Perturbative Contributions to Field Correlators in Gluodynamics.

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Abstract

The cancellation of perturbative contributions to the string tension in gluodynamics in the framework of vacuum field correlators method is shown at the order $O(g^4)$ by explicit calculation. The general pattern of these cancellations at all orders and relation with the renormalization properties of the Wilson loop is discussed.
1 Introduction

Recently the formalism of gauge-invariant field correlators (FC) has proved to be a useful tool in relating vacuum properties to the QCD string parameters and hadron observables. In particular the phenomena of confinement and deconfinement are understood as due to particular terms of FC and the string tension is obtained as an integral over those terms.

On the lattice FC have been measured both for SU(2) and SU(3) gluodynamics and for the full QCD with the four flavours of staggered fermions and the first measurement of FC in the vacuum without cooling was done recently. All these studies refer to the nonperturbative contents of FC, while the perturbative component, suppressed in the cooling process on the lattice, is an admixture important at small distances and seen clearly in [3]-[5]. Analytically only the lowest order contribution $O(g^2)$ had been known till recently, next-to-leading order terms have been found in [7] and [8]. For the bilocal correlator (see (8)) the exact structure found in [7, 8] looks like:

$$\langle \alpha_s F(x) F(0) \rangle \sim \frac{a + b \ln x}{x^4}$$

where $a$ and $b$ are constants. The renormalization properties of FC (and the values of constants $a$ and $b$) cannot be entirely explained by charge renormalization and contain some additional contributions (see discussion in [8]). These results bring several questions, which are important for the whole formalism of FC and which we try to answer below.

Firstly, what are renormalization properties of FC and how they are connected with those of the Wilson loops? Secondly, the $O(g^4)$ contribution to the bilocal correlator $\langle \alpha_s F F \rangle$ formally leads to the (divergent) contribution to the string tension, which physically has no sense and should be cancelled by other terms. What is the exact mechanism of this cancellation? And thirdly, one should see the general pattern of these cancellations at all
The paper is organized as follows: Sect. 2 is devoted to the definitions of the essential ingredients of the formalism, in Sect. 3 the exact relation between quadratic and triple correlators is used to demonstrate the mechanism of cancellation of perturbative contributions to the string tension at the order $O(g^4)$. The Sect. 4 concludes the paper with a discussion of the cancellation for higher orders and the relation between renormalization properties of Wilson loops and FC.

2 General definitions

We start with the nonabelian Stokes theorem [9, 10] for the Wilson loop average $W(C)$:

$$W(C) = \frac{1}{N_c} \langle \text{Tr} P \exp \left( ig \int_C A_\mu dx^\mu \right) \rangle =$$

$$= \frac{1}{N_c} \langle \text{Tr} P \exp \left( ig \int_S F_{\mu\nu}(z, x_0) d\sigma_{\mu\nu}(z) \right) \rangle$$  \hspace{1cm} (1)

Here appears the basic quantity of the FC method - the field strength operator $F_{\mu\nu}(z)$ covariantly transported with the help of the operators

$$\Phi(z, x_0) = P \exp \left( ig \int_{x_0}^{z} A_\mu(u) du_\mu \right)$$

to some chosen reference point $x_0$.

$$F_{\mu\nu}(z, x_0) = \Phi(x_0, z) F_{\mu\nu}(z) \Phi(z, x_0)$$  \hspace{1cm} (2)

The averaging process, denoted in (1) by angular brackets is the standard integration in the QCD partition function, containing gauge fixing and ghost terms. For our purposes we neglect quark degrees of freedom and assume the
perturbative expansion of the partition function, yielding perturbative series for \( W(C) \) and FC.

The cluster expansion theorem for (1) reads:

\[
W(C) = \frac{1}{N_c} \text{Tr} \exp \left( \sum_{n=1}^{\infty} \frac{(ig)^n}{n!} \int d\sigma(1) \cdots d\sigma(n) \langle \langle F(1) \cdots F(n) \rangle \rangle \right)
\]

(3)

where we have suppressed the indices, \( F(k) = F_{\mu\nu}(z_k, x_0) \), and we have used irreducible cumulants instead of averages, denoting them with the double angular brackets \([11]\). Note also that cumulants are unit matrices in colour space and ordering operator in (3) is not needed in contrast to (1).

