HYBRID EXCITATIONS OF THE QCD STRING WITH QUARKS

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Abstract

The model for constituent glue is presented starting from the perturbation theory in nonperturbative QCD vacuum. Green function is constructed for a system containing $q\bar{q}$ pair and gluon propagating in the vacuum background field, and Hamiltonian approach to the problem is formulated. The masses of lowest $q\bar{q}g$ hybrids with light quarks are evaluated, and implication of results for the spectroscopy of exotic states is discussed.

There is no doubts now that the gluonic degrees of freedom are not only present in the QCD Lagrangian, but should also exhibit themselves at the constituent level in such a way that the pure gluonium and hybrid states should exist. Experimental candidates for such exotic states appear from time to time, causing quite natural excitement in the situation when the standard $q\bar{q}$ nonets are already overpopulated. On the other hand, the state-of-art of the QCD in the strong coupling nonperturbative regime still does not allow to predict unambiguously the mass scale of the non-$q\bar{q}$ states.

Estimates for the masses of the hybrid $q\bar{q}g$ mesons were carried out within the various QCD–inspired approaches, such as bag models [1], constituent gluon models [2], flux-tube model [3,4], and QCD sum rules [5], with rather controversial results. For example, the predictions for the masses of lightest hybrids with $u$ and $d$ quarks vary from 1.5 GeV to 2.5 GeV. Having in mind that model approaches usually contain both theoretical and parameter uncertainties, we claim that more direct QCD–motivated treatment is desirable to handle the problem.

In the present letter we study the $q\bar{q}g$ system in the framework of Vacuum Background Correlator method [6], constructing the hadronic Green functions starting from the QCD Lagrangian, with the dynamics of interaction defined by the averages of corresponding Wilson

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loop operators. It may be shown [6,7] that under reasonable assumptions on the behaviour of background vacuum correlators the area law asymptotics for Wilson loop averages at large distances appears, giving rise to the string–type interaction of constituents. The $q\bar{q}$ system with this interaction was described in [8], and the picture of hybrids in this approach was first considered in [9].

We introduce a model for constituent glue starting from the perturbation theory expansion in nonperturbative QCD vacuum [10]. As it is usually done in the background field formalism, we split the gluonic field $A_\mu$ into the background part $B_\mu$ and the perturbation $a_\mu$ over background:

$$ A_\mu = B_\mu + a_\mu. \tag{1} $$

Ascribing the inhomogeneous part of gauge transformation to the field $B_\mu$ we can form gauge invariant states with the field $a_\mu$ involved. One–gluon hybrid is represented as

$$ \Psi(x_q, x_{\bar{q}}, x_g) = \bar{\psi}_\alpha(x_{\bar{q}}) \Phi^\alpha_\beta(x_g, x_q) a^\beta_\gamma(x_g) \Phi^\gamma_\delta(x_g, x_q) \psi_\delta(x_q), \tag{2} $$

where $\alpha...\delta$ are the colour indices in the fundamental representation, spin indices are omitted, $\psi(\bar{\psi})$ stands for the quark (antiquark) field, $a^\beta_\gamma = (\lambda_a)_{\beta}^{\gamma} a_a$, and parallel transporter $\Phi$ contains only background field:

$$ \Phi^\alpha_\beta(x, y) = P \exp \int_y^x B^\alpha_\beta dz_\mu \tag{3} $$

The evolution of the state (2) is described by the Green function

$$ G(x_q x_{\bar{q}} x_g, y_q y_{\bar{q}} y_g) = < \Psi(y_q, y_{\bar{q}}, y_g) \Psi^+(x_q, x_{\bar{q}}, x_g) >_B, \tag{4} $$

where the averaging over background field configurations is implied.

