QCD: Confinement, Hadron Structure and DIS

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To the memory of my first teacher
Isaac Yakovlevich Pomeranchuk.

Abstract
The main features of QCD, e.g. confinement, chiral symmetry breaking, Regge trajectories are naturally and economically explained in the framework of the Field Correlator Method (FCM). The same method correctly predicts the spectrum of hybrids and glueballs. When applied to DIS and high-energy scattering it leads to the important role of higher Fock components in the Fock tower of moving hadron, containing primarily gluonic excitations.

1 Introduction
The QCD is the (only) internally selfconsistent theory, defined by the only scale (parameter the string tension $\sigma$ or $\Lambda_{QCD}$ can be used for this purpose), which eventually explains fundamental structure of baryons and other hadrons and by this some 98% of all the mass in our world. The features of QCD are unique in their complexity: confinement, chiral symmetry breaking, string structure of hadrons demonstrating nonperturbative (NP) interactions, on one hand and on another hand the applicability of perturbation

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theory for large $Q$ and large momentum processes, where NP effects enter as corrections. The most popular theoretical approaches to these phenomena look fragmentary, i.e. magnetic monopoles are used for confinement, instantons for chiral symmetry breaking and pure perturbative expansions for high energy processes.

The situation has improved with the introduction of the QCD sum rules \cite{1}, where NP contributions are encoded in the form of local condensates. In this talk I shall describe a general method which allows to consider all NP effects on one ground – with the help of nonlocal Field Correlators (FC) \cite{2}. It will be argued that the simplest of these correlators – the quadratic one – is dominant and is responsible for confinement, chiral symmetry breaking, and the structure of meson and baryon spectra.

When supplemented with the Background Perturbation Theory (BPT) the method allows to treat valence gluons in the confining background \cite{3}. This yields the spectrum of hybrids and glueballs in good agreement with lattice simulations, and the new perturbation series for high energy processes, without IR renormalons and Landau ghost poles. As the new and unexpected element, it will be argued that hybrids play a more fundamental role in DIS and high-energy scattering – as the building blocks of the colliding hadron wave functions.

The talk is organized as follows. In section 2 FC are introduced and confinement is related to their properties. In section 3 the Hamiltonian for valence components is written down and properties of meson spectrum are discussed. In section 4 this Hamiltonian is extended to the systems with valence gluons and spectrum of hybrids and glueballs is discussed. In section 5 the Fock towers are introduced for hadrons and the general matrix Hamiltonians is written down. In section 6 the role of higher (hybrid) Fock components in DIS is discussed and the gluon contribution to structure functions and proton spin is emphasized. The last section concludes the general picture of the QCD dynamics with the discussion of the selfconsistent calculation of the FC.
2 The QCD vacuum structure. Stochastic vs coherent.

The basic quantity which defines the vacuum structure in QCD is the field correlator (FC)

\[ D^{(n)}(x_1, \ldots x_n) \equiv \langle F_{\mu_1 \nu_1}(x_1) F_{\mu_2 \nu_2}(x_2) \Phi(x_2, x_3) \ldots F_{\mu_n \nu_n}(x_n) \Phi(x_n, x_1) \rangle, \]

\[ (1) \]

\[ \Phi(x, y) = P \exp ig \int_y^x A_\mu(z) dz. \]

The set of FC (1) for \( n = 2, 3, \ldots \) gives a detailed characteristic of vacuum structure, including field condensates (for coinciding \( x_1 = x_2 = \ldots x_n \)). On general grounds one can distinguish two opposite situations: 1) stochastic vacuum 2) coherent vacuum. In the first case FC form a hierarchy with the dominant lowest term \( D^{(2)}(x_1, x_2) = D^{(2)}(x_1 - x_2) \), while higher FC are fast decreasing with \( n \). We shall call this situation the Gaussian Stochastic Approximation (GSA). In the second case all FC are comparable, and expansion of physical amplitudes as the series of FC is impractical. This is the case for the gas/liquid of classical solutions, e.g. of instantons, magnetic monopoles etc. The physical picture behind the situation of nonconverging FC series is that of the coherent lump(s), when all points in the lump are strongly correlated.

To understand where belongs the QCD vacuum one can start with the Wilson loop in the representation \( D \) of the color group SU(3),

\[ W_D(C) = \langle tr_D \exp(ig \int_C d_\mu A_\mu^a T_a^{(D)}) \rangle \]

\[ (2) \]

The Stokes theorem and the cluster expansion identity allow to obtain the basic equation, which is used in most applications of the FCM (for more details see [2])

\[ W(C) = \exp \sum_n \frac{(ig)^n}{n!} \int D^{(n)}(x_1, \ldots x_n) d_\sigma_{\mu_1 \nu_1}(x_1) \ldots d_\sigma_{\mu_n \nu_n}(x_n). \]

\[ (3) \]

Here integration is performed over the minimal surface \( S_{\text{min}} \) inside the contour \( C \) defined in (2) and \( D_n \) is the so-called cumulant or the connected correlator, obtained from the FC Eq. (1) by subtracting all disconnected averages. From (3) one easily obtains that the Wilson loop has the area-law asymptotics, \( W(C) \sim \exp(-\sigma S_{\text{min}}) \), for any finite number of terms \( n_{\text{max}} \); \( n \leq n_{\text{max}} \) in the exponent (3).
The string tension is expressed through $\tilde{D}^{(n)}$.

$$\sigma = \frac{1}{2} \int D^{(2)}(x_1 - x_2) d^2(x_1 - x_2) + 0(\tilde{D}^{(n)}, n \geq 4) = \sigma_2 + \sigma_4 + \ldots \quad (4)$$

