Relations between correlators in gauge field theory

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Abstract

Exact relations between gauge-invariant vacuum correlators in QCD are derived. Derivatives of the correlators are expressed in terms of higher orders correlators. The behaviour of the correlators at large and small distances due to these relations is discussed.

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Understanding and description of nonperturbative interactions is one of the most important problems in the modern field theory. Some progress on this way has been achieved with the method of vacuum correlators (MVC) [1, 2, 3]. Nonperturbative interactions are expressed in this method through gauge-invariant correlators of the gluon field strength operators, acting at the different space-time points. All gauge-invariant quantities like, for example, string tension and hadron masses may be written in terms of these correlators [3]. On the other hand, the latter can be computed on the lattice [4], and thus play a role in the theory as intermediate step for representing the dynamical information.

To have a self-consistent theoretical scheme for MVC, it is necessary to derive the equations for correlators, or, at least, some relations between different correlators.

Concerning the first, a set of exact equations has been proposed [5], which obtained by the stochastic quantization method, but, unfortunately they look quite difficult to be solved and there are still no physical results derived in this framework.

We choose the second way in this paper and derive simple relations, yielding some information about the behaviour of the correlators at small and large distances. In the previous paper [6] the equation for 2-point and 3-point correlators was found. The analysis of this equation will be continued here and the higher order correlators - 4-point ones are also included.

We are going to study gauge-invariant Green functions (n-point correlators) of the following form:

\[
\Delta_{\mu_1\nu_1,\ldots,\mu_n\nu_n}^{(n)}(z_1,\ldots,z_n) = \langle Tr(F_{\mu_1\nu_1}(z_1)\Phi(z_1,z_2)F_{\mu_2\nu_2}(z_2)\ldots F_{\mu_n\nu_n}(z_n)\Phi(z_n,z_1)) \rangle
\]

(1)

here \(F_{\mu\nu}(z) = F_{\mu\nu}^a(z)t^a\) is gluon field strength and nonabelian phase factors (parallel transporters) \(\Phi(z,z') = P \exp i \int_{z}^{z'} A_\mu dx^\mu\) - are crucial for gauge-invariance of the correlators. Vacuum averages are defined as follows:

\[
\langle O(A) \rangle = \int DA \, e^{-\frac{1}{2g^2}\int d^4x Tr(F_{\mu\nu}F_{\mu\nu})} \, O(A)
\]

Any \(n\)-point correlator may be constructed using invariant tensors \(\epsilon_{\mu_\nu\sigma\rho}, \delta_{\mu\nu}\), vectors \((z_i - z_j)_\mu\) and a set of scalar functions, depending on relative coordinates. For the simplest nontrivial 2-point correlator one has two scalar
functions \([\Pi]\):
\[
\Delta_{\mu_1\nu_1,\mu_2\nu_2}^{(2)} = < Tr(F_{\mu_1\nu_1}(z_1)\Phi(z_1, z_2)F_{\mu_2\nu_2}(z_2)\Phi(z_2, z_1)) > = \\
= \frac{1}{2} \left( \frac{\partial}{\partial z_{\mu_1}} (z_{\mu_2}\delta_{\nu_1\nu_2} - z_{\nu_2}\delta_{\mu_1\mu_2}) + \frac{\partial}{\partial z_{\nu_1}} (z_{\mu_2}\delta_{\mu_1\nu_2} - z_{\nu_2}\delta_{\mu_1\mu_2}) \right) D_1(z_1 - z_2) + \\
+ \frac{1}{2} \varepsilon_{\mu_1\nu_1ab}\varepsilon_{\mu_2\nu_2ab} D(z_1 - z_2)
\]

Function \(D(z)\) as well as \(D_1(z)\) may be written as a sum of two parts:
\[
D(z) = D_{\text{pert}}(z) + D_{\text{np}}(z)
\]

The first part \(D_{\text{pert}}(z) \to \infty\) if \(z^2 \to 0\) and physically represents the exchange of perturbative gluons. Notice here, that one-gluon exchange does not get contribution to the \(D_{\text{pert}}(z)\), the contribution to \(D_{\text{pert}}(z)\) is of the following form:
\[
D_{\text{pert}}(z) = \frac{16\alpha_s(z^2)}{3(z^2)^2} + 0(\alpha_s^2).\]