Since \( \langle \langle F(k) \rangle \rangle = 0 \), one can rewrite (3) identically as

\[
W(C) = \frac{1}{N_c} \text{Tr} \exp \left( -\frac{1}{2} \int_{S} \int_{S} d\sigma\mu\nu(u)d\sigma\rho\sigma(v)\Lambda_{\mu\nu,\rho\sigma}(u, v, C) \right)
\]

(4)

where we have defined the global correlator \( \Lambda(u, v, C) \),

\[
\Lambda_{\mu\nu,\rho\sigma}(u, v, C) \equiv g^2 \langle \langle F_{\mu\nu}(u, x_0)F_{\rho\sigma}(v, x_0) \rangle \rangle -
\]

\[
-2 \sum_{n=3}^{\infty} \frac{(ig)^n}{n!} \int d\sigma(3) \cdots d\sigma(n) \left[ \langle \langle F_{\mu\nu}(u, x_0)F_{\rho\sigma}(v, x_0)F(3) \cdots F(n) \rangle \rangle + \text{perm.}(1, 2, \ldots n) \right]
\]

(5)

and \( \text{perm.}(1, 2, \ldots n) \) stands for the sum of terms with different ordering of \( F(u, x_0) = F(1) \) and \( F(v, x_0) = F(2) \) with respect to all other factors \( F(k) \).

Since \( W(C) \) does not depend on the shape of the surface \( S \), the dependence of the global correlator on its arguments is such, that r.h.s. of (4) is independent of the choice of \( S \) too but depends on the contour \( C \). This circumstance explains the name ”global correlator” used for \( \Lambda(u, v, C) \) in contrast to local correlators which enter in the r.h.s. of (3) (note however, that the name ”local” should not be misunderstood - correlators \( \langle \langle F(1) \cdots F(k) \rangle \rangle \) depend on the points \( z_1, \ldots, z_k \) as well as on the paths, entering in the definition (2) via transporters \( \Phi(z, x_0) \)).
Let us now fix one of the integration points in the exponent of (4) and denote the rest integral as follows

\[ Q_{\mu\nu}(u, C) \equiv \frac{1}{2} \int_S d\sigma_{\rho\sigma} \Lambda_{\mu\nu,\rho\sigma}(u, v, C) \]  

(6)

In the confining phase one expects for large contours \( C \) the minimal area law of Wilson loop, which implies that \( Q_{\mu\nu} \) in this limit does not depend on the point \( u \) when \( S \) is the minimal area surface and simply coincides with the string tension \( \sigma \), while for the arbitrary surface one can identify \( Q_{\mu\nu} \) as

\[ Q_{\mu\nu}(u, C) = P_{\mu\nu} \cdot \sigma \]  

(7)

where \( P_{\mu\nu} \) projects onto the minimal surface. Conversely if \( Q_{\mu\nu} \) does not have constant limit for large \( S \) then the area law of Wilson loop does not hold.

To calculate \( Q_{\mu\nu}(u, C) \) one can for simplicity take \( x_0 \) in (4) to coincide with \( u \). Then the lowest order FC in (5) depends only on two points (and on the straight line, connecting them.) In what follows we concentrate on contributions to \( \sigma \) and therefore take for simplicity a planar contour \( C \) with the minimal surface \( S \) lying in the plane.