Eq. (4) has several consequences: 1) confinement appears naturally for $n = 2$, i.e. in the GSA 2) the lack of confinement can be due to vanishing of all FC, or due to the special cancellation between the cumulants, as it happens for the instanton vacuum $[4, 3]$ for static quarks in the representation $D$ of the color group SU(3), the string tension $\sigma_2^{(D)}$ is proportional to the quadratic Casimir factor (the Casimir scaling)

$$\sigma_2^{(D)} = \frac{C_2^{(D)}}{C_2^{(fund)}} \sigma_2^{(fund)}, C_2^{(D)} = \frac{1}{3}(\mu^2 + \mu\nu + \nu^2 + 3\mu + 3\nu) \quad (5)$$

However for larger $n, n \geq 4$ the Casimir scaling is violated:

$$\sigma_n^{(D)} = a_1 C_2^{(D)} + a_2 (C_2^{(D)})^2 + a_3 C_3^{(D)} + \ldots \quad (6)$$

It is remarkable that perturbative interaction of static quarks $V^{(D)}(r)$ satisfies the Casimir scaling to the order $O(g^6)$ considered so far $[5]$, so the total potential $V^{(D)}(r) = V_{pert}^{(D)}(r) + \sigma^{(D)} r + const$ is also Casimir scaling, if GSA works well.

This picture was tested recently on the lattice $[6]$ and confirmed the Casimir scaling with the accuracy around 1% in the range $0.1 \leq r \leq 1.1$ fm. The full theoretical understanding of this fundamental fact is still lacking, both for the perturbative part and for the string tension. On the pedestrian level the Casimir scaling and the quadratic (Gaussian) correlator dominance implies that the vacuum is highly stochastic and any quasiclassical objects, like instantons, are strongly suppressed in the real QCD vacuum. The vacuum consists of small white dipoles of the size $T_g$ made of neighboring field strength operators. The smallness of $T_g$ might be an explanation for the Gaussian dominance since higher correlator terms in $\sigma$ are proportional to $(FT_g^2)^n (T_g^2)^{-1}$, where $F$ is the estimate of the average nonperturbative vacuum field, $F \sim 500$ (MeV)$^2$.

Lattice calculations of FC have been done repeatedly during last decades, using cooling technic $[7]$ and with less accuracy without cooling $[8]$. (Recently

\footnote{The diagrams of the order $O(g^8)$ violating the Casimir scaling have been found by W.Wetzel (to be published).}
another approach based on the so-called gluelump states was exploited on the lattice [9] and analytically [10], which has a direct connection to FC.

The basic result of [7] is that FC consists of perturbative part \( O(1/x^4) \) at small distances and nonperturbative part \( O(\exp(-x/T_g)) \) at larger distances with \( T_g \) in the range \( T_g = 0.2 \) fm (quenched vacuum) and \( T_g = 0.3 \) fm (2 flavours).

Calculations in [8] and [9, 10], as well as sum rule estimates [11] yield a smaller value, \( T_g \approx 0.13 \) fm to 0.17 fm. This enables us in what follows to take the limit \( T_g \to 0 \) keeping \( \sigma = \text{const} \approx 0.18 \text{ GeV}^2 \), and consider \( \sigma T_g^2 \) as a small parameter of expansion, \( \sigma T_g^2 \ll 1 \). For example, the contribution of higher correlators in [4] is proportional to \( \sigma(\sigma T_g^2)^{n-1}, n = 4, 6, 8, ... \)

3 Hamiltonian for valence components

There are two possible approaches to incorporating nonperturbative field correlators in the quark-antiquark (or 3q) dynamics. The first has to deal with the effective nonlocal quark Lagrangian containing field correlators [13]. From this one obtains first-order Dirac-type integro-differential equations for heavy-light mesons [12, 14], light mesons and baryons [15]. These equations contain the effect of chiral symmetry breaking [12, 13] which is directly connected to confinement.

The second approach is based on the effective Hamiltonian for any gauge-invariant quark-gluon system. In the limit \( T_g \to 0 \) this Hamiltonian is simple and local, and in most cases when spin interaction can be considered as a perturbation one obtains results for the spectra in an analytic form, which is transparent.

For this reason we choose below the second, Hamiltonian approach [16, 17]. We start with the exact Fock-Feynman-Schwinger Representation for the \( qq \) Green’s function (for a review see [18]), taking for simplicity nonzero flavor case

\[
G^{(x,y)}_{qq} = \int_0^\infty ds_1 \int_0^\infty ds_2 (Dz)_{xy}(D\bar{z})_{xy} e^{-K_1 - K_2}
\]

\[
\langle tr\Gamma_{in}(m_1 - \hat{D}_1)W_\sigma(C)\Gamma_{out}(m_2 - \hat{D}_2)\rangle_A
\]

(7)

where \( K_i = \int_0^{\delta^i} d\tau_i (m_i + \frac{1}{4}(\dot{z}_{\mu}^{(i)})^2) \), \( \Gamma_{in,out} = 1, \gamma_5, ... \) are meson vertices, and \( W_\sigma(C) \) is the Wilson loop with spin insertions, taken along the contour \( C \).
formed by paths \((Dz)_{xy}\) and \((D\bar{z})_{xy}\),

\[
W_\sigma(C) = P_F P_A \exp (ig \int_C A_\mu dz_\mu) \times
\]
\[
\times \exp \left( g \int_0^{s_1} \sigma_\mu^{(1)} F_{\mu\nu} d\tau_1 - g \int_0^{s_2} \sigma_{\mu\nu}^{(2)} F_{\mu\nu} d\tau_2 \right).
\]