The Green function for the field $a_\mu$ propagating in the given background field $B_\mu$ may be written in the Euclidean space–time as [10]:

$$ G_{\mu\nu}^{-1} = D^2 \delta_{\mu\nu} - g F_{\nu\mu} - D_\nu D_\mu + \frac{1}{\xi} D_\mu D_\nu = M_{\mu\nu} + \frac{1}{\xi} D_\mu D_\nu, \tag{5} $$

with the background gauge fixing term $G_\mu^a = (D_\mu a_\mu)^a$, where $D_\lambda^c = \partial_\lambda \delta^c + g f^{cba} B_\lambda^b$. If the background field satisfies the equation of motion $D_\mu F_{\mu\nu} = 0$, then one has $M_{\mu\nu} D_\nu = 0$, and the Green function (5) may be rewritten as

$$ G_{\mu\nu} = (\delta_{\mu\nu} + D_\mu (\frac{1}{D^2}) D_\lambda \cdot (\xi - 1)) (D^2 - 2 g F)_{\mu\nu}^{-1}. \tag{6} $$

The choice $\xi = 0$ corresponds to the Landau gauge, in which the Green function (6) contains explicitly the projector $P_{\mu\lambda}$ onto transverse states:

$$ P_{\mu\lambda} = \delta_{\mu\lambda} - D_\mu (\frac{1}{D^2}) D_\lambda \tag{7} $$

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3The condition $D_\mu F_{\mu\nu} = 0$ is claimed to be not necessary [11] and may be removed using the Coulomb background gauge, which has an additional advantage: in this gauge the ghosts can be decoupled explicitly.
The next step is to use the Feynman–Schwinger representation [12] of Green function (4) to define the effective action for the $q\bar{q}g$ system. Here we present the simplified version of the model, neglecting spin degrees of freedom and reducing the problem to the effective scalar one; this approximation corresponds to omitting the projector in the gluon Green function (6) and colour magnetic interaction (term proportional to $gF$). In a similar way we omit the spin dependence in the quark Green function, assuming

$$G_q = (D^2 - m_q^2)^{-1}$$

Within this approximation the Feynman–Schwinger representation for the Green function (4) takes the form

$$G(x_q x_{\bar{q}} x_g ; y_q y_{\bar{q}} y_g) = \int_0^\infty ds \int_0^\infty d\bar{s} \int_0^\infty dS \int D\bar{z}DzD\bar{Z}DZ \exp(-K_q - K_{\bar{q}} - K_g) < W >_B,$$

where

$$K_q = m_q^2 s + \frac{1}{4} \int_0^s \dot{z}^2(\tau)d\tau, \quad K_{\bar{q}} = m_{\bar{q}}^2 \bar{s} + \frac{1}{4} \int_0^{\bar{s}} \dot{\bar{z}}^2(\tau)d\tau, \quad K_g = \frac{1}{4} \int_0^S \dot{Z}^2(\tau)d\tau,$$

boundary conditions read

$$z(0) = x_q, \quad \bar{z}(0) = x_{\bar{q}}, \quad Z(0) = x_g,$$

$$z(s) = y_q, \quad \bar{z}(\bar{s}) = y_{\bar{q}}, \quad Z(S) = y_g,$$

and all the dependence on the gluonic fields $B$ is contained in the Wilson loop operator

$$W = \lambda_a \Phi_{\Gamma_q}(y_q, x_q) \lambda_b \Phi_{\Gamma_{\bar{q}}}(x_{\bar{q}}, y_{\bar{q}}) \Phi_{ab, \Gamma_g}(y_g, x_g),$$

the contours $\Gamma_q, \Gamma_{\bar{q}}$ and $\Gamma_g$ run over the trajectories of the quarks and the gluon correspondingly (see Fig.), $\Phi_{ab}$ is the transporter in the adjoint representation.

