The Infrared Behavior of Gluon and Ghost Propagators in Landau Gauge QCD

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(September 17, 1997)

A solvable systematic truncation scheme for the Dyson–Schwinger equations of Euclidean QCD in Landau gauge is presented. It implements the Slavnov–Taylor identities for the three–gluon and ghost–gluon vertices, whereas irreducible four–gluon couplings as well as the gluon–ghost and ghost–ghost scattering kernels are neglected. The infrared behavior of gluon and ghost propagators is obtained analytically: The gluon propagator vanishes for small spacelike momenta whereas the ghost propagator diverges more strongly than a massless particle pole. The numerical solutions are compared with recent lattice data for these propagators. The running coupling of the renormalization scheme approaches a fixed point, \( \alpha_c \approx 9.5 \), in the infrared.


A theoretical understanding of confinement of quarks and gluons into colorless hadrons could be obtained by proving the failure of the cluster decomposition property for color–nonsinglet gauge–covariant operators. One long established idea in this direction is based on the occurrence of infrared divergences to suppress the emission of colored states from color–singlet states [1]. Such a description of confinement in terms of perturbation theory necessarily has to fail.

Thus, to study the infrared behavior of QCD amplitudes non–perturbative methods are required, and, since divergences are anticipated, a formulation in the continuum is desirable. Both of these are provided by studies of truncated systems of Dyson–Schwinger equations (DSEs), the equations of motion of QCD Green’s functions. Typically, for their truncation, additional sources of information like the Slavnov–Taylor identities, entailed by gauge invariance, are used to express vertex functions in terms of the elementary 2–point functions, i.e., the quark, ghost and gluon propagators. Those propagators can then be obtained as selfconsistent solutions to non–linear integral equations representing a closed set of truncated DSEs. Some systematic control over the truncating assumptions can be obtained by successively including higher n–point functions in selfconsistent calculations, and by assessing their influence on lower n–point functions in this way. At present, even at the level of propagators no complete solution to truncated DSEs of QCD exists. In particular, even in absence of quarks, solutions for the gluon propagator in Landau gauge rely on neglecting ghost contributions [2–5]. While this particular problem is avoided in ghost free gauges such as the axial gauge, in studies of the gluon DSE in this gauge [6], the possible occurrence of an independent second term in the tensor structure of the gluon propagator has so far been disregarded [7]. In fact, if the complete tensor structure of the gluon propagator in axial gauge is taken into account, one arrives at equations of no less complexity than the ghost–gluon system in the Landau gauge [8].

In addition to the prospect of some insight into confinement from studying the infrared behavior of QCD Green’s functions, DSEs have proved to be a highly successful tool in developing a hadron phenomenology that interpolates smoothly between the infrared (non–perturbative) and ultraviolet (perturbative) regimes [9]. In particular, a variety of models for the interactions of quarks mediated by gluons exists, which are very well suited for a dynamical description of chiral symmetry breaking from the DSE of the quark propagator in some analogy to the gap equation in superconductivity [10]. The superficial result of these studies is that for the quark self–energy to reflect a spontaneous breaking of chiral symmetry there has to be some sufficient interaction strength at low energies. Under these circumstances, the dichotomy of the pion as a Goldstone boson emerging from the Bethe–Salpeter equation for quark–antiquark bound states is very well understood and explains the smallness of its mass as compared to all other hadrons.

In this letter we present a simultaneous solution of a truncated set of DSEs for the propagators of gluons and ghosts in Landau gauge. An extension to this selfconsistent framework to include quarks dynamically is possible and subject to further studies. The behavior of the solutions in the infrared, implying the existence of a fixed point at a critical coupling \( \alpha_c \approx 9.5 \), is obtained analytically. The gluon propagator is shown to vanish for small spacelike momenta in the present truncation scheme. This behavior, though in contradiction with many previous DSE studies [11], can be partially understood from the observation that, in our present calculation, the previously neglected ghost propagator assumes an infrared enhancement similar to what was then obtained for the gluon.