The last factor in (8) defines the spin interaction of quark and antiquark. The average \(\langle W_\sigma(C) \rangle_A\) in (7) can be computed exactly through field correlators \(\langle F^{(1)}...F^{(n)} \rangle_A\), and keeping only the lowest one, \(\langle F^{(1)}F^{(2)} \rangle\), which yields according to lattice calculation [6] accuracy around 1% [5], one obtains

\[
\langle W_\sigma(C) \rangle_A \simeq \exp (-\frac{1}{2} \int_{S_{min}} ds_{\mu\nu} \int_{S_{min}} ds_{\lambda\sigma}(1) + \sum_{i,j=1}^2 \int_0^{s_i} \sigma_\mu^{(i)} d\tau_i \int_0^{s_j} \sigma_{\lambda\sigma}^{(j)} d\tau_j \langle F_{\mu\nu}(1)F_{\lambda\sigma}(2) \rangle).
\]

The Gaussian correlator \(\langle F_{\mu\nu}(1)F_{\lambda\sigma}(2) \rangle \equiv D_{\mu\nu,\lambda\sigma}(1,2)\) can be rewritten identically in terms of two scalar functions \(D(x)\) and \(D_1(x)\) [2], which have been computed on the lattice [7] to have the exponential form \(D \sim \exp(-|x|/T_g)\) with the gluon correlation length \(T_g \approx 0.2\) fm.

As the next step one introduces the einbein variables \(\mu_i\) and \(\nu\); the first one to transform the proper times \(s_i, \tau_i\) into the actual (Euclidean) times \(t_i \equiv z_4^{(i)}\). One has [17]

\[
2\mu_i(t_i) = \frac{dt_i}{d\tau_i}, \quad \int_0^\infty ds_i (D^4z^{(i)})_{xy} = \text{const} \int D\mu_i(t_i)(D^4z^{(i)})_{xy}.
\]

The variable \(\nu\) enters in the Gaussian representation of the Nambu-Goto form for \(S_{min}\) and its stationary value \(\nu_0\) has the physical meaning of the energy density along the string. In case of several strings, as in the baryon case or the hybrid case, each piece of string has its own parameter \(\nu^{(i)}\).

To get rid of the path integration in (7) one can go over to the effective Hamiltonian using the identity

\[
G_{q\bar{q}}(x,y) = \langle x | \exp(-HT) | y \rangle
\]

where \(T\) is the evolution parameter corresponding to the hypersurface chosen for the Hamiltonian: it is the hyperplane \(z_4 = \text{const}\) in the c.m. case [17].
The final form of the c.m. Hamiltonian (apart from the spin and perturbative terms to be discussed later) for the $q\bar{q}$ case is [17] [19]

$$H_0 = \sum_{i=1}^{2} \left( \frac{m_i^2 + p_i^2}{2\mu_i} + \frac{\mu_i}{2} \right) + \frac{\hat{L}^2/r^2}{2} + \frac{\sigma^2 r^2}{2} \int_0^1 \frac{d\beta}{\nu(\beta)} + \int_0^1 \frac{\nu(\beta)}{2} d\beta. \quad (12)$$

Here $\zeta = (\mu_1 + \int_0^1 \beta \nu d\beta)/(\mu_1 + \mu_2 + \int_0^1 \beta \nu d\beta)$ and $\mu_i$ and $\nu(\beta)$ are to be found from the stationary point of the Hamiltonian

$$\frac{\partial H_0}{\partial \mu_i} \bigg|_{\mu_i = \mu_i^{(0)}} = 0, \quad \frac{\partial H_0}{\partial \nu} \bigg|_{\nu = \nu^{(0)}} = 0. \quad (13)$$

Note that $H_0$ contains as input only $m_1, m_2$ and $\sigma$, where $m_i$ are current masses defined at the scale of 1 GeV. The further analysis is simplified by the observation that for $L = 0$ one finds $\nu^{(0)} = \sigma r$ from (13) and $\mu_i = \sqrt{m_i^2 + p_i^2}$, hence $H_0$ becomes the usual Relativistic Quark Model (RQM) Hamiltonian

$$H_0(L = 0) = \sum_{i=1}^{2} \sqrt{m_i^2 + p_i^2} + \sigma r. \quad (14)$$

But $H_0$ is not the whole story, one should take into account 3 additional terms: spin terms in [20] which produce two types of contributions: self-energy correction $H_{self} = \sum_{i=1}^{2} \frac{\Delta m_i^2(i)}{2\mu_i}$, $\Delta m^2_q = -\frac{4\sigma}{\pi}\eta(m_i)$, $\eta(0) = \frac{3}{4} \left(1 + \frac{D_1(0)}{D(0)}\right) \approx 1 + 0.9$

and spin-dependent interaction between quark and antiquark $H_{spin}$ [21] which is entirely described by the field correlators $D(x), D_1(x)$, including also the one-gluon exchange part present in $D_1(x)$.

