{8\varepsilon}-\frac{67}{8}\right]+d^3\left[\frac{3}{16\varepsilon^2}-\frac{11}{32\varepsilon}+\frac{5}{8}\right] . \qquad (9)
\end{aligned}
$$

We recall that for consistency of these asymptotics with the Schwinger–Dyson equation we require fulfillment of the conditions

$$U^{(1)} + W^{(1)} = 0 , \qquad U^{(2)} + W^{(2)} = 0 . \qquad (10)$$

If we substitute the gauge parameter (3) in (9), then irrespective of the value of ξ the poles of second order in ε cancel in each of the expressions (9) (and in the expressions for $W^{(1)}$ and $W^{(2)}$ the poles of first order do so as well), and we obtain the expressions

$$
\begin{aligned}
U^{(1)} = \frac{1}{8}\left[\frac{-12\xi^2+6\xi}{\varepsilon}+51+49\xi-2\xi^2\right] , \qquad & U^{(2)} = \frac{1}{32}\left[\frac{12\xi-6}{\varepsilon}-49-202\xi\right] , \\
W^{(1)} = 12\xi-6 , \qquad & W^{(2)} = -6\xi+3 . \qquad (11)
\end{aligned}
$$

Generally speaking, when dimensional regularization is used the parameters of the problem can depend of the dimension n of space. Therefore, we consider the expansion of ξ with respect to ε, in which the first two terms are important:

$$\xi = \xi_0 + \xi_1\varepsilon + \mathcal{O}(\varepsilon^2) . \qquad (12)$$

Comparing (11) and (12), we see that the choice $\xi_0 = \frac{1}{2}$ eliminates the poles with respect to ε in $U^{(1)}$ and $U^{(2)}$, and, in addition, makes $W^{(1)}$ and $W^{(2)}$ vanish. Substituting $\xi = \frac{1}{2}+\xi_1\varepsilon$, we obtain

$$U^{(1)} = \frac{3}{8}(25-2\xi_1) , \qquad U^{(2)} = \frac{3}{16}(2\xi_1-25) . \qquad (13)$$

It can be seen from this that the same value $\xi_1 = \frac{25}{2}$ makes both expressions vanish. Thus, we have found that in the gauge (3) and for value

$$\xi = \frac{1}{2} + \frac{25}{2}\varepsilon + \mathcal{O}(\varepsilon^2)\ . \tag{14}$$

of the coefficient in front of the transverse part of the three-gluon vertex the condition for cancelling of the leading infrared singularity is satisfied. Note that it is not only the total contribution of the two diagrams that vanishes but also the contributions of each of them separately: $U^{(j)} = W^{(j)} = 0$, $j = 1, 2$.

The vanishing of the $U^{(j)}$ contributions for the value (14) of ξ in gauge (3) can also be seen by calculating the one-loop subdiagram of the first two-loop diagram of Fig. 1:

$$\int \mathrm{d}^n p\ D_{\lambda\lambda'}(p)\ D_{\rho\rho'}(t{-}p)\ \Gamma_{\lambda'\rho'\tau}(p, t{-}p, -t) \sim \frac{1}{t^2}\,(t_\lambda g_{\rho\tau} - t_\rho g_{\lambda\tau})\left[\xi - \frac{1}{2} - \xi\varepsilon - 12\varepsilon + \mathcal{O}(\varepsilon^2)\right]. \tag{15}$$

We retain here the terms of order ε, since in a second integration poles in ε may arise. Note that the tensor structure in the expression (15) remains the same in any other gauge. This is explained by the fact that in the expressions (9) for $U^{(j)}$ there are no terms containing d^4.

To conclude this section, we note that if we consider the system (10), (9) by itself, for arbitrary d, then other solutions can be obtained, for example, $d = 4 - 4\varepsilon + \mathcal{O}(\varepsilon^2)$, $\xi = \frac{7}{2} + \mathcal{O}(\varepsilon)$. However, for this value of the gauge parameter the free solution for the ghost propagator does not satisfy Eq. (1), and we can say nothing about the contribution of the ghosts to the equation for the gluon propagator or about their influence on the Slavnov–Taylor identities.

3. Thus, the verification of the consistency of the infrared asymptotics of the gluon Green's functions (2), (5) and (6) and the Schwinger–Dyson equation permits the following conclusions to be drawn. First, the assumption that in the gauge (3), $d = 4/(5 - n)$, the contributions of the ghosts to the gauge identities are unimportant in the infrared region is self-consistent in the sense that the infrared-singular gluon propagator (2), $D(k) \sim M^2/k^4$, satisfies asymptotically the Schwinger–Dyson equation, on the one hand, and leads to the free solution for the ghost propagator, on the other. In this connection, we should like to emphasize once more the distinguished nature of the gauge (3), whose advantages have already been noticed more than once [12, 6] and which, in particular in the present paper, can be obtained as solution of the system (9).

Also important is the result giving the parameter value $\xi_0 = \frac{1}{2}$ in the gauge (3), this value fixing in our approach the form of the asymptotic behaviour of the gluon vertices. The value obtained for the corresponding parameter in the axial gauge was -1 [2, 9], and this illustrates, in particular, the gauge dependence of the gluon Green's functions.

We emphasize that we have considered here for the first time the two-loop terms of the Schwinger–Dyson equation for the gluon propagator that are the most singular in the infrared region. In earlier studies they were either ignored without any justification or (in the axial gauge) the projection mentioned above, to which they do not contribute, was chosen [1, 2, 8]. It must be borne in mind that if the equation is to be satisfied not only in the leading asymptotic behaviour consistency is also required for the less singular terms (in the axial light-like gauge this problem was considered in [8]). In our case we need cancellation in the following step of the leading asymptotic behaviour of the one-loop diagrams of the equation shown in Fig. 1 with the softer contributions of the two-loop diagrams that derive from the following terms in the infrared expansion of the Green's functions. The practical implementation of this iteration procedure requires even more laborious work.

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