Nonperturbative mechanisms of strong decays 
in QCD

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Abstract

Three decay mechanisms are derived systematically from the QCD Lagrangian using the field correlator method. Resulting operators contain no arbitrary parameters and depend only on characteristics of field correlators known from lattice and analytic calculations. When compared to existing phenomenological models, parameters are in good agreement with the corresponding fitted values.

1 Introduction

An enormous amount of experimental data on strong decays of mesons and baryons is partly used by theoreticians for comparison in the framework of the $^3P_0$ model [1], and its flux-tube modifications [2]. The analysis done in [3] confirmed the general validity of the model, whereas in [4] results of other forms of decay operators have been also investigated in meson decays, and in [5] in baryon decays. On the whole, the phenomenological picture seems to be satisfactory for the $^3P_0$ model with some exclusions discussed in [3] and [5]. The recent extensive study of strong decays of strange quarkonia based on the $^3P_0$ model was done in [4].

The key element which is missing in this situation is the systematic derivation of all terms in the decay Hamiltonian from the basic principles, i.e. from the QCD Lagrangian. It is the purpose of the present paper to make some progress in this direction using the Field Correlator Method (FCM) [8] and
background perturbation theory \cite{9} to treat nonperturbative (NP) QCD contributions together with perturbative ones.

In doing so one should take into account the special role of pions in the hadron decays and therefore to perform accurately the chiral bosonization of the effective quark Lagrangian, obtained from the basic QCD Lagrangian. This will give the first term in the decay Hamiltonian, and the corresponding decay mechanism will be referred to as a Chiral Decay Mechanism (CDM). At the same time one should take into account the string degrees of freedom in the original meson and the possibility of string breaking due to the $q\bar{q}$ pair creation. The corresponding term in the decay Hamiltonian will be derived below without free parameters and this second mechanism will be called the String Breaking Mechanism (SBM). As will be seen, the dominant term of SBM has the structure, which can be compared quantitatively with the phenomenological fits of the $^3P_0$ model in \cite{3,4}.

Finally, the QCD perturbation theory in the perturbative background developed in \cite{9}, allows to derive two additional terms in the decay Hamiltonian: one for the OZI-allowed decays, which has the Lorentz form of the $^3S_1$ type but proceeds through the intermediate hybrid state, and another for the OZI-forbidden decays, which proceeds through the intermediate glueball state and only at very small distances reduces to the two-gluon or three-gluon $q\bar{q}$ pair creation. We shall call these mechanisms the Hybrid Mediated Decay (HMD) and the Glueball Mediated Decay (GMD) respectively.

### 2 The Chiral Bosonization and the Chiral Decay Mechanism

One starts with the QCD Lagrangian in the Euclidean space-time and averages over the gluonic fields writing the general form of the gauge-invariant correlator (known from lattice or analytic calculations, see refs. in \cite{3}) which contains the confining part $D(x)$, namely,

$$
\frac{g^2}{N_c} \langle \text{tr}(F_{\mu\nu}(x)\Phi(x,y)F_{\lambda\sigma}(y)\Phi(y,x)) \rangle = (\delta_{\mu\lambda}\delta_{\nu\sigma} - \delta_{\mu\sigma}\delta_{\nu\lambda})D(x-y) + O(D_1) \tag{1}
$$

where $O(D_1)$ contains a relatively small nonconfining part $D_1(x)$ and $\Phi(x,y) = P \exp ig \int_y^x A_\mu dz_\mu$ is the parallel transporter.
Assuming also that all higher correlators can be neglected (as it is supported by lattice data, see [10]) one obtains the effective action [11]

\[ \mathcal{L}_{\text{EQL}}^{(2)} = \frac{1}{2N_c} \int d^4x \int d^4y f(x) f(y) \gamma^{\alpha\beta} \gamma^{\gamma\varepsilon} J(x,y) \]

where the kernel \( J(x,y) \) is expressed through \( D(x) \),

\[ J_{E(M)}(x,y) = \int_{\mathcal{Y}} d\nu_i \int_{\mathcal{Y}} d\nu_i K_{E(M)}(u,v) D(u-v), \quad i = 1, 2, 3. \]

and \( D(x) \) is connected to the string tension \( \sigma \) in the usual way, \( \sigma = \frac{1}{2} \int d^2 x D(x) \).