Finally one should take into account gluon exchange contributions, which can be divided into the Coulomb part $H_{Coul} = -\frac{4\alpha_s(x)}{3} \frac{\mu_i}{\alpha_s(x)}$, and $H_{rad}$ including space-like gluon exchanges and perturbative self-energy corrections (we shall systematically omit these corrections since they are small for light quarks to be discussed below). In addition there are gluon contributions which are nondiagonal in number of gluons $n_g$ and quarks (till now only the sector $n_g = 0$ was considered) and therefore mixing meson states with hybrids and
glueballs [22]; we call these terms \( H_{mix} \) and refer the reader to [22] and the cited there references for more discussion. Assembling all terms together one has the following total Hamiltonian in the limit of large \( N_c \) and small \( T_g \) (for more discussion see [23]):

\[
H = H_0 + H_{self} + H_{spin} + H_{Coul} + H_{rad} + H_{mix}.
\] (16)

We start with \( H_0 = H_R + H_{string} \). The eigenvalues \( M_0 \) of \( H_R \) can be given with 1% accuracy by [24]

\[
M_0^2 \approx 8\sigma L + 4\pi\sigma(n + \frac{3}{4})
\] (17)

where \( n \) is the radial quantum number, \( n = 0, 1, 2, ... \). Remarkably \( M_0 \approx 4\mu_0 \), and for \( L = n = 0 \) one has \( \mu_0(0, 0) = 0.35 \) GeV for \( \sigma = 0.18 \) GeV\(^2\), and \( \mu_0 \) is fast increasing with growing \( n \) and \( L \). This fact partly explains that spin interactions become unimportant beyond \( L = 0, 1, 2 \) since they are proportional to \( d\tau_1 d\tau_2 \sim \frac{1}{4\mu_1\mu_2} dt_1 dt_2 \) (see [10] and [23]). Thus constituent mass (which is actually "constituent energy") \( \mu_0 \) is "running". The validity of \( \mu_0 \) as a socially accepted "constituent mass" is confirmed by its numerical value given above, the spin splittings of light and heavy mesons [25] and by baryon magnetic moments expressed directly through \( \mu_0 \), and being in agreement with experimental values [26].

4 Hamiltonian and bound states of valence gluons

We now come to the gluon-containing systems, hybrids and glueballs. Referring the reader to the original papers [24]-[29] one can recapitulate the main results for the spectrum. In both cases the total Hamiltonian has the same form as in (16), however the contribution of corrections differs.

For glueballs it was argued in [29] that \( H_0 \) (12) has the same form, but with \( m_i = 0 \) and \( \sigma \rightarrow \sigma_{adj} = \frac{2}{3}\sigma \) while \( H_{self} = 0 \) due to gauge invariance.

Thus one can retain in (16) only two main terms: \( H = H_0 + H_{spin} \) while \( H_{Coul} \) was argued to be strongly decreased by loop corrections. The calculation in [29] was done for two-gluon and three-gluon glueball states and results are in surprisingly good agreement with lattice data for both systems (no fitting parameters have been used in [29]).
We now coming to the next topic of this talk: hybrids and their role in hadron dynamics. We start with the hybrid Hamiltonian and spectrum. This topic in the framework of FCM was considered in \cite{27, 28} The Hamiltonian \( H_0 \) for hybrid looks like \cite{23, 27, 28}

\[
H_{0}^{(hyb)} = \frac{m_1^2}{2\mu_1} + \frac{m_2^2}{2\mu_2} + \frac{\mu_1 + \mu_2 + \mu_g}{2} + \frac{p_\xi^2 + p_\eta^2}{2\mu} + \sigma \sum_{i=1}^{2} |r_g - r_i| + H_{str}. \tag{18}
\]

Here \( p_\xi, p_\eta \) are Jacobi momenta of the 3-body system, \( H_{self} \) is the same as for meson, while \( H_{spin} \) and \( H_{Coul} \) have different structure \cite{28}.

The main feature of the present approach based on the BPTh, is that valence gluon in the hybrid is situated at some arbitrary point on the string connecting quark and antiquark, and the gluon creates a kink on the string so that two pieces of the string move independently (however connected at the point of gluon). This differs strongly from the flux-tube model where hybrid is associated with the string excitation as a whole, but has a strong similarity to the treatment of gluons in the framework of the Lund model \cite{30}.

Results for light and heavy exotic \( 1^{-+} \) hybrids are also given in \cite{23} and are in agreement with lattice calculations. Typically an additional gluon in the exotic \( (L = 1) \) state "weights" 1.2÷1.5 GeV for light to heavy quarks, while nonexotic gluon \( (L = 0) \) brings about 1 GeV to the mass of the total \( q\bar{q}g \) system. Let us now consider the hybrid spectrum in more detail. First of all we use for that 3-body problem the hyperspherical method, which works with accuracy of few percent \cite{31, 32}. Then the whole spectrum is classified by the grand angular momentum \( K = 0, 1, 2, ... \), which is actually an arithmetic sum of all partial pair angular momenta in the system \( q\bar{q}g \). The lowest \( K = 0 \) states can be formed from the s-wave \( q\bar{q} \) pair and s-wave valence gluon \( g \), which gives the \( \rho + g \) and \( \pi + g \) systems, and vectors imply spin-one particle. In this way one obtains the classification

\[
K = 0, (\pi + g) - 1^{+-}, (\rho + g) - (2^{++}, 1^{++}, 0^{++})
\]

\[
K = 1, (\pi + (\nabla \times g)) - 1^{--}, (\rho + (\nabla \times g)) - (2^{-+}, 1^{-+}, 0^{-+}).
\]

The eigenvalues of the Hamiltonian \( (18) \) are easily obtained for the light quarks using the hypercentral (lowest \( K \)) approximation –

\[
M(K = 0) = 1.872 \text{ GeV}
\]
\[ M(K = 1) = 2.45 \text{ GeV} \]
\[ M(K = 2) = 2.90 \text{ GeV} \]
\[ M(K = 3) = 3.27 \text{ GeV}. \]

Here the (negative) self-energy part of quarks \( H_{\text{self}} \) is already added to masses. One has also in addition the color Coulomb part \( H_{\text{Coul}} \) and spin-dependent part \( H_{\text{spin}} \), which contribute approximately \( \langle H_{\text{Coul}} \rangle \sim -0.2 \text{ GeV} \) and for spin-spin interaction approximately 0.08 GeV \( \begin{pmatrix} -2 \\ -1 \\ +1 \end{pmatrix} \), where numbers inside brackets refer to \( \vec{1} + \vec{1} = \vec{0}, \vec{1}, \vec{2} \) respectively from top to bottom.