At the same time for the theories with a nontrivial vacuum like QCD, there is the second part \(D_{\text{np}}(z)\), which has no perturbative interpretation. Lattice simulations show \([\Pi]\), that the function \(D_{\text{np}}(z)\) for \(|z| \gtrsim 0.2\, fm\) falls exponentially:
\[
D_{\text{np}}(z) \sim e^{-\frac{|z|}{T_g}}
\]
with the correlation length \(T_g = 0.2\, fm\), the behaviour of \(D_{\text{np}}(z)\) for smaller \(z\) is still unknown. At the point \(z = 0\) nonperturbative components are expressed through the gluon condensate \([\Pi]\):
\[
\Delta_{\mu_1\nu_1,\mu_2\nu_2}^{(2)}(0) = < Tr(F_{\mu_1\nu_1} F_{\mu_2\nu_2}) > = \\
= \frac{1}{2} \varepsilon_{\mu_1\nu_1ab}\varepsilon_{\mu_2\nu_2ab} \cdot (D(0) + D_1(0))
\]
connect neighbouring points along the chosen trajectories. We will consider
 correlators with the points, lying on a fixed straight line, this choice is natural
for 2–point correlator and we follow it for higher correlators too.

The fields under consideration are nonabelian, that’s why the derivative
of the correlator with respect to \( z_\sigma \) is not equal to the correlator of
\( D_\sigma F_{\mu\nu}(z) \) with another fields. The difference arises from the differentiation of the
contour and it is crucial for the nontriviality of the picture. Nonabelian
transporters

\[
\Phi(z, z') = P \exp i \int_z^{z'} A_\mu dx^\mu \tag{6}
\]

are differentiated according to the rule \[10\]:

\[
\frac{\partial \Phi(z, z')}{\partial z'_\gamma} = i\Phi(z, z')A_\gamma(z') + i(z' - z)_\rho \tilde{I}_{\rho\gamma}(z, z') \tag{7}
\]

where notation is used:

\[
\tilde{I}_{\rho\gamma}(z, z') = \int_0^1 d\alpha \ \alpha \Phi(z, z + \alpha(z' - z)) F_{\rho\gamma}(z + \alpha(z' - z)) \cdot \Phi(z + \alpha(z' - z), z') \tag{8}
\]

Analogously

\[
\frac{\partial \Phi(z, z')}{\partial z_\gamma} = -iA_\gamma(z)\Phi(z, z') + i(z' - z)_\rho I_{\rho\gamma}(z, z') \tag{9}
\]

\[
I_{\rho\gamma}(z, z') = \int_0^1 d\alpha \cdot \alpha \Phi(z, z' + \alpha(z - z')) F_{\rho\gamma}(z' + \alpha(z - z')) \cdot \Phi(z' + \alpha(z - z'), z') \tag{10}
\]

We shall often omit the arguments of transporters below for the sake of
simplicity. Let us consider the derivative of the 2–point correlator (2), using
(7) and (9) one has:

\[
\frac{\partial}{\partial z_{2\xi}} \Delta^{(2)}_{\mu_1\nu_1;\mu_2\nu_2} = < Tr(F_{\mu_1\nu_1}(z_1)\Phi D_{\xi} F_{\mu_2\nu_2}(z_2)\Phi) > + \]

\[
i(z_2 - z_1)_\sigma \left( < Tr(F_{\mu_1\nu_1}(z_1)\tilde{I}_{\sigma\xi}(z_1, z_2) F_{\mu_2\nu_2}(z_2)\Phi(z_2, z_1)) > - \right)
\]

4
\[ - < \text{Tr}(F_{\mu_1 \nu_1}(z_1) \Phi(z_1, z_2) F_{\mu_2 \nu_2}(z_2) I_{\sigma \xi}(z_2, z_1)) > \] (11)

The l.h.s. of this equation is:

\[ \varepsilon_{\mu_2 \nu_2 \xi \rho} \frac{\partial}{\partial z_2} \Delta^{(2)}_{\mu_1 \nu_1 \mu_2 \nu_2} = 4 \varepsilon_{\mu_1 \nu_1 \xi \rho} \frac{dD(z)}{dz^2} z_\xi \] (12)