Here the string kernel \( J_{E(M)} \) contains color electric (magnetic) fields and the former is dominant at large distances, and therefore magnetic part will be omitted for simplicity. At this point one should specify the choice of integration contours, the kernel \( K \) and the point \( \mathcal{Y} \) in (3). As it is known, the full result of the integration over gluon fields does not depend on the gauge and on the shape of the contours (when all correlators are taken into account); however to write the contribution of bilocal correlator (1) in the gauge-invariant form, one has to use one of the variants of the contour gauge, e.g. [12], where \( K_{E} \equiv 1 \), and to choose the contour corresponding to the minimal string which minimizes the contribution of higher correlators. In this section we consider the following geometry: the quarks at the points \( x, y \) in (2) are at one end of the string (they are dynamically close [11]), while the point \( \mathcal{Y} \) is at the position of the heavy antiquark, as it was introduced in [11]. In the next section we shall consider a more general geometry, when the points \( \mathcal{Y} \) in (3) are different and integration over \( du_i \) and \( dv_i \) runs over two different pieces of the broken string. For the results of the present section the exact definition of the point \( \mathcal{Y} \) is inessential since the pion is emitted from the end of the string under consideration while another end of string is a spectator.

As the next step the bosonization of the Lagrangian (2) can be done in the usual way, however with nonlocal mesonic fields with the result [13]

\[ \Delta L = i \int d^4x d^4y \psi^+(x) \hat{M}(x,y) \psi(y) \]

where the kernel \( \hat{M}(x,y) \) can be written as a nonlinear form for PS meson fields

\[ \hat{M}(x,y) = M_S(x,y) \exp(i \gamma_5 \hat{\phi}(x,y)) + ..., \quad \hat{\phi} = \frac{\hat{\phi}_A \lambda^A}{F_\pi}, \quad F_\pi = 93\text{MeV}. \]
Here ellipsis implies all other omitted terms, including those containing isovector scalars, vector and pseudovectors.

One should note, that in our case when confinement is present the constant condensate of scalar-isoscalar field always enters multiplied by $J(x,y)$ and thus produces the scalar confining potential of the string in $M_S(x,y)$ \[13\]. Namely, for long enough string, i.e. for $|x - Y| \gg T_g$, where $T_g$ is the gluonic correlation length in $D(x)$, one has approximately \[13\]

$$M_S(x,y) \approx \sigma|x|\delta^4(x - y),$$

This is different from the instanton, or the NJL model, where the scalar-isoscalar field acquires a nonzero condensate, which is constant in all space-time.

Inserting (6) into (5) one obtains a localized form of the quark-meson Lagrangian, describing interaction on one end of the string; to the lowest order in PS field,

$$\Delta L^{(1)} = \int \overline{\psi}(x)\sigma|x - Y|\gamma_5\phi^A\lambda^A\frac{\phi^A\lambda^A}{F^2}\psi(x)dtd^3x,$$ \[7\]

One can visualize in (7) the simultaneous presence of quark fields $\psi, \bar{\psi}$, together with the string $\sigma|x - Y|$ and Nambu-Goldstone (NG) fields $\phi^A, A = 1, \ldots, n_f^2 - 1$.

Using Dirac equation for the quark field $\psi(x)$ one arrives as in \[13\] at the familiar Weinberg Lagrangian \[14\]

$$\Delta L^{ch} = g_A^q tr(\overline{\psi}\gamma_\mu\gamma_5\omega_\mu\psi), \omega_\mu = \frac{i}{2}(u\partial_\mu u^+ - u^+\partial_\mu u),$$ \[8\]

where $u(x) = \exp(\frac{i}{2}\gamma_5\phi(x,x))$. In our derivation $g_A^q$ is uniquely defined in the local limit,

$$g_A^q \equiv 1$$ \[9\]

which agrees with large $N_c$ limit, discussed in \[14\].

Note that both Lagrangians (7), (8) are local limits of nonlocal expression (4), and for not very long string with the length $L \sim T_g \sim 0.2$ fm the nonlocality is essential. However string is apparently not present in (8), both quark operators there are solutions of Dirac equation with the string entering as a scalar potential. It is clear that the Lagrangian (8) describes the pion field emission both from the quark at one end of the string and from the
antiquark at another end of the string. In case of baryons one should sum up in (8) over all three quarks.

The form (8) was used in [15] for the calculation of pionic transitions in the heavy-light mesons with \( g_A \) playing role of a fitting parameter, which turned out to be around 0.7. This difference from (9) can be considered as an indication of a possible role of nonlocality.