As a result, neglecting rather small contribution from \( H_{\text{rad}} \) and \( H_{\text{mix}} \) in (16), and the string correction, taking into account the moment of inertia of the string [17], yielding around (-100 MeV), one has the approximate mass estimates for the lowest mass states of \( K \) multiplets:

\[ M_{\text{low}}(K = 0) \approx 1.42 \text{ GeV} \]
\[ M_{\text{low}}(K = 1) \approx 1.9 \text{ GeV} \]
\[ M_{\text{low}}(K = 2) \approx 2.45 \text{ GeV}. \]

The mass value \( M_{\text{low}}(K = 1) \) agrees well with the lattice results for the mass of the \( 1^{-}+ \) state [33]. The recent unquenched calculation [34] yields the value which is somewhat lower.

One should stress that the hybrid states, which start at the mass around 1.4 GeV, have a high multiplicity which grows exponentially with mass, as well as excited string states in bossonic string theory [35]. This fact has a very important consequence for high-energy processes, where the hybrid excitation is argued to be the dominant physical mechanism.

5 Hamiltonian and Fock states

As was mentioned above the QCD Hamiltonian is introduced in correspondence with the chosen hypersurface, which defines internal coordinates \( \{ \xi_k \} \) lying inside the hypersurface, and the evolution parameter, perpendicular to it. Two extreme choices are frequently used, 1) the c.m. coordinate system with the hypersurface \( x_4 = \text{const.} \), which implies that all hadron constituents
have the same (Euclidean) time coordinates $x_4^{(i)} = \text{const}, i = 1, \ldots, n$, 2) the light-cone coordinate system, where the role of $x_4$ and $x_4^{(i)}$ is played by the $x_+^+, x_+^{(i)}$ components, $x_+ = \frac{x_0 + x_3}{\sqrt{2}}$.

To describe the structure of the Hamiltonian in general terms we first assume that the bound valence states exist for mesons, glueballs and baryons consisting of minimal number of constituents. To form the Fock tower of states starting with the given valence state, one can add gluons and $q\bar{q}$ pairs keeping the $J^{PC}$ assignment intact. At this point we make the basic simplifying approximation assuming that the number of colors $N_c$ is tending to infinity, so that one can do for any physical quantity an expansion in powers of $1/N_c$. Recent lattice data confirm a good convergence of this expansion for $N_c = 3, 4, 6$ and all quantities considered \[36\] (glueball mass, critical temperature, topological susceptibility etc.).

Then the construction of the Fock tower is greatly simplified since any additional $q\bar{q}$ pair enters with the coefficient $1/N_c$ and any additional white (e.g. glueball) component brings in the coefficient $1/N_c^2$. In view of this in the leading order of $1/N_c$ the Fock tower is formed by only creating additional gluons in the system, i.e. by the hybrid excitation of the original (valence) system. Thus all Fock tower consists of the valence component and its hybrid equivalents and each line of this tower is characterized by the number $n$ of added gluons. Then, the internal coordinates $\{\xi\}_n$ describe coordinates and polarizations of $n$ gluons in addition to those of valence constituents.

We turn now to the Hamiltonian $H$, assuming it to be either the total QCD Hamiltonian $H_{\text{QCD}}$, or the effective Hamiltonian $H^{(\text{eff})}$, obtained from $H_{\text{QCD}}$ by integrating out short-range degrees of freedom. We shall denote the diagonal elements of $H$, describing the dynamics of the $n$-th hybrid excitation of $s$-th valence state ($s = m\{f\bar{f}\}, gg, 3g, b\{f_1f_2f_3\}$ for mesons, 2-gluon and 3-gluon glueballs and baryons respectively with $f_i$ denoting flavour of quarks) as $H^{(s)}_{nn}$. For nondiagonal elements we confine ourselves to the lowest order operators $H^{(s)}_{n,n+1}$ and $H^{(s)}_{n-1,n}$ describing creation or annihilation of one additional gluon, viz.

$$H_{qq} = g \int \bar{q}(x,0)\hat{a}(x,0)q(x,0)d^3x$$

$$H_{gg} = \frac{g}{2}f^{abc} \int (\partial_\mu a^a_\nu - \partial_\nu a^a_\mu)a^b_\mu a^c_\nu d^3x,$$

and we disregard for simplicity the terms $H_{g3g}$. 

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As it is clear from (19), (20), the first operator refers to the gluon creation from the quark line, while the second refers to the creation of 2 gluons from the gluon line. In what follows we shall be mostly interested in the first operator, which yields dominant contribution at large energies, and physically describes addition of one last cross-piece to the ladder of gluon exchanges between quark lines, while (20) corresponds in the same ladder to the \( \alpha_s \) renormalization graphs, and to the graphs with creation of additional gluon line.

The effective Hamiltonian in the one-hadron sector can be written as follows

\[
\hat{H} = \hat{H}^{(0)} + \hat{V}
\]

where \( \hat{H}^{(0)} \) is the diagonal matrix of operators,

\[
\hat{H}^{(0)} = \{ H_0^{(s)}, H_1^{(s)}, H_2^{(s)}, \ldots \}
\]

while \( \hat{V} \) is the sum of operators (19) and (20), creating and annihilating one gluon. In (22) \( H_n^{(s)} \) is the Hamiltonian operator for what we call the ”n-hybrid”, i.e. a bound state of the system, consisting of \( n \) gluons together with the particles of the valence component. In this way the \( n \)-hybrid for the valence \( \rho \)-meson is the system consisting of \( q \bar{q} \) plus \( n \) gluons ”sitting” on the string connecting \( q \) and \( \bar{q} \).

