Widths of tetraquarks with open charm

S.M. Gerasyuta 1,2 and V.I. Kochkin 1

1 Department of Theoretical Physics, St. Petersburg State University, 198904, St. Petersburg, Russia
2 Department of Physics, LTA, 194021, St. Petersburg, Russia

E-mail: gerasyuta@SG6488.spb.edu

Abstract

In the framework of coupled-channel formalism the relativistic four-quark equations are found. The dynamical mixing of the meson-meson states with the four-quark states is considered. The four-quark amplitudes of the tetraquarks with open charm, including u, d, s, c quarks, are constructed. The poles of these amplitudes determine the masses and widths of tetraquarks.

Keywords: Tetraquarks; coupled-channel formalism.

PACS number: 11.55.Fv, 12.39.Ki, 12.39.Mk, 12.40.Yx.

I. Introduction.

The discovery by the Belle Collaboration [1] of the very narrow X(3872) resonance has triggered the interest in the charmonium-like states, both experimentally and theoretically. The observation of X(3872) has been confirmed by CDF [2], D0 [3] and BaBar Collaboration [4]. Belle Collaboration observed the X(3940) in double-charmonium production in the reaction $e^+e^- \rightarrow J/\psi + X$ [5]. The state, designated as X(4160), was reported by the Belle Collaboration in Ref. 6. The fact that the newly found states do not fit quark model calculations [7] has opened the discussion about the structure of such states. Maiani et al. advocate a tetraquark explanation for the X(3872) [8, 9]. On the other hand, the mass of X(3872) is very close to the threshold of $D^*D$ and, therefore, it can be interpreted as molecular state [10 – 15]. Several review papers, as for example [16, 17], discuss the difficulty of interpreting these resonances as charmonium states.

In the present paper the relativistic four-quark equations are found in the framework of coupled-channel formalism. The dynamical mixing between the meson-meson states and the four-quark states is considered [18 – 20]. Taking the X(3872) and X(3940) as input [21] we predict the masses and widths of S-wave tetraquarks with open charm (Table I).

II. Four-Quark Amplitudes for the Tetraquarks with Open Charm.

We derive the relativistic four-quark equations in the framework of the dispersion relations technique.

The correct equations for the amplitude are obtained by taking into account all possible subamplitudes. It corresponds to the division of complete system into subsys-
tems with the smaller number of particles. Then one should represent a four-particle amplitude as a sum of six subamplitudes:

\[ A = A_{12} + A_{13} + A_{14} + A_{23} + A_{24} + A_{34}. \]  

(1)

This defines the division of the diagrams into groups according to the certain pair interaction of particles. The total amplitude can be represented graphically as a sum of diagrams.

We need to consider only one group of diagrams and the amplitude corresponding, for example, \( A_{12} \). The relativistic generalization of the Faddeev-Yakubovsky approach [22, 23] for the tetraquark is obtained. We shall construct the four-quark amplitude of \( c\bar{u}u\bar{u} \) tetraquark in which the quark amplitudes with quantum numbers of \( 0^{--} \) and \( 1^{--} \) are included. The set of diagrams associated with the amplitude \( A_{12} \) can further be broken down into four groups corresponding to subamplitudes: \( A_1(s, s_{12}, s_34), A_2(s, s_{23}, s_{14}), A_3(s, s_{23}, s_{123}), A_4(s, s_{14}, s_{124}) \), if we consider the tetraquark with the spin-parity \( J^{pc} = 0^{++} (c\bar{u}u\bar{u}) \).

Here \( s_{ik} \) is the two-particle subenergy squared, \( s_{ijk} \) corresponds to the energy squared of particles \( i, j, k \) and \( s \) is the system total energy squared.