The Wilson loop configuration (10) can be disentangled using the relation between ordered exponents along the gluon path $\Gamma_g$ in the adjoint and fundamental representation:

$$\frac{1}{2} \Phi_{ab, \Gamma_g}(x, y) = tr(\lambda_a \Phi_{\Gamma_q}(x, y) \lambda_b \Phi_{\Gamma_{\bar{q}}}(y, x)),$$

(path $\Gamma_{\bar{g}}$ is directed oppositely to $\Gamma_g$), with the result

$$< W >_B = \frac{1}{2} < (W_1 W_2 - \frac{1}{N_c} W) >_B,$$

where $W_1, W_2$ and $W$ are the Wilson loops in the fundamental representation along the closed contours $C_1, C_2$ and $C$ shown at the Figure. The averaging over background can be done using the cluster expansion method [6]. The average of two Wilson loops was calculated in [7], and it was demonstrated that, assuming the existence of finite gluonic correlation length $T_g$, the generalized area law asymptotics may be obtained. For the Wilson loop average (12) one has (when both contours $C_1$ and $C_2$ are large):

$$< W > = \frac{N_c^2 - 1}{2} exp(-\sigma(S_1 + S_2)),$$
where $\sigma$ is the string tension in the fundamental representation, and $S_1$ and $S_2$ are the minimal areas inside the contours $C_1$ and $C_2$. The area law (13) holds for practically all the reasonable configurations in the $q\bar{q}g$ system, apart from the special case of the contours $C_1$ and $C_2$ embedded into the same plane, when the expression (13) is replaced by

$$\langle W \rangle = \frac{N_c^2 - 1}{2} \exp(-\sigma(S_1 - S_2) - \sigma_{\text{adj}} S_2), \quad S_1 > S_2 \quad (14)$$

$\sigma_{\text{adj}}$ is the string tension in the adjoint representation. Both regimes (13) and (14) match smoothly each other at the average distances between the contours $C_1$ and $C_2$ of order of correlation length $T_g$. Note that to derive the expression (13) we don’t use the standard procedure [13] of appealing to the limit $N_c \to \infty$ and replacing the adjoint line in (10) by the double fundamental one. In actual calculations we, however, assume the regime (13) to be valid everywhere in the $q\bar{q}g$ configuration space, arriving in such a way to the picture of the $q\bar{q}g$ hybrid as a system of a gluon with two minimal strings attached, each having a quark (or antiquark) at the end.

To reduce the four–dimensional dynamics in (9) to the three–dimensional one we follow the procedure outlined in [8], using the parametrization for which

$$z_\mu = (\tau, \vec{r}_q), \quad \bar{z}_\mu = (\tau, \vec{\bar{r}}_q), \quad Z_\mu = (\tau, \vec{r}_g),$$

and introducing new dynamical variables

$$\mu_1(\tau) = \frac{T}{2s} \dot{z}_0(\tau), \quad \mu_2(\tau) = \frac{T}{2s} \dot{\bar{z}}_0(\tau), \quad \mu_3(\tau) = \frac{T}{2s} \dot{Z}_0(\tau)$$

with $0 \leq \tau \leq T$. The representation (9) for the Green function is now rewritten as

$$G = \int D\vec{r}_q D\vec{\bar{r}}_q D\vec{r}_g D\mu_1 D\mu_2 D\mu_3 \exp(-A), \quad (15)$$

with the effective action

$$A = \int_0^T d\tau \left\{ \frac{m_q^2}{2\mu_1} + \frac{m_{\bar{q}}^2}{2\mu_2} + \frac{\mu_1 r_q^2}{2} + \frac{\mu_2 r_{\bar{q}}^2}{2} + \frac{\mu_3 r_g^2}{2} + \right.$$

$$+ \sigma \int_0^1 d\beta_1 \sqrt{\dot{w}_1^2 w_1'^2 - (\dot{w}_1 w_1')^2} + \sigma \int_0^1 d\beta_2 \sqrt{\dot{w}_2^2 w_2'^2 - (\dot{w}_2 w_2')^2},$$

where the surfaces $S_1$ and $S_2$ are parametrized by the coordinated $w_{i\mu}$, and

$$\dot{w}_{i\mu} = \frac{\partial w_{i\mu}}{\partial \tau}, \quad w_{i\mu}' = \frac{\partial w_{i\mu}}{\partial \beta_i}, \quad i = 1, 2.$$