Before applying the stationary perturbation theory in \( \hat{V} \) to the Hamiltonian (21), one should have in mind that there are two types of excitations of the ground state valence Fock component: 1) Each of the operators \( H_n^{(s)} \), \( n = 0, 1, \ldots \) has infinite amount of excited states, when radial or orbital motion of any degree of freedom is excited, 2) in addition one can add a gluon, which means exciting the string and this excitations due to the operator \( \hat{V} \) transforms the \( n \)-th Fock component \( \psi_n^{(s)} \) into \( \psi_{n+1}^{(s)} \).

The wave equation for the Fock tower \( \Psi_N \{ P, \xi \} \) has the standard form

\[
\hat{H} \Psi_N = (\hat{H}^{(0)} + \hat{V}) \Psi_N = E_N \Psi_N,
\]

or in the integral form

\[
\Psi_N = \Psi_N^{(0)} - G^{(0)} \hat{V} \Psi_N
\]

where \( G^{(0)} \) is diagonal in Fock components,

\[
G^{(0)}(E) = \frac{1}{\hat{H}^{(0)} - E}, \quad G_n^{(0)}(E) = \delta_{nn} \frac{1}{H_n^{(s)} - E},
\]

or in the integral form

\[
\Psi_N = \Psi_N^{(0)} - G^{(0)} \hat{V} \Psi_N
\]
and $\Psi_N^{(0)}$ is the eigenfunction of $\hat{H}^{(0)}$,

$$\hat{H}^{(0)}\Psi_N^{(0)} = E_N^{(0)}\Psi_N^{(0)}$$  \hspace{1cm} (26)$$

and since $\hat{H}^{(0)}$ is diagonal, $\Psi_N^{(0)}$ has only one Fock component, $\Psi_N^{(0)} = \psi_n(P, \{\xi\}_n)$, $n = 0, 1, 2, ...$, and the eigenvalues $E_N^{(0)}$ contain all possible excitation energies of the $n$-hybrid, with the number $n$ of gluons in the system fixed,

$$E_N^{(0)} = E_n^{(0)}(P) = \sqrt{P^2 + M_{n[k]}^2}.$$  \hspace{1cm} (27)$$

Here $\{k\}$ denotes the set of quantum numbers of the excited $n$-hybrid.

From (24) one obtains in the standard way corrections to the eigenvalues and eigenfunctions.

As a first step one should specify the unperturbed functions $\Psi_N^{(0)}$, introducing the set of quantum numbers $\{k\}$ defining the excited hybrid state for each $n$-hybrid Fock component $\psi_n(P, \{\xi\})$; we shall denote therefore:

$$\Psi_N^{(0)} = \psi_{n\{k\}}(P, \{\xi\}_n), \hspace{0.5cm} n = 0, 1, 2, ...$$  \hspace{1cm} (28)$$

The set of functions $\psi_{n\{k\}}$ with all possible $n$ and $\{k\}$ is a complete set to be used in the expansion of the exact wave-function (Fock tower) $\Psi_N$:

$$\Psi_N = \sum_{m\{k\}} c_{n\{k\}}^{m\{k\}} \psi_{m\{k\}}.$$  \hspace{1cm} (29)$$

Using the orthonormality condition

$$\int \psi_{n\{p\}}^{+} \psi_{n\{p\}} d\Gamma = \delta_{mn}\delta_{\{k\}\{p\}}$$  \hspace{1cm} (30)$$

where $d\Gamma$ implies integration over all internal coordinates and summing over all indices, one obtains from (24) an equation for $c_{n\{k\}}$ and $E_N$,

$$c_{n\{p\}}(E_N - E_n^{(0)}) = \sum_{m\{k\}} c_{m\{k\}}^{n\{p\}} V_{n\{p\},m\{k\}}$$  \hspace{1cm} (31)$$

where we have defined

$$V_{n\{p\},m\{k\}} = \int \psi_{n\{p\}}^{+} \hat{V} \psi_{m\{k\}} d\Gamma.$$  \hspace{1cm} (32)$$

Consider now the Fock tower built on the valence component $\psi_{\nu\{k\}}$, where $\nu$ can be any integer. For $\nu\{k\} = 0\{0\}$ this valence component corresponds
to the unperturbed hadron with minimal number of valence particles. For higher values of $\nu\{\kappa\}$ the Fock component $\psi_{\nu\{\kappa\}}$ corresponds to the hybrid with $\nu$ gluons which after taking into account the interaction is "dressed up" and acquires all other Fock components, so that the number $N$ in (29) contains the "bare number" $\nu\{\kappa\}$ as its part $N = \nu\{\kappa\}$, ... (at least for small perturbation $\hat{V}$).

One can impose on $\Psi_N$ the orthonormality conclusion

$$\int \Psi_N^+ \Psi_M d\Gamma = \sum_{m(k)} c_{m(k)}^N c_{m(k)}^M = \delta_{NM}. \quad (33)$$

Expanding now in powers of $\hat{V}$, one has

$$c_{m(k)}^{N(\nu\{\kappa\})} = \delta_{mn} \delta_{\{k\}\{\kappa\}} + c_{m(k)}^{N(1)} + c_{m(k)}^{N(2)} + ... \quad (34)$$

$$E_{N(\nu\{\kappa\})} = E_{\nu\{\kappa\}}^{(0)} + E_{N}^{(1)} + E_{N}^{(2)} + ... \quad (35)$$