In order to represent the subamplitudes \( A_1(s, s_{12}, s_34), A_2(s, s_{23}, s_{14}), A_3(s, s_{23}, s_{123}), A_4(s, s_{14}, s_{124}) \) in the form of dispersion relations it is necessary to define the amplitudes of quark-antiquark interaction \( a_n(s_{ik}) \). The pair quarks amplitudes \( q\bar{q} \rightarrow q\bar{q} \) are calculated in the framework of the dispersion \( N/D \) method with the input four-fermion interaction [24 – 26] and the quantum numbers of the gluon [27]. The regularization of the dispersion integral for the \( D \)-function is carried out with the cutoff parameter \( \Lambda \). The four-quark interaction is considered as an input [27]:

\[ g_V (\bar{q} \gamma \gamma \gamma q)^2 + g_V^{ss} (\bar{q} \gamma \gamma \gamma q) (\bar{s} \gamma \gamma \gamma s) + g_V^{ss} (\bar{s} \gamma \gamma \gamma s)^2. \]  

(2)

Here \( I_f \) is the unity matrix in the flavor space \((u, d)\). \( \lambda \) are the color Gell-Mann matrices. Dimensional constants of the four-fermion interaction \( g_V, g_V^{ss} \) and \( g_V^{ss} \) are parameters of the model. At \( g_V = g_V^{ss} = g_V^{ss} \) the flavor SU(3)_f symmetry occurs. The strange quark violates the flavor SU(3)_f symmetry. In order to avoid an additional violation parameters, we introduce the scale shift of the dimensional parameters [27]:

\[ g = \frac{m^2}{\pi^2} g_V = \left( \frac{m + m_s}{4\pi^2} \right) g_V^{ss} = \frac{m_s^2}{\pi^2} g_V^{ss}. \]  

(3)

\[ \Lambda = \frac{4\Lambda (ik)}{(m_i + m_k)^2}. \]  

(4)

Here \( m_i \) and \( m_k \) are the quark masses in the intermediate state of the quark loop. Dimensionless parameters \( g \) and \( \Lambda \) are supposed to be constants which are independent of the quark interaction type. The applicability of Eq. (2) is verified by the success of De Rujula-Georgi-Glashow quark model [28], where only the short-range part of Breit potential connected with the gluon exchange is responsible for the mass splitting in hadron multiplets.

We use the results of our relativistic quark model [27] and write down the pair quarks amplitude in the form:

\[ a_n(s_{ik}) = \frac{G_n^2(s_{ik})}{1 - B_n(s_{ik})}, \]  

(5)
Here \( G_n(s_{ik}) \) are the quark-antiquark vertex functions. The vertex functions are determined by the contribution of the crossing channels. The vertex functions satisfy the Fierz relations. All of these vertex functions are generated from \( g_V \), \( g_V^{(s)} \) and \( g_V^{(ss)} \). \( B_n(s_{ik}) \), \( \rho_n(s_{ik}) \) are the Chew-Mandelstam functions with cutoff \( \Lambda \) and the phase spaces, respectively.

Here \( n = 1 \) determines a \( q\bar{q} \)-pairs with \( J^{pc} = 0^{-+} \) in the \( 1_c \) color state, \( n = 2 \) corresponds to a \( q\bar{q} \)-pairs with \( J^{pc} = 1^{--} \) in the \( 1_c \) color state, and \( n = 3 \) defines the \( q\bar{q} \)-pairs corresponding to tetraquarks with quantum numbers: \( J^{pc} = 0^{++}, 1^{++}, 2^{++} \).

In the case in question, the interacting quarks do not produce a bound state; therefore, the integration in Eqs. (7) – (10) is carried out from the threshold \((m_i + m_k)^2\) to the cutoff \( \Lambda(i,k) \). The coupled integral equation systems (the tetraquark state with \( n = 3 \) and \( J^{pc} = 0^{++} \) for the \( c\bar{c}u\bar{u} \) can be described as:

\[
A_1(s, s_{12}, s_{34}) = \frac{\lambda_1 B_2(s_{12}) B_2(s_{34})}{[1 - B_2(s_{12})][1 - B_2(s_{34})]} + 2 \hat{J}_2(s_{12}, s_{34}, 2, 2) A_3(s, s'_{23}, s'_{123})
\]
\[
+ 2 \hat{J}_2(s_{12}, s_{34}, 2, 2) A_1(s, s'_{14}, s'_{124}),
\]
\[
A_2(s, s_{23}, s_{14}) = \frac{\lambda_2 B_1(s_{23}) B_1(s_{14})}{[1 - B_1(s_{23})][1 - B_1(s_{14})]} + 2 \hat{J}_2(s_{23}, s_{14}, 1, 1) A_3(s, s'_{34}, s'_{234})
\]
\[
+ 2 \hat{J}_2(s_{23}, s_{14}, 1, 1) A_1(s, s'_{12}, s'_{123}),
\]
\[
A_3(s, s_{23}, s_{123}) = \frac{\lambda_3 B_3(s_{23})}{[1 - B_3(s_{23})]} + 2 \hat{J}_3(s_{23}, 3) A_1(s, s'_{12}, s'_{34}) + \hat{J}_3(s_{23}, 3) A_2(s, s'_{12}, s'_{34})
\]
\[
+ \hat{J}_1(s_{23}, 3) A_4(s, s'_{34}, s'_{234}) + \hat{J}_1(s_{23}, 3) A_3(s, s'_{12}, s'_{123}),
\]
\[
A_4(s, s_{14}, s_{124}) = \frac{\lambda_4 B_4(s_{14})}{[1 - B_4(s_{14})]} + 2 \hat{J}_4(s_{14}, 3) A_1(s, s'_{13}, s'_{24}) + 2 \hat{J}_4(s_{14}, 3) A_2(s, s'_{13}, s'_{24})
\]
\[
+ 2 \hat{J}_1(s_{14}, 3) A_3(s, s'_{14}, s'_{34}) + 2 \hat{J}_1(s_{14}, 3) A_4(s, s'_{14}, s'_{134}),
\]

where \( \lambda_i, i = 1, 2, 3, 4 \) are the current constants. They do not affect the mass spectrum of tetraquarks. We introduce the integral operators:

\[
\hat{J}_1(s_{12}, l) = \frac{G_l(s_{12})}{[1 - B_l(s_{12})]} \int_{(m_1 + m_2)^2}^{[\Lambda]^2} \frac{ds'_{12} G_l(s'_{12}) \rho_l(s'_{12})}{\pi} \frac{1}{s'_{12} - s_{12}} \int_1^{+1} \frac{dz_1}{2};
\]
\[
\hat{J}_2(s_{12}, s_{34}, l, p) = \frac{G_l(s_{12}) G_p(s_{34})}{[1 - B_l(s_{12})][1 - B_p(s_{34})]} \int_{(m_1 + m_2)^2}^{[\Lambda]^2} \frac{ds'_{12} G_l(s'_{12}) \rho_l(s'_{12})}{\pi} \frac{1}{s'_{12} - s_{12}}
\]
the final state, is taken in the c.m. of particles 1 and 2. In Eq. (10): 

\begin{equation}
\frac{(m_3 + m_4)^2 \Lambda}{\pi} \int \frac{ds'_{34}}{s'_{34} - s_{34}} + \frac{1}{2} \int \frac{dz_3}{\sqrt{1 - z_3^2}} \int \frac{dz_4}{\sqrt{1 - z_4^2}} ,
\end{equation}

\begin{equation}
\hat{J}_3(s_{12}, l) = \frac{G_i(s_{12}, \tilde{\Lambda})}{[1 - B_i(s_{12}, \tilde{\Lambda})]} \frac{1}{4\pi} \frac{(m_1 + m_2)^2 \Lambda}{\pi} \int \frac{ds'_{12}}{s'_{12} - s_{12}} \rho_l(s'_{12}) \rho_l(s'_{12})
\end{equation}

\begin{equation}
\times \int \frac{+1}{z_2^2} \frac{dz_1}{2} \int \frac{dz}{1 - z^2} \int \frac{dz_2}{\sqrt{1 - z^2 - z_2^2 - 2z_2z_1z_2^2}} .
\end{equation}

here \( l, p \) are equal to 1 – 3.