The exact form of the two-point correlator may be written in the following way [2]:

\[ D_{\mu\nu,\rho\sigma}(u - v) = Tr\{gF_{\mu\nu}(u)\Phi(u, v)gF_{\rho\sigma}(v)\Phi(v, u)\} = \]

\[ = \left( D(z^2) + D_1(z^2) + \frac{z^2}{2} \frac{dD_1(z^2)}{dz^2} \right) \Delta_{\mu\nu,\rho\sigma}^{(1)} - \frac{z^2}{2} \frac{dD_1(z^2)}{dz^2} \Delta_{\mu\nu,\rho\sigma}^{(2)} \]  

(8)

where two tensor structures

\[ \delta_{\mu\rho}\delta_{\nu\sigma} - \delta_{\nu\rho}\delta_{\mu\sigma} = \Delta_{\mu\nu,\rho\sigma}^{(1)} \]  

(9)

and

\[ \Delta_{\mu\nu,\rho\sigma}^{(1)} - 2 \left( \frac{z^2}{2} \delta_{\nu\sigma} \delta_{\mu\rho} - \frac{z^2}{2} \delta_{\mu\rho} \delta_{\nu\sigma} + \frac{z^2}{2} \delta_{\mu\sigma} \delta_{\nu\rho} - \frac{z^2}{2} \delta_{\nu\rho} \delta_{\mu\sigma} \right) = \Delta_{\mu\nu,\rho\sigma}^{(2)} \]  

(10)
were introduced. Note, that

\[ \Delta^{(2)}_{\rho\sigma\mu\nu} \delta_{\mu\rho} \delta_{\nu\sigma} = 0 \]

therefore only the part proportional to \( \Delta^{(1)} \) contributes to the condensate \( \langle \alpha s F_{\mu\nu} F_{\rho\sigma} \rangle \).

It was shown in [2], that the correlator \( D_1(z) \) does not contribute to the string tension, while the contribution of \( D \) is

\[ \sigma^{(2)} = \frac{1}{2} \int d^2 v D(u - v) \]

where the subscript \( (2) \) refers to the quadratic correlator, so that the total contribution of \( \Lambda_{\mu\nu,\rho\sigma} \) can be written as the sum over contributions of correlators of order \( n \),

\[ \sigma = \sum_{n=2}^{\infty} \sigma^{(n)} \]

The perturbative studies of [7, 8] have revealed that the lowest order contribution to \( D(z) \) occurs at the \( O(g^4) \) order. This implies that \( \sigma^{(2)} \) is nonzero at this order (actually it diverges), contrary to physical expectations. It will be shown in the next section that there is another term at the same order of perturbation theory which exactly cancels \( \sigma^{(2)} \). For that purpose one needs some relation between quadratic and triple correlators. The relation of this type, namely the exterior derivative of the function \( D(z) \) from (8) expressed through the triple correlators was found in [12]:

\[ \varepsilon_{\mu_1\nu_1\rho\sigma} \frac{d D(z^2)}{d z^2} = \frac{i}{4} \varepsilon_{\mu_2\nu_2\xi\rho} \left( < Tr(F_{\mu_1\nu_1}(z_1) \tilde{I}_{\sigma\xi}(z_1, z_2) F_{\mu_2\nu_2}(z_2) \Phi(z_2, z_1)) > - < Tr(F_{\mu_1\nu_1}(z_1) \Phi(z_1, z_2) F_{\mu_2\nu_2}(z_2) I_{\sigma\xi}(z_2, z_1)) > \right) \]  \hspace{1cm} (11)

where

\[ \tilde{I}_{\rho\gamma}(z, z') = \int_0^1 d\alpha \alpha \Phi(z, z + \alpha(z' - z)) F_{\rho\gamma}(z + \alpha(z' - z)). \]
\[ \Phi(\alpha(z' - z), z') \] (12)

and analogously

\[ I_{\rho\gamma}(z, z') = \int_0^1 d\alpha \cdot \alpha \Phi(z', \alpha(z' + \alpha(z - z')) F_{\rho\gamma}(z', \alpha(z - z')). \]

\[ \cdot \Phi(z' + \alpha(z - z'), z') \] (13)

We have taken into account the Bianchi identity \( \varepsilon_{\mu_2\nu_2\xi_\rho} D_{\xi} F_{\mu_2\nu_2}(z) = 0 \) and denoted \( z_2 - z_1 = z \). These relations will be important in what follows.

In the framework of the described formalism it is natural to separate perturbative and nonperturbative contributions to the functions \( D(z) \) and \( D_1(z) \) and take them into account differently for different processes. We are concentrating in the present paper on perturbative parts of the bilocal and higher correlators to explain several specific features the perturbation theory has in field strength formulation.