Assuming the straight–line approximation for the minimal surfaces,

$$w_{1\mu}(\tau, \beta_1) = z_\mu(\tau) \beta_1 + Z_\mu(\tau)(1 - \beta_1), \quad w_{2\mu}(\tau, \beta_2) = \bar{z}_\mu(\tau) \beta_2 + \bar{Z}_\mu(\tau)(1 - \beta_2),$$

we arrive to the Lagrangian

$$\mathcal{L} = \frac{m_q^2}{2\mu_1} + \frac{m_{\bar{q}}^2}{2\mu_2} + \frac{\mu_1 r_q^2}{2} + \frac{\mu_2 r_{\bar{q}}^2}{2} + \frac{\mu_3 r_g^2}{2} + \quad (17)$$
\[ + \sigma \rho_1 \int_0^1 d \beta_1 \sqrt{1 + l_1^2} + \sigma \rho_2 \int_0^1 d \beta_2 \sqrt{1 + l_2^2}; \]

\[ l_1^2 = (\beta_1 \dot{r}_q + (1 - \beta_1) \dot{r}_g)^2 - \frac{1}{\rho_1^2} (\beta_1 (\dot{r}_q \bar{p}_1) + (1 - \beta_1) (\dot{r}_g \bar{p}_1))^2; \]

\[ l_2^2 = (\beta_2 \dot{r}_q + (1 - \beta_2) \dot{r}_g)^2 - \frac{1}{\rho_2^2} (\beta_2 (\dot{r}_q \bar{p}_2) + (1 - \beta_2) (\dot{r}_g \bar{p}_2))^2; \]

\[ \bar{p}_1 = r_q - \bar{r}_q, \quad \bar{p}_2 = r_q - \bar{r}_q. \]

The Lagrangian (17) is the straightforward generalization of the Lagrangian obtained in [8] for the \( q\bar{q} \) string.

To deal with the square root terms in (17) the auxiliary field approach was used in [8]. Here we are interested only in the masses of lowest states, so we adopt another strategy, which allows simple variational treatment of the problem.

First, the experience borrowed from the results of papers [8] tells us that for low values of relative angular momenta the square roots in (17) can be expanded up to the second order in the transverse velocities \( \vec{l}_i \), and it was proved to be a good approximation even for the massless constituents. This approximation corresponds to the potential–like regime with the linear potential \( V = \sigma \rho_1 + \sigma \rho_2 \), while the terms \( \sim l_i^2 \) (string corrections) can be taken into account as a perturbation. With this simplification the Lagrangian (17) becomes quadratic in velocities, the centre–of–mass motion is decoupled, and the Hamiltonian (in the Minkowsky space–time) can be easily obtained from (17):

\[ H_0 = \frac{m_q^2}{2 \mu_1} + \frac{m_\bar{q}^2}{2 \mu_2} + \frac{\mu_1 + \mu_2 + \mu_3}{2} + \frac{p^2}{2 \mu_p} + \frac{Q^2}{2 \mu_Q} + \sigma \rho_1 + \sigma \rho_2 \quad (18) \]

where the Jacobi coordinates \( \vec{r} \) and \( \vec{\rho} \) and conjugated momenta \( \vec{\rho} \) and \( \bar{Q} \) are introduced:

\[ \vec{r} = \vec{r}_q - \vec{r}_g, \quad \vec{\rho} = \vec{r}_g - \frac{\mu_1 \vec{r}_q + \mu_2 \vec{r}_g}{\mu_1 + \mu_2}; \]

\[ \mu_p = \frac{\mu_1 \mu_2}{\mu_1 + \mu_2}, \quad \mu_Q = \frac{\mu_3 (\mu_1 + \mu_2)}{\mu_1 + \mu_2 + \mu_3}; \]

\[ \bar{p}_1 = -\vec{\rho} + \frac{\mu_2}{\mu_1 + \mu_2} \vec{r}, \quad \bar{p}_2 = -\vec{\rho} - \frac{\mu_1}{\mu_1 + \mu_2} \vec{r}. \]