In such a manner one obtains a connection between the 2–point and 3-point correlators:

\[ \varepsilon_{\mu_1 \nu_1 \sigma \rho} \frac{dD(z)}{dz^2} = \frac{i}{4} \varepsilon_{\mu_2 \nu_2 \xi \rho} \left( < \text{Tr}(F_{\mu_1 \nu_1}(z_1) \tilde{I}_{\sigma \xi}(z_1, z_2) F_{\mu_2 \nu_2}(z_2) \Phi(z_2, z_1)) > - < \text{Tr}(F_{\mu_1 \nu_1}(z_1) \Phi(z_1, z_2) F_{\mu_2 \nu_2}(z_2) I_{\sigma \xi}(z_2, z_1)) > \right) \] (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 \( z_2 - z_1 = z \).

This relation is true for the perturbative as well as for nonperturbative parts of the correlators. We will focus our attention on the latter, assuming regular behaviour if their arguments tend to zero. The expression (13) is simplified if \( z \to 0 \), because in this case \( \tilde{I}_{\sigma \xi}(z_1, z_1) = I_{\sigma \xi}(z_1, z_1) = \frac{1}{2} F_{\sigma \xi}(z_1) \).

Thus one obtains:

\[ \left. \frac{dD(z)}{dz^2} \right|_{z=0} = \frac{f^{abc}}{96} < F^a_{\mu_1 \nu_1} F^b_{\mu_1 \nu_2} F^c_{\nu_2 \mu_1} > . \] (14)

Relations (13) and (14) have been derived in [6]. Let us take up a question, what tensor structures of the 3–point correlator contribute to the r.h.s. of (14). There are two independent Croneker tensors in this case, namely:

\[ \Delta^{(3)}_{\mu_1 \nu_1, \mu_2 \nu_2, \mu_3 \nu_3} = -i < \text{Tr}(F_{\mu_1 \nu_2} \Phi F_{\mu_2 \nu_2} \Phi F_{\mu_3 \nu_3} \Phi) >= \]

\[ = \frac{1}{6} \varepsilon_{\mu_1 \nu_1 a b c} \varepsilon_{\mu_2 \nu_2 b c} \varepsilon_{\mu_3 \nu_3 a} \cdot D_3 + \]

\[ + \frac{1}{6} \left( \varepsilon_{\mu_1 \nu_1 a} \varepsilon_{\mu_3 \nu_3 b} - \varepsilon_{\mu_1 \nu_2 a} \varepsilon_{\mu_3 \nu_3 b} + \text{permut.} \right) \cdot D_2 + \]

\[ + \text{nonkroneker terms.} \] (15)

It is easy to see, that the part, proportional to \( D_3 \) does not contribute, while \( D_2 \) does:

\[ f^{abc} < F^a_{\mu_1 \nu_1} F^b_{\nu_1 \nu_2} F^c_{\nu_2 \mu_1} > = 48D_2(0). \] (16)
Formulas (16), (17) have a simple interpretation in terms of the dual Meissner effect. Introducing three-dimensional notations for chromoelectric and chromomagnetic fields $F^a_0 = E^a_\alpha$, $F^a_\alpha = \varepsilon_{\alpha\beta\gamma} H^a_\gamma$, left hand side of (16) may be presented as:

$$ f^{abc} < F^a_{\mu_1\nu_1} F^b_{\mu_2\nu_2} F^c_{\mu_3\nu_3} >= f^{abc} \varepsilon_{\alpha\beta\gamma} (3 < E^a_\alpha E^b_\beta H^c_\gamma > + < H^a_\alpha H^b_\beta H^c_\gamma >) $$

Thus a nonzero vacuum averages of triple correlators imply, that spontaneous magnetic fluxes may be created due to electric or magnetic lines splitting. The existence of noncorrelated magnetic fluxes of this type is equivalent to the presence of monopole currents and is the key point for the possibility of confinement in such vacuum. Notice also, that colour structure of the 3-point condensate is determined by antisymmetric constants $f^{abc}$, symmetric colour term is absent:

$$ d^{abc} < F^a_{\mu_1\nu_1} F^b_{\mu_2\nu_2} F^c_{\mu_3\nu_3} >= 0 $$

or

$$ < F^a_{\mu_1\nu_1} F^b_{\mu_2\nu_2} F^c_{\mu_3\nu_3} >= \frac{f^{abc}}{6} \Delta^{(3)}_{\mu_1\nu_1,\mu_2\nu_2,\mu_3\nu_3}(0) $$