In Eqs. (11) and (13) \( z_1 \) is the cosine of the angle between the relative momentum of the particles 1 and 2 in the intermediate state and the momentum of the particle 3 in the final state, taken in the c.m. of particles 1 and 2. In Eq. (13) \( z \) is the cosine of the angle between the momenta of particles 3 and 4 in the final state, taken in the c.m. of particles 1 and 2. \( z_2 \) is the cosine of the angle between the relative momentum of particles 1 and 2 in the intermediate state and the momentum of the particle 4 in the final state, is taken in the c.m. of particles 1 and 2. In Eq. (10): \( z_3 \) is the cosine of the angle between relative momentum of particles 1 and 2 in the intermediate state and the relative momentum of particles 3 and 4 in the intermediate state, taken in the c.m. of particles 1 and 2. \( z_4 \) is the cosine of the angle between the relative momentum of the particles 3 and 4 in the intermediate state and that of the momentum of the particle 1 in the intermediate state, taken in the c.m. of particles 3, 4.

We can pass from the integration over the cosines of the angles to the integration over the subenergies [29].

Let us extract two-particle singularities in the amplitudes \( A_1(s, s_{12}, s_{34}) \), \( A_2(s, s_{23}, s_{14}) \), \( A_3(s, s_{23}, s_{123}) \) and \( A_4(s, s_{14}, s_{124}) \):

\begin{equation}
A_1(s, s_{ik}, s_{lm}) = \frac{\alpha_1(s, s_{ik}, s_{lm})B_2(s_{ik})B_2(s_{lm})}{[1 - B_2(s_{ik})][1 - B_2(s_{lm})]} ,
\end{equation}

\begin{equation}
A_2(s, s_{ik}, s_{lm}) = \frac{\alpha_2(s, s_{ik}, s_{lm})B_1(s_{ik})B_1(s_{lm})}{[1 - B_1(s_{ik})][1 - B_1(s_{lm})]} ,
\end{equation}

\begin{equation}
A_j(s, s_{ik}, s_{ikl}) = \frac{\alpha_j(s, s_{ik}, s_{ikl})B_3(s_{ik})}{1 - B_3(s_{ik})} , \quad j = 3 - 4 .
\end{equation}

We do not extract three-particles singularities, because they are weaker than two-particle singularities.

We used the classification of singularities, which was proposed in paper [30]. The construction of the approximate solution of Eqs. (7) – (10) is based on the extraction of the leading singularities of the amplitudes. The main singularities in \( s_{ik} \approx (m_i + m_k)^2 \) are from pair rescattering of the particles \( i \) and \( k \). First of all there are threshold square-root singularities. Also possible are pole singularities which correspond to the bound states. The amplitudes apart from two-particle singularities have triangular singularities and the singularities defining the interactions of four particles. Such classification allows us to search the corresponding solution of Eqs. (7) – (10) by taking into account some definite number of leading singularities and neglecting all the weaker ones. We
consider the approximation which defines two-particle, triangle and four-particle singularities. The functions \( \alpha_1(s, s_{12}, s_{34}), \alpha_2(s, s_{23}, s_{14}), \alpha_3(s, s_{23}, s_{123}) \) and \( \alpha_4(s, s_{14}, s_{124}) \) are the smooth functions of \( s_k, s_{ik}, s \) as compared with the singular part of the amplitude, hence they can be expanded in a series in the singularity point and only the first term of this series should be employed further. Using this classification, one defines the reduced amplitudes \( \alpha_1, \alpha_2, \alpha_3, \alpha_4 \) as well as the \( B \)-functions in the middle point of physical region of Dalitz-plot at the point \( s_0 \):

\[
s_{0}^{ik} = 0.25(m_i + m_k)^2 s_0,
\]

\[
s_{123} = 0.25 s_0 \sum_{i,k=1}^{3} (m_i + m_k)^2 - \sum_{i=1}^{3} m_i^2, \quad s_0 = \frac{s + 2 \sum_{i=1}^{4} m_i^2}{0.25 \sum_{i,k=1}^{4} (m_i + m_k)^2}.
\]