The Hamiltonian (18) contains the fields \( \mu_i(\tau) \), and one is to integrate over these fields in the path integral representation. Instead we proceed in a way suggested in [14]: we find eigenvalues \( \varepsilon_0(\mu_i) \) of the Hamiltonian \( H_0 \) variationally, assuming \( \mu_i \) independent of \( \tau \), and then minimize \( \varepsilon_0(\mu_i) \) with respect to \( \mu_i \). Such procedure works nicely for the ground states, with the accuracy about several percent, and, apart from being very simple technically, allows for the approximate solution of the problem of defining the physical transverse states.

Indeed, we have neglected the spin dependence in the gluon Green function (6). Thus we do not expect the eigenfunctions of the Hamiltonian (18) to be transverse and they are really not. On the other hand, within the framework of the variational procedure described above we can impose the constraint

\[ \mu_3 \Psi_0 = \mu_3 (\vec{r}_g \bar{\psi}) = 0, \quad (19) \]
projecting out the physical hybrid state $\Psi_\lambda = (\Psi_0, \vec{\Psi})$ (we restore for the moment the gluon spin index $\lambda$ in the state (2)). The constraint (19) is compatible with the projector $P_{\mu\lambda}(\text{eq.}(7))$(after averaging over background and introducing the variables $\mu_i$). In accordance with the condition (19) and neglecting string corrections we choose the physical states to be transverse with respect to the (three–dimensional) gluon momentum:

$$\vec{p}_3 \vec{\Psi} = 0, \quad \Psi_0 = 0,$$

so that the states contain electric or magnetic gluons:

$$\vec{\Psi}^e_j \sim \vec{\Psi}^m_j \sim \sqrt{\frac{j+1}{2j+1}} Y_{jjm}^j(\hat{Q}) + \sqrt{\frac{j}{2j+1}} Y_{jjm+1}(\hat{Q})$$

where $j$ is the total momentum in the gluonic subsystem. (See the paper [15], where the gluonium spectrum was calculated within the similar assumptions).

We stress that the choice (20), though looking quite natural, is nothing but the variational ansatz, and, in principle, cannot substitute for rigorous treatment of the problem of transverse and longitudinal gluonic degrees of freedom. The latter should be done with inclusion of spin variables into the path integral representation for the gluon Green function.

As the illustrative example of our approach we present the results for the masses of lowest hybrids with $u$ and $d$ quarks, assuming the latters to be massless. The radial part of the three–body wave function was chosen to be of the simple cluster form $\sim \exp(-\alpha \rho_1^2 - \alpha \rho_2^2)$ (up to the centrifugal barrier), and the angular momenta were taken to be the lowest compatible with the ansatz (21) or (22). So the ground state hybrid contains a $1^+ \text{ or } 1^- \text{ gluon and a quark–antiquark pair with relative angular momentum } l = 0$. More elaborated calculations for massive quarks, more realistic trial wave function and with inclusion of string corrections and short–range Coulomb force are in progress now and will be reported elsewhere.

The parity and charge conjugation of the hybrids are given by

$$P = (-1)^{l+j}, \quad C = (-1)^{l+s+1}$$

for the states with electric gluon (21), and

$$P = (-1)^{l+j+1}, \quad C = (-1)^{l+s+1}$$

for the states with magnetic gluon (22), where $s$ is the total spin of quarks. So the possible quantum numbers for the ground state hybrids are

$$J^{PC} = 0^{++}, \quad 1^{++}, \quad 2^{++}, \quad 1^{+-}.$$  

With the chosen trial function these states are degenerate, and we expect this degeneracy to be removed mainly by spin–dependent effects. For the string tension $\sigma = 0.2 GeV^2$ the eigenenergy of the ground state is estimated to be $E_0 = 3.3 GeV$.