Besides all elementary 2–point functions, i.e., the quark, ghost and gluon propagators, the DSE for the gluon propagator also involves the 3– and 4–point ver-
text functions which obey their own DSEs. These equations involve successively higher n–point functions. A first step towards a truncation of the gluon equation is to neglect all terms with 4–gluon vertices. These are the momentum independent tadpole term, an irrelevant constant which vanishes perturbatively in Landau gauge, and explicit 2–loop contributions to the gluon DSE. The latter are subdominant in the ultraviolet and will thus not affect the behavior of the solutions for asymptotically high momenta. In the infrared it has been argued that the singularity structure of the 2–loop terms does not interfere with the one–loop terms [12]. Without contributions from 4–gluon vertices (and quarks) the renormalized equation for the inverse gluon propagator in Euclidean momentum space is given by [13],

\[
D^{-1}_{\mu\nu}(k) = Z_3 D^{-1}_{\mu\nu}(k) + g^2 N_c Z_1 \frac{1}{2} \int \frac{d^4q}{(2\pi)^4} \left( \Gamma_{\mu\rho}(k, -p, q) D_{\alpha\beta}(q) D_{\rho\sigma}(p) \Gamma_{\sigma\nu}(q, -p, -k) \right) - g^2 N_c \tilde{Z}_1 \int \frac{d^4q}{(2\pi)^4} i k_\mu D_G(p) D_G(p) G_\nu(q, p),
\]

(1)

where \( p = k + q, \) \( D^i \) and \( \Gamma^i \) are the tree level propagator and 3–gluon vertex, \( D_G \) is the ghost propagator and \( \Gamma \) and \( G \) are the fully dressed 3–point vertex functions. The equation for the ghost propagator in Landau gauge QCD, without any truncations, is given by

\[
D_G^{-1}(k) = -\tilde{Z}_3 k^2 + g^2 N_c \tilde{Z}_1 \int \frac{d^4q}{(2\pi)^4} i k_\mu D_G(k - q) G_\nu(q, k) D_G(q).
\]

(2)

The renormalized propagators for ghosts and gluons and the renormalized coupling are defined from the respective bare quantities by introducing multiplicative renormalization constants, \( \tilde{Z}_3 D_G \) := \( \tilde{Z}_3 D^0_G \), \( Z_3 D_{\mu\nu} \) := \( Z_3 D^0_{\mu\nu} \) and \( Z_3 g := g_0 \). Furthermore, \( Z_1 = Z_g Z_3^{1/2}, \tilde{Z}_1 = Z_g Z_3^{1/2} \tilde{Z}_3, \) and we use that \( \tilde{Z}_1 = 1 \) in Landau gauge [14]. The ghost and gluon propagators are parameterized by their respective renormalization functions \( G \) and \( Z \),

\[
D_G(k) = -\frac{G(k^2)}{k^2}, \quad D_{\mu\nu}(k) = \left( \delta_{\mu\nu} - \frac{k_\mu k_\nu}{k^2} \right) \frac{Z(k^2)}{k^2}.
\]

(3)

In order to arrive at a closed set of equations for the functions \( G \) and \( Z \), we use a form for the ghost–gluon vertex which is based on a construction from its Slavnov–Taylor identity (STI) neglecting irreducible 4–ghost correlations in agreement with the present level of truncation [15],

\[
G_\mu(p, q) = i q_\mu \left( \frac{G(k^2)}{G(q^2)} + i p_\mu \left( \frac{G(k^2)}{G(p^2)} - 1 \right) \right).
\]

(4)

With this result, we can construct the 3–gluon vertex according to general procedures from previous studies [16],

\[
\Gamma_{\mu\nu\rho}(p, q, k) = \frac{1}{2} A_+(p^2, q^2; k^2) \delta_{\mu\nu} \delta_{\rho\sigma} i(p + q)_\sigma + \frac{1}{2} A_-(p^2, q^2; k^2) \delta_{\mu\nu} \delta_{\rho\sigma} i(p - q)_\sigma + \text{cyclic permutations},
\]