### 3 Cancellation of the perturbative contributions to the string tension at the order \( O(g^4) \).

It has already been mentioned, that the results of [7, 8] imply that scalar functions \( D(z) \) and \( D_1(z) \) both receive the perturbative contributions at the order \( O(g^4) \) while at the tree level only \( D_1(z) \) is nonzero. The absence of perturbative contributions to the function \( D(z) \) (and therefore to the string tension) at the tree level in \( SU(2) \) gluodynamics was also noticed in different respect in [13].

To look for cancellation at the given order \( O(g^4) \) one must identify all terms of this order in \( \Lambda(u, v, C) \) and \( Q_{\mu\nu}(u, C) \). The \( O(g^4) \) contribution comes from the quadratic and triple terms in (3) which we write in ”polar” coordinates, \( u = s_1z_1, v = s_2z_2, 0 \leq s_i \leq 1 \):

\[ \Lambda_{\nu\rho\mu\sigma} = g^2 \langle \langle F_{\nu\rho}(s_1z_1, x_0) F_{\mu\sigma}(s_2z_2, x_0) \rangle \rangle + \]
\begin{equation}
+ \int dz_3^\phi \int_0^1 ds_3 s_3 \langle \{ F_{\nu\rho}(s_1 z_1, x_0) F_{\mu\sigma}(s_2 z_2, x_0) F_{\xi\phi}(s_3 z_3, x_0) \} \rangle + (14)
\end{equation}

The term \( O(\langle \langle FFFF \rangle \rangle) \) starts from the quartic cumulant and is \( O(g^6) \), therefore it does not contribute to the function \( \Lambda_{\nu\rho\mu\sigma} \) at the \( g^4 \)-order we are interested in at the moment. According to (6) we need to calculate

\begin{equation}
Q_{\nu\rho} = \frac{1}{2} \int dV_\beta \epsilon_{\mu\sigma\kappa\beta} \frac{\partial}{\partial \nu^\kappa} \Lambda_{\nu\rho\mu\sigma}(u, v)
\end{equation}

and show this quantity to be equal to zero at the desired order. From the Stokes theorem point of view it means disappearance of the area term in the Wilson loop.

Since bilocal and triple cumulants coincide with the usual correlators due to \( \langle F_{\mu\nu}(z, x_0) \rangle = 0 \) one gets (omitting for simplicity of notation the reference point \( x_0 = 0 \) and phase factors \( \Phi(x_0, z) \) in all correlators):

\begin{equation}
\epsilon_{\mu\sigma\kappa\beta} \frac{\partial}{\partial z_2^\kappa} \left( \langle F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) \rangle + \right.
\end{equation}

\begin{equation}
+ \int dz_3^\phi \int_0^1 ds_3 s_3 \langle F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\xi\phi}(s_3 z_3) \rangle \bigg) = \epsilon_{\mu\sigma\kappa\beta} (L_{DF} + L_3 + L_4)_{\kappa\nu\rho\mu\sigma}
\end{equation}

where we have denoted

\begin{equation}
L_{DF} = \langle F_{\nu\rho}(s_1 z_1) D_{\kappa} F_{\mu\sigma}(s_2 z_2) \rangle
\end{equation}

\begin{equation}
L_3 = \langle F_{\nu\rho}(s_1 z_1) (s_2)^2 z_2^\gamma \int_0^1 d\alpha F_{\gamma\kappa}(\alpha s_2 z_2) F_{\mu\sigma}(s_2 z_2) \rangle - \langle F_{\nu\rho}(s_1 z_1) (s_2)^2 z_2^\gamma \int_0^1 d\alpha F_{\rho\sigma}(s_2 z_2) F_{\gamma\kappa}(\alpha s_2 z_2) \rangle +
\end{equation}
\[ L_4 = \int dz_3 \int ds_3 s_3 \gamma^2 (F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(s_3 z_2)) + \]

\[ + \int ds_3 z_2 \gamma (F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(s_3 z_2)) + \]

\[ + \int ds_3 s_3 \gamma^2 (F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(s_3 z_2)) + \]

\[ + \int du \gamma (F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(s_3 z_2)) - \]

\[ - \int du \gamma (F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(s_3 z_2)) \]