Now we enter the rather delicate point of setting the absolute scale for the mass of a hybrid. It is well–known that in the potential models the large and negative additive constant in the potential is needed to fit the meson spectra. The same holds for the lowest $q\bar{q}$’s in the suggested
picture: the negative constant $C \approx -0.7 \div -0.8$ GeV is needed to describe the $S-,P-$ and $D-qq$ levels. At present there is no reasonable model for this constant term, and the best way would be to adjust it by some hybrid candidate mass. In principle, the constant may be attributed to the perimeter term in the formula for Wilson loop, or to the hadronic shift (creation of additional hadronic loops due to the string breaking). For the $q\bar{q}g$ system with two strings the most naive recipe to account for such mechanisms is that the constant is twice as large as for the $q\bar{q}$ system. Taking this as an educational guess we end up with the mass

$$M_0 = 1.7 - 1.8\text{GeV}$$

for the ground $q\bar{q}g$ state.

The value (26) should not of course be taken seriously. However, it appears to be rather close to the mass range where the hybrid candidates are believed to be placed, and to the value obtained in the flux tube model [3,4]. Ideologically, our approach and the flux tube model have common features: both models account for the transverse motion of the $q\bar{q}$ string, though the idea is realized in different ways. The most important difference is in quantum numbers: there is no constituent gluon in the flux tube, so instead of (25) one has [3]

$$J^{PC} = 0^{\mp\pm}, 1^{\mp\pm}, 2^{\mp\pm}, 1^{\pm\pm}$$

(27)

It is not easy to distinguish between models using the existing data in light quark sector. Indeed, at present several hybrid candidates in the mass range 1.5-2 GeV are under discussion: $1^{-+}$ exotic signal from BNL experiment [16], the $0^{-+}$ and $2^{-+}$ states seen by VES collaboration [17], the $\pi(1775)$ state [18] which may be $1^{-+}$, but $0^{-+}$ or $2^{-+}$ assignment is not excluded, and additional vector states seen in $e^+e^-$ annihilation around 1.5-1.7 GeV [19]. The hybrid assignment is attached to these states basing on the very clear signature for hybrid decays in the flux tube model [20]: the decay of a hybrid into two $S$-wave mesons is suppressed, and it is really the case for the listed states. Nevertheless, just the same signature exists for a hybrid with electric constituent gluon [21,22], that means that the $P$-odd ground state hybrids (25) exhibit the same decay pattern. So to tell one model from another one is to study the resonance activity in $0^{++}, 1^{++}, 2^{++}, 1^{+-}/0^{-+}, 1^{+-}, 2^{+-}, 1^{++}$ channels around 1.5-2.0 GeV, a rather challenging experimental task.

More promising is to use another selection rule, valid only for the decay of constituent hybrids [21,22]. There are two decay mechanisms. One proceeds via string breaking, leads to the $q\bar{q} - q\bar{q}g$ final state and is forbidden for the ground state because of phase space. Another is due to the conversion of a gluon into $q\bar{q}$ pair. The constituent gluon behaves as the massive particle, transferring its spin to the created pair, and it means that not only the total angular momentum but also the total spin of constituents conserved, giving rise to the set of selection rules for hybrid decays, similar to those obtained in [22].

To conclude, we have demonstrated that the model for QCD quark–antiquark string can be extended to account for hybrid meson excitations. The arising model of constituent glue appears to be compatible with the data on light quark meson spectroscopy.

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Figure caption

Wilson loop configuration corresponding to the propagation of the hybrid state.