The equation (13) will be examined here in more details. One can prove, that it is equivalent to:

$$ z^a \frac{dD(z)}{dz^2} = \frac{1}{2} \int_0^z t dt \left( D_2(z - t, t) + D_2(-t, t - z) \right) $$

Let us make an assumption, that function $D$ and $D_2$ may be parametrized, taking into account their cyclic symmetry, as:

$$ D(z) = D(0) h(z) h(-z); \ D_2(z - t, t) = D_2(0) h_2(z - t) h_2(t) h_2(-z) $$

Then, it is easy to check, that the solutions of the following equation satisfy also (19):

$$ h(-z) \frac{dh(z)}{dz} = -\frac{a^2}{4} h_2(-z) \int_0^z dt \ h_2(t) \cdot h_2(z - t) $$
where \( a^2 = -\frac{2D_2(0)}{D(0)} \). Naively one can think, that \( h(z) = h_2(z) \) for all \( z \).

Equation (20) may be solved in this case and the regular at \( z = 0 \) solution has the form:

\[
h(z) = 2 \frac{J_1(az)}{az}
\]

where \( J_1(az) \) - Bessel function of the first type. This behaviour contradicts the lattice measurements, it means, that the relation \( h(z) \equiv h_2(z) \) does not hold in the real gluodynamic vacuum. Still this solution leads to the definite value of the string tension, \( \sigma \):

\[
\sigma = \int_0^\infty d^2z D(z) = \frac{2\pi (D(0))^2}{D_2(0)}
\]

The important point here is the negative sign of the fraction \( \frac{D_2(0)}{D(0)} \), i.e. \( a \) should be real, overwise correlation functions would unphysically increase with a distance. Indeed, there are several grounds to think, (see [4] and ref. in [4]) that the 3-point condensate is negative in the real world.

Let us consider a more realistic case \( h(z) \neq h_2(z) \). Lattice results allow to look for the solution in the form:

\[
D(z) = D(0) e^{-\frac{|z|}{T_g}} \sum_{n=0}^{\infty} \frac{d_n}{n!} \left( \frac{z}{T_g} \right)^n
\]

\[
D_2(x, z - x) = D_2(0) e^{-\frac{|z|}{T_g}} \sum_{m=0}^{\infty} \frac{b_m}{m!} \left( \frac{x}{T_g} \right)^m \cdot e^{-\frac{|z-x|}{T_g}} \sum_{m=0}^{\infty} \frac{b_m}{m!} \left( \frac{z-x}{T_g} \right)^m
\]

The substitution of (21) in Eq. (19) gives a relation between \( d_n \) and \( b_m \):

\[
d_{n+2} = d_{n+1} - \frac{(a T_g)^2}{2} \sum_{m=0}^{n} b_m b_{n-m}
\]

with the obvious conditions \( d_0 = b_0 = d_1 = 1 \). There are infinitely many sets of coefficients \( b_m, d_m \), which solve the above equation. What solution is realized in the real world is a question for future investigations.

3

It remains to be seen whether the above picture is compatible with the Gaussian model for QCD vacuum [4], [5]. In the latter all correlators of the odd
orders, including $\Delta^{(3)}$ are assumed to be zero, while the even correlators are factorized in a product of the 2-point ones. Among main arguments in favour of it is the fact, that potential of confinement for higher representations of a gauge group is proportional to the quadratic Casimir operator in the Gaussian vacuum. This fact is confirmed with a good accuracy by lattice computations (see discussion in [9]). The Gaussian model for QCD vacuum was introduced in [1] and intensively used in [7] for the high energy scattering amplitudes. Eq. (19) shows clearly, that such picture is not selfconsistent for any nontrivial function $D(z)$. Thus one should introduce an extended Gaussian ensemble and postulate factorization in a product of the 2-point and 3-point correlators. If 3-point correlator is zero, this definition coincides with the ordinary Gaussian ensemble, because all odd order correlators are also zero, in the opposite case a new stochastic ensemble arises, it is natural to call it the minimally extended Gaussian ensemble. At the same time even in the ordinary Gaussian vacuum function $D(z)$ may be a nontrivial one. We shall deal in this case with the second derivative of the 2-point correlator at $z = 0$:

$$\frac{\partial^2 \Delta^{(2)}}{\partial z_2 \partial z_1} \varepsilon_{\mu_1 \nu_1 \rho \eta} \varepsilon_{\mu_2 \nu_2 \xi \gamma} \bigg|_{z_1 \rightarrow z_2} = \varepsilon_{\mu_1 \nu_1 \rho \eta} \varepsilon_{\mu_2 \nu_2 \xi \gamma} \times$$

$$\times \left[ -\frac{i}{2} < \text{Tr}(F_{\mu_1 \nu_1} [F_{\mu_2 \nu_2}] > + \frac{i z_\sigma}{6} < \text{Tr}(F_{\mu_1 \nu_1} [D_\rho F_{\sigma \xi} F_{\nu_2 \nu_2}] > -

-\frac{z_\sigma z_\phi}{24} \left( 6 < \text{Tr}(F_{\mu_1 \nu_1} F_{\sigma \xi} [F_{\nu_2 \nu_2}] > + 6 < \text{Tr}(F_{\mu_1 \nu_1} [F_{\nu_2 \nu_2} F_{\sigma \xi}] F_{\phi \rho}) > +

+ < \text{Tr}(F_{\mu_1 \nu_1} [F_{\phi \rho} F_{\nu_2 \nu_2}] > + < \text{Tr}(F_{\mu_1 \nu_1} F_{\nu_2 \nu_2} [F_{\phi \rho} F_{\sigma \xi}] ) > \right) \right] \right] \quad (23)$$

Taking into account the following:

$$\frac{\partial^2 \Delta^{(2)}}{\partial z_2 \partial z_1} \varepsilon_{\mu_1 \nu_1 \rho \eta} \varepsilon_{\mu_2 \nu_2 \xi \gamma} = 16 \frac{d^2 D(z)}{(dz)^2} \left( z_\gamma z_\eta - \delta_\gamma \delta_\eta z^2 \right) - 24 \frac{dD}{dz} \delta_\gamma \delta_\eta \quad (24)$$

one has:

$$-48 D''|_0 = \lim_{z \rightarrow 0} \left[ \frac{i z_\sigma}{6 z^2} \varepsilon_{\mu_1 \nu_1 \rho \eta} \varepsilon_{\mu_2 \nu_2 \xi \eta} < \text{Tr}(F_{\mu_1 \nu_1} [D_\rho F_{\sigma \xi}, F_{\nu_2 \nu_2}] > -

-\frac{z_\sigma z_\phi}{24 z^2} \left( 6 < \text{Tr}(F_{\mu_1 \nu_1} F_{\phi \rho} [F_{\sigma \xi} F_{\nu_2 \nu_2}] ) > + 6 < \text{Tr}(F_{\mu_1 \nu_1} [F_{\nu_2 \nu_2} F_{\sigma \xi}] F_{\phi \rho}) > +$$

8
\[ + \langle Tr(F_{\mu_1\nu_1}[F_{\sigma\xi}F_{\phi\rho}]F_{\mu_2\nu_2}) \rangle + \langle Tr(F_{\mu_1\nu_1}F_{\mu_2\nu_2}[F_{\phi\rho}F_{\sigma\xi}]) \rangle \]  

(25)

The first and the second terms in the r.h.s. of (23) are absent in the Gaussian vacuum, 4–point condensate has to be introduced for the analysis of the others:

\[ \Delta^{(4)} = \langle Tr(F_{\mu_1\nu_1}\Phi F_{\mu_2\nu_2}\Phi F_{\mu_3\nu_3}\Phi F_{\mu_4\nu_4}\Phi) \rangle = D^{\text{connect}}_4 + \]

\[ + (\alpha_0 \cdot \varepsilon_{\mu_1\nu_1ab} \varepsilon_{\mu_2\nu_2bc} \varepsilon_{\mu_3\nu_3cd} \varepsilon_{\mu_4\nu_4de} + \alpha_1 \cdot \varepsilon_{\mu_1\nu_1ab} \varepsilon_{\mu_3\nu_3ba} \varepsilon_{\mu_2\nu_2cd} \varepsilon_{\mu_4\nu_4de} + \]

\[ + \alpha_2 \cdot \varepsilon_{\mu_1\mu_2ab} \varepsilon_{\mu_4\nu_4ba} \varepsilon_{\mu_2\nu_2cd} \varepsilon_{\mu_3\nu_3de} \rangle \cdot D^{\text{Gauss}}_4 + \text{nonkron. terms} \]  

(26)

here the connected 4–point function \( D^{\text{connect}}_4 \) should also be zero in the Gaussian vacuum.