It is easy to see that $E_{N\{1\}}^{(1)} \equiv 0$, while for $c^{(1)}$ one obtains from (31) the standard expression

$$c_{n\{p\}}^{N(1)} = \frac{V_{n\{p\},\nu\{\kappa\}}}{E_{\nu\{\kappa\}}^{(0)} - E_{n\{p\}}^{(0)}}. \quad (36)$$

In what follows we shall be interested in the high Fock components, $\nu+l, \{k\}$, obtained by adding $l$ gluons to the valence component $\nu\{\kappa\}$. Using (31) and (34) one obtains

$$c_{\nu+l\{k\}}^{N(\nu\{\kappa\})} = \sum_{\{k_1\}...\{k_l\}} \frac{V_{\nu+l\{k_1\},\nu+l-1\{k_1\}}}{E_{\nu\{\kappa\}}^{(0)} - E_{\nu+l\{k_1\}}^{(0)}} \frac{V_{\nu+l-1\{k_1\},\nu+l-2\{k_2\}}}{E_{\nu+l\{k_1\}}^{(0)} - E_{\nu+l-1\{k_1\}}^{(0)}} ... \quad (37)$$

Since $\hat{V}$ is proportional to $g$, one obtains in (34) the perturbation series in powers of $\alpha_s$ for $c^N$ and hence for $\Psi_N$ (29). One should note that $\alpha_s(Q^2)$ is the background coupling constant, having the property of saturation for positive $Q^2$ and the background perturbation series has no Landau ghost pole and is defined in all Euclidean region of $Q^2$.

The estimate of the mixing between meson and hybrid was done earlier in the framework of the potential model for the meson in [27]. In [24]
the mixing between hybrid, meson and glueball states was calculated in the framework of the present formalism and we shortly summarize the results. One must estimate the matrix element \( \langle \text{32} \rangle \) between meson and hybrid wave functions taking the operator \( \hat{V} \) in the form of \( \langle \text{19} \rangle \), where the operator of gluon emission at the point \((x, 0)\) can be approximated as

\[
\mathbf{a}_\mu(x, t) = \sum_{k, \lambda} \frac{1}{\sqrt{2\mu(k)V}} \times
\left[ \exp(ik \cdot x - i\mu t)\epsilon^{(\lambda)}_\mu(k) + \epsilon^{(\lambda)}_\mu(k) \exp(-ik \cdot x + i\mu t) \right]
\]

(38)

Omitting for simplicity all polarization vectors and spin-coupling coefficients which are of the order of unity, one has the matrix element

\[
V_{Mh} = \frac{g}{\sqrt{2\mu_g}} \int \varphi_M(r)^\mu \psi^+(h, 0, r)d^3r
\]

(39)

where \( \varphi_M(r) \), \( \psi^+(r_1, r_2) \) are meson and hybrid wave functions respectively, and in \( \langle \text{39} \rangle \) it is taken into account that the gluon is emitted (absorbed) from the quark position.

Using realistic Gaussian approximation for the wave functions in \( \langle \text{39} \rangle \) one obtains the estimate \( \langle \text{24} \rangle \)

\[
V_{Mh} \approx g \cdot 0.08 \text{ GeV.}
\]

(40)

A similar estimate is obtained in \( \langle \text{22} \rangle \) for the hybrid-glueball mixing matrix element, while the meson-glueball mixing is second-order in \( \langle \text{10} \rangle \).

Hence the hybrid admixture coefficient \( \langle \text{36} \rangle \) for the meson is

\[
C_{Mh} = \frac{V_{Mh}}{E^{(0)}_M - E^{(0)}_h} = \frac{V_{Mh}}{\Delta M_{Mh}}
\]

(41)

and for the ground state low-lying mesons when \( \Delta M_{Mh} \sim 1 \text{ GeV} \) it is small, \( C_{Mh} \sim 0.1 - 0.15 \), yielding a 1-2\% probability. However for higher states in the region \( M_M \gtrsim 1.5 \text{ GeV} \), the mass difference \( \Delta M_{Mh} \) of mesons and hybrids with the same quantum numbers can be around \( 200 \text{ MeV} \), and the mixing becomes extremely important, also for meson-glueball mixing, which can be written as

\[
C_{MG} = \sum_h \frac{V_{Mh}V_{hG}}{\Delta M_{Mh}\Delta M_{hG}}
\]

(42)

and \( V_{Mh} \sim V_{hG} \). It is clear that the iterative scheme described above can be useful only because hybrid excitation by one additional gluon "costs" around 1 GeV increase in mass, hence the coefficient \( c^N_n \langle \text{36} \rangle \) can be small.
6 Hybrid states and DIS

As was stressed in the previous section, in the large \( N_c \) limit the higher Fock components which are excited by the external current (or incident hadron) are the multihybrid (or \( n \)-hybrid), states. It is convenient to consider these states in the light-cone formalism, following the line of derivation and most notations in [37].

Consider the \( n \)-hybrid with quark at the point \( z^{(a)}_{\mu} \), antiquark at the point \( z^{(b)}_{\mu} \) and gluons at the points \( z^{(k)}_{\mu}, k=1,...,n \). We also define \( \rho^{(i)} = z^{(i)} - z^{(i-1)} \), with \( z^{(0)} \equiv z^{(a)} \) and \( z^{(n+1)} \equiv z^{(b)} \).