Such a choice of point \( s_0 \) allows us to replace integral Eqs. (7) – (10) by the algebraic equations (18) – (21) respectively:

\[
\alpha_1 = \lambda_1 + 2\alpha_3 JB_1(2, 2, 3) + 2\alpha_4 JB_2(2, 2, 3),
\]

\[
\alpha_2 = \lambda_2 + 2\alpha_3 JB_3(1, 1, 3) + 2\alpha_4 JB_4(1, 1, 3),
\]

\[
\alpha_3 = \lambda_3 + 2\alpha_1 JC_1(3, 2, 2) + 2\alpha_2 JC_2(3, 1, 1) + \alpha_4 JA_1(3) + \alpha_3 JA_2(3),
\]

\[
\alpha_4 = \lambda_4 + 2\alpha_1 JC_3(3, 2, 2) + 2\alpha_2 JC_4(3, 1, 1) + \alpha_3 JA_3(3) + \alpha_4 JA_4(3).
\]

We use the functions \( JA_1(l), JB_1(l, p, r), JC_1(l, p, r) \) \( (l, p, r = 1 - 3) \), which are determined by the various \( s_0^{ik} \) (Eq. 17). These functions are similar to the functions:

\[
JA_4(l) = \frac{G_l^2(s_0^{12}) B_l^2(s_0^{23})}{B_l(s_0^{12})} \frac{(m_1 + m_2)^2 \Lambda}{(m_1 + m_2)^2} \int \frac{ds_{12}'}{\pi} \rho_l(s_{12}') \int_{-1}^{+1} \frac{dz_1}{2} \frac{1}{1 - B_l(s_{23}')},
\]

\[
JB_1(l, p, r) = \frac{G_l^2(s_0^{12}) G_p^2(s_0^{34}) B_r(s_0^{23})}{B_l(s_0^{12}) B_p(s_0^{34})} \frac{(m_1 + m_2)^2 \Lambda}{(m_1 + m_2)^2} \int \frac{ds_{12}'}{\pi} \rho_l(s_{12}') \int_{-1}^{+1} \frac{dz_3}{2} \frac{dz_4}{2} \frac{1}{1 - B_r(s_{23}')},
\]

\[
JC_3(l, p, r) = \frac{G_l^2(s_0^{12}, \Lambda) B_p(s_0^{23}) B_r(s_0^{23}),}{1 - B_l(s_0^{12}, \Lambda)} \frac{1}{4\pi} \frac{(m_1 + m_2)^2 \Lambda}{(m_1 + m_2)^2} \int \frac{ds_{12}'}{\pi} \rho_l(s_{12}') \int_{-1}^{+1} \frac{dz_1}{2} \frac{dz_2}{2} \frac{dz_3}{2} \frac{1}{1 - z_1^2 - z_2^2 - z_3^2 + 2z_1z_2z_3}.
\]
\[ \frac{1}{[1 - B_p(s_{23}')]^2[1 - B_r(s_{14}')]}, \]  
\[ \tilde{\Lambda}(ik) = \begin{cases} 
\Lambda(ik), & \text{if } \Lambda(ik) \leq (\sqrt{s_{123}} + m_3)^2 \\
\left(\sqrt{s_{123}} + m_3\right)^2, & \text{if } \Lambda(ik) > (\sqrt{s_{123}} + m_3)^2 
\end{cases} \]  

The other choices of point \( s_0 \) do not change essentially the contributions of \( \alpha_1, \alpha_2, \alpha_3 \) and \( \alpha_4 \), therefore we omit the indices \( s_0^{ik} \). Since the vertex functions depend only slightly on energy it is possible to treat them as constants in our approximation.

The solutions of the system of equations are considered as:

\[ \alpha_i(s) = F_i(s, \lambda_i)/D(s), \]  

where zeros of \( D(s) \) determinants define the masses of bound states of tetraquarks. \( F_i(s, \lambda_i) \) determine the contributions of subamplitudes to the tetraquark amplitude.