(17)

The phase factors have been differentiated according to [14]. The terms \( L_{DF} \) containing \( D_{\mu} \tilde{F}_{\mu\sigma} \) vanish because of Bianchi identity. The terms of the order \( \langle F^3 \rangle \) may be rewritten as (in the radial gauge for simplicity, \( \Phi(0, x) = 1 \)):

\[ \epsilon_{\mu\sigma\kappa\beta} z_2^\beta \left( \int u du \langle F_{\nu\rho}(s_1 z_1) [F_{\gamma\kappa}(u z_2) F_{\mu\sigma}(s_2 z_2)] \rangle + \right) \]

\[ \int u du \langle F_{\nu\rho}(s_1 z_1) F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(u z_2) \rangle \]

(18)

At the \( O(g^4) \)-order one can easily observe antisymmetric color structure of the tree-point correlator:

\[ \langle F^a(x) F^b(y) F^c(z) \rangle \propto f^{abc} D_3 \]

(19)

The above expression will be used to demonstrate vanishing of (18).

Taking into account the identity \( f^{abc} t^a t^b t^c = i/4 (N^2 - 1) \tilde{1} \) and performing the symmetrization with respect to integrations over \( s_2 \) and \( u \) one obtains that the integrand in (18) is proportional to

\[ \langle F_{\nu\rho}(s_1 z_1) \{ F_{\mu\sigma}(s_2 z_2) F_{\gamma\kappa}(u z_2) \} \rangle \]

where \{..\} denote anticommutator. This average is zero due to (13).
Situation with the $\langle F^4 \rangle$ terms in (17) is simpler, since only disconnected parts of the quartic correlator contribute at the $g^4$ order. Hence phase factors may be omitted and the correlator is factorized:

$$
\langle F_1 F_2 F_3 F_4 \rangle = \langle F_1 F_2 \rangle \langle F_3 F_4 \rangle + \langle F_1 F_3 \rangle \langle F_2 F_4 \rangle + \langle F_1 F_4 \rangle \langle F_2 F_3 \rangle
$$

It is easy to observe, that two last terms in (17) cancel each other at this order. This finishes the proof of the stated cancellation at the $g^4$ order.

## 4 Renormalization properties of the Wilson loops and field correlators.

The Wilson loop renormalization properties were studied in [15, 16] and for smooth contour $C$ the result is:

$$
W(C) = Z \cdot W_{\text{ren}}(C) \quad (20)
$$

where the (infinite) $Z$-factor contains linear divergencies arising from the integrations over the contour while all logarithmic divergencies are absorbed into the renormalized charge $g_{\text{ren}}(\mu)$ defined at the corresponding dynamical scale $\mu$.

To connect the property (20) with FC one can write the perturbative series for $W(C)$ in the form of the cluster expansion:

$$
W(C) = \frac{1}{N_c} \langle \text{Tr} \ Pexp \left( ig \int_C A_\mu dx^\mu \right) \rangle =
$$

$$
= \frac{1}{N_c} \text{Tr} \ exp \left( -\frac{1}{2} \int_C \int_C dz_\mu du_\nu A_{\mu\nu}(z, u, C) \right) \quad (21)
$$

where $A_{\mu\nu}(z, u, C)$ is defined as (note an analogy with the definition of $\Lambda_{\mu\nu,\rho\sigma}$ in (5))

$$
A_{\mu\nu}(z, u, C) \equiv g^2 \langle \langle A_\mu(z) A_\nu(u) \rangle \rangle
$$
\[-2 \sum_{n=3}^{\infty} \frac{(ig)^n}{n!} \int dz(3) \ldots dz(n) \left[ \langle A_\mu(z)A_\nu(u)A(3) \ldots A(n) \rangle + \text{perm.} \right] \tag{22}\]

It is clear that \( \langle A_\mu(z) \rangle = 0 \) and the ordering operator \( P \) is not needed in (21) due to the color neutrality of the vacuum.