3 The String Breaking Mechanism

This mechanism was considered in some detail in [16] (see also refs. therein). In the present section we shall consider the pair creation vertex due to the nonperturbative QCD configurations. In the \( ^3P_0 \) model [1] this vertex was modelled by an adjustable constant and in [16] by some function \( F \). It is our purpose here to derive this vertex from the basic \( 4q \) effective action (2).

To describe the creation of the \( q\bar{q} \) pair in the presence of the string which connects quark \( Q \) at the point \( X \) and antiquark \( \bar{Q} \) at the point \( \bar{X} \), we can also choose the contours of integration along the strings: e.g. \( A_4(x) = \int_X^x du_1 E_1(u_1), A_4(y) = \int_X^y dv_1 E_1(v_1) \), and the string is along the axis 1. (This change of contours from (3) can be traced to the effect of cancellation in the sum of contours from the quark \( Q \) to the point \( Y \) and with opposite sign –from the antiquark \( \bar{q} \) to the point \( Y \), which results in the contour integral between positions \( Q(X) \) and \( \bar{q}(x) \)). As a result the kernel \( J_E \) is

\[
J_E(x, x_4; y, y_4) = \int_x^X du_1 \int_y^X dv_1 D(u_1 - v_1) \simeq \frac{\sigma}{\pi} \exp\left(-\frac{(x_4 - y_4)^2}{4T_g^2}\right) \tag{10}
\]

where for \( D(x) \) the Gaussian form was used, c.f. [14]. As in [9] one can find the effective mass operator (due to color electric fields) \( M(x, y) \), using the definition [12],

\[
M(x, y) = -i\gamma_4 S(x, y)\gamma_4 J(x, y) \tag{11}
\]
and the estimate of the quark Green’s function \( S(x, y) \) for a long string, \( L \gg T_g \), done in \cite{11, 13}, gives \( S(x, y) \sim i\delta^{(3)}(x - y) \).

As a result one obtains for the effective Lagrangian the same form as in \cite{4},

\[
\Delta L^{(SBM)} = i \int d^4x d^4y \bar{\psi}^+(x)M^{(br)}(x, y)\psi_a(y)
\]  

(12)

where

\[
M^{(br)}(x) = \frac{\sigma}{\pi}\delta^{(3)}(x - y)e^{-\frac{(x_4 - y_4)^2}{4T_g^2}}
\]  

(13)

Integrating over \( d(x_4 - y_4) \) one gets in the Minkowskian space-time

\[
\Delta L^{(SBM)} = \frac{2T_g\sigma}{\sqrt{\pi}} \int d^4x \bar{\psi}(x)\psi(x)
\]  

(14)

The form (14) coincides with that assumed in the \(^3P_0\) model. At this point one should take into account that the pair in (14) is created at any point \( x \) in the space, and by the derivation this point should lie on the world surface of the string, deforming it from the general minimal area shape to the minimal area surface passing through the point \( x, x_4 \).

The probability amplitude for such a deformation is equal to \( \exp(-\Delta A(x)) \), where \( \Delta A(x) = \sigma\Delta S \) is the increase of the area \( \Delta S \). Assuming the latter is described by the cone with the height \( h \) and radius \( R \), \( h \ll R \), one obtains an additional factor to be inserted in (14), namely,

\[
F^P \equiv \exp(-\Delta A) \simeq \exp(-\sigma\frac{\pi}{2} h^2)
\]  

(15)

This factor is similar to the one suggested in the \(^3P_0\) flux-tube model \cite{2}. Integrating over \( h \) in (14) yields the final result, which we write in the same form as in the \(^3P_0\) model \cite{4},

\[
H^{(SBM)} = g \int d^3x \bar{\psi}(x)\psi(x); \quad g = 2T_g\sigma/\sqrt{\pi}; \quad \gamma = \frac{T_g\sigma}{\sqrt{\pi}\mu} = \frac{g}{2\mu},
\]  

(16)

where \( \mu \) is the constituent quark mass computed through \( \sigma \) \cite{18}. Taking \( T_g \sim (0.2-0.3) \) fm, \( \sigma \approx 0.2 \) GeV\(^2\) and \( \mu \approx 0.35 \) GeV one obtains \( \gamma \sim 0.3-0.5 \), i.e. the values in the same ballpark as in the phenomenological analysis \cite{3, 4}. Several remarks are now in order. Firstly, above the simplified Gaussian form of \( D(x) \) was used in \cite{11}, which requires the redefinition of \( T_g \); secondly, the contribution of magnetic correlators and the term \( D_1 \) in (1) has been neglected. These additional contributions will be considered elsewhere.
4 “Perturbative” pair creation via the hybrid or glueball formation