The action is

\[
A = K + \sigma S_{\text{min}}
\]

where the kinetic operator \( K \) and the minimal area \( S_{\text{min}} \) can be written as

\[
K = \frac{m_a^2}{2\mu_a} + \frac{m_b^2}{2\mu_b} + \frac{1}{2} \int_0^T dz_+ [\mu_a((\dot{z}_1^{(a)})^2 + 2\dot{z}_-^{(a)}) + \mu_b((\dot{z}_1^{(b)})^2 + 2\dot{z}_-^{(b)})
+ \sum_{i=1}^n \mu_i((\dot{z}_1^{(i)})^2 + 2\dot{z}_-^{(i)})]
\]

\[
\sigma S_{\text{min}} = \frac{1}{2} \int_0^T dz_+ \sum_{i=1}^{n+1} \int_0^1 d\beta_i \left[ \nu_i \left( \left( \dot{w}^{(i)} \right)^2 - \frac{(\dot{w}^{(i)} w'^{(i)})^2}{(w'^{(i)})^2} \right) + \frac{\sigma^2 (w'^{(i)})^2}{\nu_i} \right]
\]

and we have defined

\[
w^{(i)}_\mu = z^{(i-1)}_\mu (1 - \beta_i) + \beta_i z^{(i)}_\mu, \quad \dot{w}^{(i)} = \dot{z}^{(i-1)} - \beta_i \dot{z}^{(i)},
\]

\[
w'^{(i)} = z^{(i)} - z^{(i-1)} \equiv \rho^{(i)}.
\]

As in [37] we introduce the total momentum \( P_+ \),

\[
P_+ = \sum_{i=1}^n \mu_i + \sum_{i=1}^{n+1} \int_0^1 \nu_i d\beta + \mu_a + \mu_b.
\]

At this point one can make a new important step and introduce the parton’s quota \( x_i \) of the total momentum \( P_+ \) to be associated with the Feynman variables \( x_i, \ i = 1,...,n, \ x_a \) and \( x_b \), which can be written as

\[
x_i = \frac{\mu_i + \int_0^1 \nu_i \beta d\beta + \int_0^1 \nu_i (1 - \beta) d\beta}{P_+}
\]
\[ x_a \equiv x_0 = (\mu_a + \int_0^1 \nu_1(1 - \beta)d\beta)/P_+ \] (50)

\[ x_b \equiv x_{n+1} = (\mu_b + \int_0^1 \nu_{n+1}\beta d\beta)/P_+. \] (51)

One can notice that \( x_i \) consists of three pieces: 1) the \((+)\) -component of momentum of the valence gluon \((\mu_i)\), 2) the \((+)\) momentum of the preceeding piece \((i)\) of string weighted with the factor \(\beta\) which takes into account that the string \((i)\) is deformed by the motion of the gluon \((i)\) while another end of the string is fixed 3) the \((+)\) momentum of the string \((i + 1)\) weighted with the factor \((1 - \beta)\) taking into account motion of the \((i \neq 1)\) string due to the \(i\)-th gluon. Note that the parameter \(\beta\) in all strings \((i), i = 1, ...n + 1\) grows in one direction, e.g. from the left to the right.

In this way the momentum of each piece of the string \((i)\) is shared by two adjacent gluons: \((i - 1)\) and \((i)\), so that each parton quota \(x_i\) contains momentum of the parton (gluon or quark or antiquark) itself and of pieces of adjacent strings.

These results imply several nontrivial consequences. First of all, one can see that the einbein factors \(\mu_i\), which played in the c.m. system the role of constituent mass (energy) of gluon and quarks, and \(\nu_i\) played the role of energy density along the string, in the l.c. system they enter directly the Feynman parameters of gluons. In this way one can for the first time see the connection of the standard constituent quark (gluon) picture with the parton picture and calculate as in (50), (51) parton parameters through the (Lorentz boosted) constituent energies of quarks and gluons. Secondly in the l.c. wave-function of the n-hybrid the average values of \(\mu_i\) and \(\nu_i\) are equal for large \(n\), \(\bar{\mu}_i = \bar{\nu}_i, i = 1, 2, ...n\), while for \(n = 1\) one obtains \(\bar{\mu}_g = \sqrt{2}\bar{\mu}_q = \sqrt{2}\bar{\mu}_{\bar{q}}\) [38].

Hence gluons carry more momentum on average than quarks in the n-hybrid state. Therefore one expects that in DIS at large enough energy when the \(n\)-hybrid component of the hadron wave-function is excited, the contribution of gluons to the momentum sum rule and to the proton spin should be dominant. This expectation is consistent with the experimental data. Thus in the neutrino-isoscalar scattering the momentum sum rules for the quark part of \(F_2(x,Q^2)\) yield [39] \(0.44 \pm 0.003\), which implies that gluons carry more than 50% of the total momentum.
For the proton spin one has the relation

\[ \frac{1}{2} = \frac{1}{2} \Delta \Sigma(\mu) + L_q(\mu) + J_g(\mu) \]  

(52)

where the quark sigma-term experimentally is \( \Delta \Sigma(\mu = 1 \text{ GeV}) = 0.2 \pm 0.1 \), and the most part of the difference between the l.h.s. and the r.h.s. is presumably due to the gluon spin contribution \( J_g(\mu) \). The detailed estimates of hybrid contributions to DIS and high-energy scattering will be published elsewhere.

7 Conclusions

It was explained above that the Field Correlator Method is a powerful tool for investigation of all nonperturbative effects in QCD. In particular it provides a natural mechanism of confinement, compatible with all lattice data, and explains the close connection of confinement and chiral symmetry breaking. The spectrum of mesons, glueballs and hybrids is calculated with \( \sigma, \alpha_s \) and current masses as the only fixed parameters used and this spectrum is in good agreement with lattice data and experiment. The latest development concerns the dominant role of hybrids in DIS and high-energy scattering and here the first qualitative results are consistent with experimental evidence. The author is grateful to the organizers of the Pomeranchuk International Conference for their excellent job, and to A.M.Badalian, K.G.Boreskov, A.B.Kaidalov and O.V.Kancheli for many stimulating discussions.

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