One can now use one of the coordinate gauges [10] and connect \( A_\mu \) and \( F_{\mu\nu} \), in the simplest Fock-Schwinger gauge one has:

\[ A_\mu(x) = \int_0^1 s x_\nu F_{\nu\mu}(s) ds \]

As a consequence \( A_{\mu\nu} \) is expressed through FC as follows:

\[ A_{\mu\nu}\sigma(z, t, C) = \int \int \frac{\partial u_\mu}{\partial z_\nu} \frac{\partial v_\nu}{\partial t_\sigma} \Lambda_{\mu\nu,\rho\sigma}(u, v, C) du_\rho dv_\lambda \tag{23} \]

Now following the procedure of [15, 16] and comparing the perturbation series for \( A_{\mu\nu} \) and \( \Lambda_{\mu\nu,\rho\sigma} \) one can see that only those terms in \( \Lambda_{\mu\nu,\rho\sigma} \) which take the form of the full derivatives in \( u, v \) have counterparts in \( A_{\mu\nu} \) (i.e. the terms \( D_1 \) and the similar structures for higher correlators) while the Kronecker type terms (proportional to \( \Delta^{(1)} \)) are cancelled since these terms are not present in \( A_{\mu\nu} \). It is also clear, that all contributions arising from perturbative expansion of parallel transporters in the gauge-covariant definition (2) are exactly cancelled at each given order \( O(g^n) \) in \( A_{\mu\nu} \) (see also discussion of the related points in [8]). Hence perturbative contributions to the string tension are cancelled at any finite order \( O(g^n) \) as well as those logarithmic contributions to the coupling constant renormalization which arise from the phase factors’ perturbative expansion. Our explicit calculation in Section 3 is the demonstration of this general statement in the special case \( n = 4 \).

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References

[1] H.G. Dosch, Progr.Part.Nucl.Phys. 33 (1994) 121.
   A. DiGiacomo, H.G. Dosch, Yu.A. Simonov, in preparation.
   Yu.A. Simonov, Physics-Uspekhi 39 (1996) 313;

[2] H.G. Dosch, Phys.Lett. B190 (1987) 177.
   H.G. Dosch, Yu.A. Simonov, Phys.Lett. B205 (1988) 339.

[3] A. DiGiacomo, H. Panagopoulos, Phys.Lett. B 285 (1992) 133.

[4] A. DiGiacomo, E. Meggiolaro, H. Panagopoulos, Nucl.Phys.Proc.Suppl. 54 A (1997) 343.

[5] M. D’Elia, A. DiGiacomo, E. Meggiolaro, hep-lat/9705032.

[6] G. Bali, N. Brambilla, A. Vairo, hep-lat/9709079.

[7] M. Eidemüller, M. Jamin, Phys.Lett.B 416 (1998) 415.

[8] V.I. Shevchenko, hep-ph/9802274.

[9] M. Halpern, Phys.Rev.D19 (1979) 517.
   Y. Aref’eva, Theor.Math.Phys. 43 (1980) 353;
   N.Bralic, Phys.Rev. D22 (1980) 3090;
   Yu.A. Simonov, Sov.J.Nucl.Phys. 50 (1989) 134;
   M.Hirayama, S.Matsubara, hep-th/9712120

[10] V.I. Shevchenko, Yu.A. Simonov, hep-th/9802134

[11] N.G. Van Kampen, Stochastic Processes in Physics and Chemistry,
    North-Holland Physics Publishing, 1984.
[12] Yu.A. Simonov, Yad.Phys. 50 (1987) 213.
V.I. Shevchenko, Yu.A. Simonov, Phys.of Atom.Nucl., 60 (1997) 1201.

[13] D. Antonov, D. Ebert, Mod.Phys.Lett. A 12 (1997) 2047.

[14] Yu.M. Makeenko, A.A. Migdal, Nucl.Phys, B188 (1981) 269.

[15] V.S. Dotsenko, S.N. Vergeles, Nucl.Phys. B169 (1980) 527.

[16] R.A. Brandt, F. Neri, M.-A. Sato, Phys.Rev. D24 (1981) 879.