### III. Calculation results.

Our calculations do not include the new parameters. We use the cutoff \( \Lambda = 10.0 \) and the gluon coupling constant \( g = 0.794 \), which are determined by fixing the tetraquark masses for the states with the hidden charm \( J^{pc} = 1^{++} X(3872) \) and \( J^{pc} = 2^{++} X(3940) \) [21]. The quark masses of model \( m_{u,d} = 385 \text{ MeV} \) and \( m_s = 510 \text{ MeV} \) coincide with the ordinary meson model ones [27]. In order to fix anyhow \( m_c = 1586 \text{ MeV} \), we use the tetraquark mass for the \( J^{pc} = 2^{++} X(3940) \). The masses and widths of meson-meson states with the spin-parity \( J^{pc} = 0^{++}, 1^{++}, 2^{++} \) are given in Table I. In our paper we predicted the scalar tetraquark with the mass \( M = 2610 \text{ MeV} \) and the width \( \Gamma_0^{++} = 180 \text{ MeV} \) (channel \( D^0\eta \)). We calculated the scalar tetraquark with the mass \( M = 2691 \text{ MeV} \) and the width \( \Gamma_0^{++} = 110 \text{ MeV} \) (channels \( D^+_\eta \) and \( D^0K^+ \)). The other scalar tetraquark is predicted as \( M = 2805 \text{ MeV} \) and the width \( \Gamma_0^{++} = 80 \text{ MeV} \) (channel \( D^+_sK^- \)). The tetraquarks with the spin-parity \( J^{pc} = 1^{++}, 2^{++} \) (Table I) have only the weak decays.

The functions \( F_i(s, \lambda_i) \) (Eq. (26)) allow us to obtain the overlap factors \( f \) for the tetraquarks. We calculated the overlap factors \( f \) and the phase spaces \( \rho \) for the reactions \( X \rightarrow M_1 M_2 \) (Table II). The widths of the scalar tetraquarks with open charm are obtained (Table I). We considered the formula \( \Gamma \sim f^2 \times \rho \) [31], there \( \rho \) is the phase space. The widths of the tetraquarks are fitted by the fixing width \( \Gamma_2^{++} = (39 \pm 26) \text{ MeV} \) [21] for the tetraquark \( X(3940) \) (\( c\bar{c}u\bar{u} \)) with the spin-parity \( J^{pc} = 2^{++} \).

In the open charm sector the scalar tetraquarks have relatively small width \( \sim 100 - 200 \text{ MeV} \), so in principle these exotic states could be observed. The low-lying tetraquarks with the open charm were calculated in other works [32, 33]. These states appear as narrow states. In our model the tetraquarks with the open charm and the spin-parity \( J^{pc} = 1^{++}, 2^{++} \) can decay only in the weak channels.

### Acknowledgments.

The work was carried with the support of the Russian Ministry of Education (grant 2.1.1.68.26).
Table I. Masses and widths of tetraquark with open charm.

Parameters of model [21]: quark masses $m_{u,d} = 385 \text{ MeV}$, $m_s = 510 \text{ MeV}$ and $m_c = 1586 \text{ MeV}$; cutoff parameter $\Lambda = 10.0$, gluon coupling constant $g = 0.794$.

| Tetraquark   | $J^{pc}$ | Mass (MeV) | Width (MeV) | $J^{pc}$ | Mass (MeV) | $J^{pc}$ | Mass (MeV) |
|--------------|----------|------------|-------------|----------|------------|----------|------------|
| $(c\bar{u})(u\bar{u})$ | 0$^{++}$ | 2610       | 180         | 1$^{++}$ | 2672       | 2$^{++}$ | 2736       |
| $(c\bar{s})(u\bar{u})$ | 0$^{++}$ | 2691       | 110         | 1$^{++}$ | 2770       | 2$^{++}$ | 2851       |
| $(c\bar{u})(s\bar{s})$ | 0$^{++}$ | 2805       | 80          | 1$^{++}$ | 2890       | 2$^{++}$ | 2975       |

Table II. Overlap factors $f$ and phase spaces $\rho$ of tetraquarks with open charm.

| Tetraquark (channels) | $J^{pc}$ | $f$  | $\rho$ |
|-----------------------|----------|------|--------|
| $(c\bar{u})(u\bar{u})$ | $D^0\eta$ | 0$^{++}$ | 0.396  | 0.325  |
| $(c\bar{s})(u\bar{u})$ | $D_s^+\eta$ | 0$^{++}$ | 0.246  | 0.300  |
| $(c\bar{u})(u\bar{s})$ | $D^0K^+$ | 0$^{++}$ | 0.183  | 0.414  |
| $(c\bar{s})(s\bar{s})$ | $D_s^0\eta_s$ | 0$^{++}$ | 0.192  | –      |
| $(c\bar{s})(s\bar{u})$ | $D_s^+K^-$ | 0$^{++}$ | 0.237  | 0.407  |
References.