(5)

Some additionally possible terms, transverse with respect to all three gluon momenta, cannot be constrained by its STI and are thus disregarded. For the fermion vertex in QED as constructed from its Ward–Takahashi identity it is well known that additional transverse terms, with the further constraint not to introduce kinematic singularities, are essential for multiplicative renormalizability [17]. Based on this requirement such terms have been obtained explicitly for quenched QED in ref. [18]. Similar constructions for the vertices in QCD are presently not available. However, the full Bose (exchange) symmetry of the 3–gluon vertex alleviates this problem since, combined with the STI, it puts much tighter constraints on this vertex than those obtained for fermion vertices.

Instead of a direct numerical solution of the coupled system of integral equations resulting from the present truncation scheme we use a one–dimensional approximation: For integration momenta \( q^2 < k^2 \) we use the angle approximation replacing \( G((k - q)^2) \to G(k^2) \) and \( Z((k - q)^2) \to Z(k^2) \). Since this preserves the limit \( q^2 \to 0 \), it is suitable for an analytic discussion of the solutions in the infrared. For \( q^2 > k^2 \) we replace all arguments (including the external \( k^2 \)) by the integration momentum \( q^2 \). The justification for this is the weak logarithm–like dependence of \( G \) and \( Z \) at high momenta [19]. The DSEs (1) and (2) then simplify to

\[
\frac{1}{Z(k^2)} = Z_3 + \frac{g^2}{16\pi^2} \left\{ \int_0^{k^2} \frac{dq^2}{k^2} \left( \frac{7}{2} q^4 - \frac{17}{2} q^2 - \frac{9}{8} \right) Z(q^2) G^2(q^2) \left( Z(q^2) G(q^2) + \int_0^{k^2 \ll q^2} \frac{dq^2}{q^2} \frac{7}{8} q^2 - 7 \right) \right\},
\]

(6)

\[
\frac{1}{G(k^2)} = \tilde{Z}_3 - \frac{g^2}{16\pi^2} \left[ \left( \frac{1}{2} \right) Z(k^2) G(k^2) + \int_{k^2}^{\Lambda_{QCD}^2} \frac{dq^2}{q^2} Z(q^2) G(q^2) \right],
\]

(7)
We introduced an $O(4)$-invariant momentum cutoff $A_{0\nu}$ to account for logarithmic ultraviolet divergences which are absorbed by the renormalization constants $Z_1$ and $\bar{Z}_3$. $Z_1$ has to be ultraviolet finite [20]. This is inconsistent with gauge invariance implying $Z_1 = Z_3/\bar{Z}_3$. While this problem, appearing at order $g^4$ in a perturbative expansion, is quite natural for a truncation scheme neglecting explicit 4-gluon couplings at the same order, its remedy could provide information on purely transverse terms in the 3-gluon vertex. For details of the renormalization and the numerical procedure see [20].

To deduce the infrared behavior of the propagators we make the Ansatz that for $x := k^2 \to 0$ the product $Z(x)G(x) \to cx^\kappa$ with $\kappa \neq 0$ and some constant $c$. The special case $\kappa = 0$ leads to a logarithmic singularity in eq. (7) for $x \to 0$ which precludes the possibility of a selfconsistent solution. In order to obtain a positive definite function $G(x)$ for positive $x$ from an equally positive $Z(x)$, as $x \to 0$, we obtain the further restriction $0 < \kappa < 2$. Eq. (7) then yields,

$$G(x) \to \left( g^2 \gamma_0^{G} \left( \frac{1}{\kappa} - \frac{1}{2} \right) \right)^{-1} c^{-1} x^{-\kappa} \Rightarrow (8)$$