Two decay mechanisms discussed in previous sections are of purely nonperturbative origin, now we turn to the mechanisms which contain as a limit a purely perturbative $q\bar{q}$ pair creation by a gluon. At large distances one has to know how this process is modified by the presence of nonperturbative confining fields and to this end we shall use the Background Perturbation Theory (BPTh) developed in [10], where the total gluonic field $A_\mu$ is separated into valence gluon field $a_\mu$ and background $B_\mu$, $A_\mu = B_\mu + a_\mu$. The field $B_\mu$ saturates correlator $D(x)$ and contains therefore its own mass scale, while perturbation theory is done in powers of $g a_\mu$. The main physical outcome of the analysis of [10] is that the valence gluons are propagating inside the film of the string worldsheet, so that all Feynman diagrams in the coordinate space can be considered as filled inside by this film on the minimal surface with boundaries specified by quark and gluon trajectories.

At this point it is clear that one should use the path integral representation for the quark and gluon Green’s function, namely the Fock-Feynman-Schwinger (FFS) formalism [19]. The FFS method has proved useful in conjunction with the BPTh to study meson, hybrid and glueball Green’s functions (see [18] for review). In [20] this method has been exploited to calculate mixing between meson, hybrid and glueball states, and in what follows we shall pursue the same way to study matrix elements of decays proceeding via hybrid and glueball intermediate states.

We start with the OZI allowed planar pair creation mechanism by a gluon, propagating inside the film in a hybrid state $\varphi^{(H)}$, and therefore the matrix element for the meson decay via the hybrid states can be written as

$$W^{(H)} = \sum_n \lambda_n^{(MH)} W_n^{(H)}$$

(17)

where $\lambda_n^{(MH)}$ is the dimensionless mixing coefficient of the $n$-th hybrid state in the given initial meson, which according to [20] can be written as

$$\lambda_n^{(MH)} = \frac{V^{(\mu)}_{on}}{\sqrt{2\mu_g(n)|M_H^{(n)} - M_M|}}$$

(18)

Here $M_H^{(n)}$, $M_M$ are hybrid and meson masses respectively and $\mu_g(n)$ is the constituent gluon mass in the $n$-th hybrids state, computed through the string tension as in [18], [20].
The matrix element $V_{on}^{(\mu)}$, introduced in [20], is that of the pair creation operator $H_1$,
\[ H_1 = g \int \bar{q}(\mathbf{x}, 0) \gamma^\mu \gamma^0 q(\mathbf{x}, 0) d^3x \] (19)
between the meson state $\Phi_{\alpha\beta}^{(M)}(\mathbf{r})$ and the hybrid state $\Phi_{\mu,\gamma\delta}^{(H)}(\mathbf{r}_1, \mathbf{r}_2)$, where we have specified quark Dirac indices $\alpha\beta$, $\gamma\delta$ and gluon component index $\mu$ (this is the $4 \times 4$ representation typical for the Bethe-Salpeter wave functions which is used to build up $2j + 1$ components of meson and hybrid wave functions, for discussion and refs. see [20]).

In a simplified form $V_{on}$ can be written as [21]
\[ V_{on} = g \int d^3r \phi^{*(H)}(0, \mathbf{r}) \Gamma \phi^{(M)}(\mathbf{r}) d^3\mathbf{r} + \text{perm.} \] (20)
where "perm" implies another term with gluon emitted by antiquark, and $\Gamma = \gamma_{\mu}$.

Similarly the term $W_n^{(H)}$ in (17) is the decay amplitude of the $n$-th hybrid state into two mesons, which can be written as
\[ W_n^{(H)} = \frac{V_{n,12}}{\sqrt{2\mu_g(n)V}}, \quad N_{n,12} \equiv \int \Phi^{*(M_1)}(\mathbf{r}_1) \Phi^{*(M_2)}(\mathbf{r}_2) \Gamma^{(H)}(\mathbf{r}_1, \mathbf{r}_2) d^3\mathbf{r}_1 d^3\mathbf{r}_2 \] (21)
Normalization of wave functions is (summation over repeating Dirac indices, which are omitted, is implied)
\[ \int |\Phi^{(H)}|^2 d^3\mathbf{r}_1 d^3\mathbf{r}_2 = 1, \quad \int |\Phi^{(M_i)}|^2 d^3\mathbf{r} = 1 \] (22)

Finally two-body decay probability is
\[ w = 2\pi |W^{(H)}|^2 \delta(E_1 - E_2 - M) \frac{V d^3k}{(2\pi)^3} \] (23)

As it is seen from (17), (18), the decay probability is strongly enhanced when the decaying meson mass is in the vicinity of some hybrid level. One should stress that this situation is standard in the mass range above 1.4 GeV, where the density of hybrid ground and excited states is fast growing with mass, (see e.g. second ref. in [3]).