1. S.K. Choi et al. (Belle Collaboration), Phys. Rev. Lett. 91, 262001 (2003).
2. D. Acosta et al. (CDF Collaboration), Phys. Rev. Lett. 93, 072001 (2004).
3. V.M. Abazov et al. (D0 Collaboration), Phys. Rev. Lett. 93, 162002 (2004).
4. B. Aubert et al. (BaBar Collaboration), Phys. Rev. D71, 071103 (2005).
5. K. Abe et al. (Belle Collaboration), Phys. Rev. Lett. 98, 082001 (2007).
6. I. Adachi et al. (Belle Collaboration), arXiv: 0708.3812 [hep-ex].
7. S. Godfrey and N. Isgur, Phys. Rev. D32, 189 (1985).
8. L. Maiani, F. Piccinini, A.D. Polosa and V. Riequer, Phys. Rev. D71, 014028 (2005).
9. L. Maiani, A.D. Polosa and V. Riequer, Phys. Rev. Lett. 99, 182003 (2007).
10. N.A. Tornqvist, Phys. Lett. B590, 209 (2004).
11. F.E. Close and P.R. Page, Phys. Lett. B628, 215 (2005).
12. E.S. Swanson, Int. J. Mod. Phys. A21, 733 (2006).
13. T. Barnes, Int. J. Mod. Phys. A21, 5583 (2006).
14. S.H. Lee, K. Morita and M. Nielsen, arXiv: 0808.3168 [hep-ph].
15. Y. Dong, A. Faessler, T. Gutsche and V.E. Lyubovitskij, arXiv: 0802.3610 [hep-ph].
16. E.S. Swanson, Phys. Rept. 429, 243 (2006).
17. S. Godfrey and S.L. Olsen, arXiv: 0801.3867 [hep-ph].
18. S.M. Gerasyuta and V.I. Kochkin, Z. Phys. C74, 325 (1997).
19. S.M. Gerasyuta and V.I. Kochkin, Nuovo Cim. A110, 1313 (1997).
20. S.M. Gerasyuta and V.I. Kochkin, arXiv: 0804.4567 [hep-ph].
21. S.M. Gerasyuta and V.I. Kochkin, arXiv: 0809.1758 [hep-ph].
22. O.A. Yakubovsky, Sov. J. Nucl. Phys. 5, 1312 (1967).
23. S.P. Merkuriev and L.D. Faddeev, Quantum scattering theory for system of few particles (Nauka, Moscow 1985) p. 398.
24. Y. Nambu and G. Jona-Lasinio, Phys. Rev. 122, 365 (1961): ibid. 124, 246 (1961).
25. T. Appelqvist and J.D. Bjorken, Phys. Rev. D4, 3726 (1971).
26. C.C. Chiang, C.B. Chiu, E.C.G. Sudarshan and X. Tata, Phys. Rev. D25, 1136 (1982).
27. V.V. Anisovich, S.M. Gerasyuta and A.V. Sarantsev, Int. J. Mod. Phys. A6, 625 (1991).
28. A.De Rujula, H.Georgi and S.L.Glashow, Phys. Rev. D12, 147 (1975).
29. S.M. Gerasyuta and V.I. Kochkin, Yad. Fiz. 59, 512 (1996) [Phys. At. Nucl. 59, 484 (1996)].
30. V.V. Anisovich and A.A. Anselm, Usp. Phys. Nauk. 88, 287 (1966).
31. J.J. Dudek and F.E. Close, Phys. Lett. B583, 278 (2004).
32. E.E. Kolomeitsev and M.F.M. Lutz, Phys. Lett. B582, 39 (2004).
33. F.K. Guo, P.N. Shen and H.C. Chiang, Phys. Lett. B647, 133 (2007).