$$Z(x) \to \left( g^2 \gamma_0^{G} \left( \frac{1}{\kappa} - \frac{1}{2} \right) \right) c^2 x^{2\kappa}, \quad (9)$$

where $\gamma_0^{G} = 9/(64\pi^2)$ is the leading perturbative coefficient of the anomalous dimension of the ghost field. Using (8) and (9) in eq. (6), we find that the 3-gluon loop contributes terms $\sim x^\kappa$ to the gluon equation for $x \to 0$ while the dominant (infrared singular) contribution $\sim x^{-2\kappa}$ arises from the ghost–loop, i.e.,

$$Z(x) \to g^2 \gamma_0^{G} \frac{9}{4} \left( \frac{1}{\kappa} - \frac{1}{2} \right)^2 \left( \frac{3}{2} \right) \frac{1}{\kappa - 3} + \frac{1}{4\kappa} \right)^{-1} c^2 x^{2\kappa}. \quad (10)$$

Comparing this to (9) we obtain a quadratic equation with a unique solution $\kappa = (61 - \sqrt{1897})/19 \approx 0.92$ for the exponent $\kappa < 2$. The leading behavior of the gluon and ghost renormalization functions is entirely due to ghost contributions. The details of the approximations to the 3-gluon loop have no influence on these considerations. In particular, additional transverse terms of the 3-gluon vertex, free of kinematical singularities, will yield contributions that are even further suppressed in the infrared. Compared to the Mandelstam approximation, in which the 3-gluon loop alone determines the infrared behavior of the gluon propagator and the running coupling in Landau gauge [2–5], this shows the importance of ghosts. The result presented here implies an infrared stable fixed point in the non-perturbative running coupling of our subtraction scheme, defined by

$$\alpha_S(s) = g^2/4\pi Z(s)G^2(s) \to \frac{16\pi}{9} \left( \frac{1}{\kappa} - \frac{1}{2} \right)^{-1} \approx 9.5 \quad (10)$$

for $s \to 0$. This is qualitatively different from the infrared singular coupling of the Mandelstam approximation [5].
contributions in Landau gauge [2–5]. This shows that ghosts are important, in particular, at low energy scales relevant to hadronic observables.

We thank F. Coester, F. Lenz, M. R. Pennington and H. Reinhardt for helpful discussions. This work was supported by DFG under contract Al 279/3-1, by the Graduiertenkolleg Tübingen and the US-DOE, Nuclear Physics Division, contract # W-31-109-ENG-38.

[11] An infrared enhanced gluon propagator was found in Landau gauge in Mandelstam approximation [2–5] as well as in some studies of its simplified axial gauge DSE [6,7].
[13] We use positive definite metric, $g_{\mu\nu} = \delta_{\mu\nu}$. Color indices are suppressed and the number of colors is fixed, $N_c = 3$.
[15] In [20] we derive a Slavnov–Taylor identity for the ghost–gluon vertex from the usual BRS invariance. This together with the symmetry of the ghost–gluon vertex fully determines its form at the present level of truncation. There are no undetermined transverse terms in this case.
[19] A similar assumption underlies the Mandelstam approximation. For $q^2 > k^2$ it can be further justified from a detailed study of the solutions in the ultraviolet [20]. It will nevertheless be important to assess the sensitivity of the results to the modified angle approximation in future.
[25] This is supported by the qualitative similarity of our solutions to the infrared behavior obtained from studies of the influence of a complete gauge fixing by D. Zwanziger, Nucl. Phys. B378, 525 (1992); ibid. B412, 657 (1994).
[26] We are indebted to F. Coester for pointing this out.

![FIG. 1. The numerical result for the gluon propagator from Dyson–Schwinger equations (solid line) compared to lattice data from fig. 3 in [23].](image1)

![FIG. 2. The numerical result for the ghost propagator from Dyson–Schwinger equations (solid line) compared to data from fig. 1 in [24] for the $24^4$ lattice up to $x \approx 1$, and a fit as obtained in ref. [24] for $x \geq 2$.](image2)