We now turn to the OZI forbidden, i.e. nonplanar $q\bar{q}$ pair creation via valence gluons. The essential step in this mechanism is the creation of the new
flavour quark-antiquark pair with the new string between them, hence the creation operator contains at least a two-gluon exchange. In the confining background trajectories of these two gluons are connected by the adjoint string (or, equivalently at large $N_c$ by a double fundamental string) and therefore a new meson is created by the two- or more gluon glueball, which is emitted virtually from the original meson. The amplitude for this process can be written as follows,

$$W^{(G)} = \sum_n \frac{V_{nMM_2} V_{nGM_1}}{M - E_n^G - E(M_2)}$$

(24)

where notation is used

$$V_{nMM_2} = \langle \phi^{(M_2)} | \Lambda^{(G)}_n | \phi^{(M)} \rangle,$$

$$V_{nGM_1} = (\phi^{(M_1)} \Lambda^{(G)}_n)^*$$

(25)

and the two-gluon glueball vertex is

$$\Lambda^{(G)}_{n,\alpha\beta}(x_1, x_2) = \Psi^{(G)}_{n,\nu_1\nu_2}(x_1, x_2) (\gamma_{\nu_1} S^{(f)}(x_1, x_2) \gamma_{\nu_2})_{\alpha\beta}$$

$$\Lambda^{(G)}_{n,\alpha\beta}(y_1, y_2) = (\gamma_{\mu_1} S^{(g)}(y_1, y_2) \gamma_{\mu_2})_{\alpha\beta} \Psi^{(G)}_{n,\mu_1\mu_2}(y_1, y_2)$$

(26)

Here $\Psi^{(G)}$ is the glueball wave function, and $S^{(g,f)}$ - the Green’s function of quarks with flavor $(g, f)$.

When all distances $|y_i - x_k|, |x_1 - x_2|, |y_1 - y_2|$ are small as compared to $T_g$, the two-gluon Green’s function reduces to the two-gluon exchange of free gluons. This limiting perturbative mechanism is known for a long time [21]. In the opposite limit only the lowest mass term survives in the sum (24).

The most interesting case occurs when the denominator in (24) becomes small, which is possible when the mass of some glueball state is close to the mass of the emitted meson $M_1$. This amplification may thus occur in the $0^{++}$ channel for $M_1$ around 1.5 GeV, or in other channels for $M_1 \gtrsim 2$ GeV [22]. A similar mechanism may take place in hadron-hadron scattering with creation of the $c\bar{c}$ states in the mass range 3 - 4 GeV, where also glueball states are predicted in lattice and in analytic calculations [22].

5 Conclusion

In the present paper the first qualitative step was done aimed at the systematic derivation of strong decay amplitudes from the QCD Lagrangian. We
have considered three decay mechanisms, where nonperturbative contribu-
tion is very important.

The first one, the CDM, was already successfully checked in heavy-light meson decays [15]. The next step would be to apply CDM to the pionic and kaonic decays of light-light mesons, and investigate double and more-pion decays, which are given by the nonlinear Lagrangian \( \mathcal{L} \). The second mechanism, the SBM, turns out to be mostly \( ^3P_0 \) (additional terms due to the smaller nonscalar component in the quark Green’s function \( S \) considered in [14] have been neglected above) and its amplitude is close to the phenomenological fits [4]. The third type of mechanisms with hybrid or glueball in the intermediate state is the nonperturbative background generalization of the original purely perturbative mechanisms of the \(^3S_1\) type [23] and of two-gluon exchange [21], respectively. It is argued that a strong enhancement of decay amplitudes is possible when the corresponding levels of a hybrid or a glueball are close to the mass of the original or the final meson respectively. The paper does not contain quantitative predictions, which are planned for the future, but rather is concentrated on the general discussion of possible decay mechanisms as they emerge from the basic QCD Lagrangian.

One of immediate extensions of the present result is the inclusion of the baryon decays, where all three mechanisms discussed above are present in the same form with the only replacement of the simple mesonic string by the \( Y \)-shaped baryonic string.

Finally, one should have in mind that all three mechanisms discussed above enter the decay amplitude additively and therefore one can expect in general some interference effects, which make the analysis of data more complicated and perhaps more interesting.

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