Using (5):

\[ \langle F^a_{\mu_1\nu_1} F^b_{\mu_2\nu_2} \rangle = \delta^{ab} \frac{8}{8} \varepsilon_{\mu_1\nu_1\rho\sigma} \varepsilon_{\mu_2\nu_2\rho\sigma} (D(0) + D_1(0)) \]  

(27)

it is easy to get:

\[ \alpha_0 = \frac{16}{3}; \alpha_1 = \frac{2}{3}; D^{\text{Gauss}}_4 = \left( \frac{D(0) + D_1(0)}{8} \right)^2 \]  

(28)

Straightforward calculations lead to the result:

\[ \frac{d^2 D(z)}{(dz)^2} \bigg|_{z=0} = \frac{7}{2} \left( \frac{D(0) + D_1(0)}{4} \right)^2 \]  

(29)

We see, that even for the Gaussian ensemble the presence of the nonperturbative condensate

\[ D(0) + D_1(0) \neq 0 \]  

(30)

automatically provides the dynamical dependence on \( z \) for function \( D(z) \).

4

The relation (29) demonstrates the existence of a scale, which determines the behaviour of \( \Delta^{(2)}(z) \) at small distances and it may be quite different from
the asymptotic correlation length \( T_g \). Indeed, taking from \[4\] \( D_1(z) \approx \frac{1}{3} D(z) \) and defining \( D(0) + D_1(0) \) from the standard gluon condensate \[8\], we get

\[
\frac{D''(0)}{D(0)} \approx 0.012 \text{ Gev}^4
\]

whence it follows that (parametrizing function \( D(z) \) near zero as a gaussian)

\[
D(z) = D(0)e^{-z^2/r_0^2}
\]

and \( r_0 \approx 0.6 \text{ fm} \). This value of \( r_0 \) is larger than \( T_g \approx 0.2 \text{ fm} \), characterizing the exponential asymptotic decreasing of the correlator for \( z > 0.2 \text{ fm} \):

\[
D(z) = D(0)e^{-|z|/T_g}
\]

if \( z > 0.2 \).

This situation is possible in the so called string limit:

\[
T_g^2(D(0) + D_1(0)) \sim \sigma = \text{const}
\]

with \( T_g \to 0 \). It is seen from (29), that also:

\[
\frac{T_g^4}{D(0) + D_1(0)} \left. \frac{d^2D(z)}{(dz^2)^2} \right|_{z=0} \to 0 \quad (31)
\]

Substituting in (31) values for \( D(0) + D_1(0) \) and \( T_g \) from \[4\], we get on the l.h.s. \( 0.04 \sim \left( \frac{T_g}{r_0} \right)^4 \), that makes string limit in the nonperturbative QCD credible.

Let us make another important notice about the Gaussian approximation for the gluodynamic vacuum. The Taylore expansion for even function \( D(z) \) is as follows:

\[
D(z) = \sum_{n=0}^{\infty} \frac{d^nD(z)}{(dz^2)^n} \left|_{z=0} \right. \frac{z^{2n}}{n!}
\]

The coefficients of the above expansion are expressed through the \((2 + n)\)-th and less orders condensates. We have seen that for \( n = 1 \) and \( n = 2 \) the derivatives at \( z = 0 \) are zero and positive respectively for the Gaussian ensemble. It may be interesting to check, are there negative coefficients in this Taylor expansion for the Gaussian vacuum or not. The latter would be
obviously incompatible with the observed decreasing of the function $D(z)$ with $z$ and would be a rigorous proof of nongaussian nature of gluodynamic vacuum.

The behaviour of the function $D(z)$ for $z \lesssim 0.2\text{fm}$ is unknown at the moment. Its determination and higher correlators measurements would take a chance to check the relations (13), (14), (22) and make a conclusion, what type of vacuum is realized in gluodynamics and whether one can approximate this vacuum as the Gaussian or the minimally extended Gaussian ensemble of the stochastic fields. We have said nothing about perturbative parts of correlators and their connection with the nonperturbative ones. This important question will be considered in its own right.

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