Chiral Symmetry Breaking and Scalar Confinement

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We address the old difficulty in accommodating the scalar quark-antiquark confining potential together with chiral symmetry breaking. We develop a quark confining potential inspired in the QCD scalar flux tube. The coupling to quarks consists in a double vector vertex. We study the Dirac and spin structure of this potential. In the limit of massless quarks the quark vertex is vector. Nevertheless symmetry breaking generates a new scalar quark vertex. In the heavy quark limit the coupling is mostly scalar. We solve the mass gap equation and find that this potential produces spontaneous chiral symmetry breaking for light quarks. The quantitative results of this model are encouraging.

I. INTRODUCTION

Spontaneous Chiral Symmetry Breaking (S\(\chi\)SB) is accepted to occur in low energy hadronic physics. Another important feature of hadronic physics, suggested by the spectroscopy of hadrons, by lattice simulations and by models of confinement is scalar confinement. However S\(\chi\)SB and scalar confinement are apparently conflicting, since the first requires a chiral invariant coupling to the quarks, like the vector coupling of QCD. We address a recent quest of Bjorken \cite{1}, “how are the many disparate methods of describing hadrons which are now in use related to each other and to first principles of QCD?”. Although the vector confinement of quarks is not yet ruled out \cite{2,3,4,5}, here we try to solve this old chiral symmetry versus scalar confinement conflict of hadronic physics, which remained open for many years.

In this paper we explore scalar confinement from the perspective of chiral symmetry breaking. In Section II we motivate the importance of both S\(\chi\)SB and scalar confinement. We show these features of hadronic physics to have some subtle weaknesses that we capitalize to construct a model. The potential used in our Quark Model (QM) is defined in Section III. The self consistent mass gap equation for the quarks is derived in Section IV. In Section V we solve numerically the mass gap equation and calculate the quark condensate. Finally, in Section VI we present some conclusions.

II. MATCHING CHIRAL SYMMETRY BREAKING WITH SCALAR CONFINEMENT

The QCD lagrangian is chiral invariant in the limit of vanishing quark masses. Nambu and Jona-Lasinio showed that including chiral symmetry in fermionic systems provides a natural explanation for the small pion mass, which is much lighter than all the other isovector hadrons. Because of this crucial fact the mechanism of S\(\chi\)SB is accepted to occur in low energy hadronic physics for the light flavors \(u, d \) and \(s\), where \(m_u, m_d \ll m_s \ll \Lambda_{QCD} < M_N/3\). Similarly to the vector Ward identities in gauge symmetry, the axial Ward identities constitute a powerful tool of chiral symmetry. The techniques of current algebra led to beautifully correct theorems, the Partially Conserved Axial Current (PCAC) theorems. The different variants of QMs are widely used as simplifications of QCD. They are convenient to study quark bound states and hadron scattering. Recently \cite{6,5} we have shown these beautiful PCAC theorems, like the Weinberg theorem for \(\pi - \pi\) scattering, to be reproduced by QMs with S\(\chi\)SB. Another important benefit of having S\(\chi\)SB in the QM is the reduced number of parameters. The mass gap equation generates a dynamical constituent quark mass, which is no longer an independent parameter, even for quarks with a vanishing current mass. The quark-quark, quark-antiquark, antiquark-antiquark potentials, and the quark-antiquark annihilation and creation interactions are all originated in the same chiral invariant Bethe-Salpeter kernel. Therefore any QM for light quarks should comply with the S\(\chi\)SB. Moreover the microscopic coupling of the quark to the confining interaction should include the vector coupling which is present in the quark-gluon vertex of QCD. However there is some evidence that a vector quark-quark potential is not sufficient to provide the expected scale of S\(\chi\)SB of the order of 200 – 300 MeV which is present both in the constituent quark mass and in the quark condensate. It was realized by Adler \cite{7} that a linear confinement with vector couplings was not sufficient to provide the correct quark condensate. Moreover the gluon propagator extracted from the lattice, when used in a one gluon exchange truncation of the quark mass gap equation, is not able to provide the expected quark condensate \cite{8}. This also happens with the gluon propagator extracted from the solution of truncated Schwinger-Dyson equations of QCD \cite{15}. Importantly, these gluon propagators exhibit a non-perturbative mass. This mass should produce a Meissner effect in Yang-Mills fields, and this is expected to produce a confining string for the quarks. Here we will estimate the effect of the confining string on the quark condensate.

On the other hand the confining potential for con-
stinent quarks is probably scalar. We can learn much by comparing simply the spectrum of the hydrogen atom with the masses of all hadron families. In a perturbative QCD scenario, the hadron spectroscopy would be dominated by the one gluon exchange, which is qualitatively similar to the one photon exchange interaction that explains in detail atomic physics. It turns out that all the hadronic families, say of mesons or baryons, with light or heavy flavors, show similar differences with the hydrogen spectrum. It is remarkable that the Spin-Orbit potential (also called fine interaction in atomic physics) turns out to be suppressed in hadronic spectra since it is smaller than the Spin-Spin potential (hyperfine interaction). This constitutes an evidence of non-perturbative QCD. Another evidence of non-perturbative QCD is present in the angular and radial excitations of hadrons, which fit linear trajectories in Regge plots, and suggest a long range, probably linear, confining potential for the quarks. This led Henriques, Kellett and Moorhouse [14] to develop a QM where a short range vector potential and a long range scalar potential partly cancel the Spin-Orbit contribution. The short range potential is Coulomb-like (inspired in the one gluon exchange) and its quark vertex has a vector coupling structure $\langle \bar{\psi} \gamma^\mu \gamma^5 \psi \rangle$. The long range potential has scalar coupling $\bar{\psi} \psi$ and is a linear potential [14, 11, 12, 13].

The same scalar confinement picture is extracted from lattice simulations. In quenched lattices which simulate the quark-antiquark potential in the heavy quark limit, the pattern of spin-spin, tensor, and spin-orbit interactions is compatible with a scalar confinement [14, 15, 16]. Moreover the presently favored confinement picture in the literature is the flux tube, or string picture, with tension $\sigma \sim 200$ MeV/Fm. Quantum mechanics suggests that a thin string, in its ground-state, should be a scalar object [17]. Only higher energy excitations of the string would have angular momentum. This was capitalized by Isgur and Paton in the flux tube model [18].

In this paper we will assume the quarks are coupled to a scalar object which provides a linear confinement. On top of that we want this coupling to use, as a microscopic building block, the vector gluon-quark which is present in QCD. Notice that to get a Lorentz scalar coupling a simple one gluon vertex is not enough. The coupling needs at least two vertices. The simplest way to achieve this is to have the string emitting two effective gluons that couple to the same quark line. This double vertex, coupling the string to a quark line via two intermediate gluon propagators, is presented in Fig. 1. Effectively this double vertex is similar to the vertices that couple a quark to a gluon ladder in the soft pomeron models [19, 20]. It is also related to the light quark vertex in the heavy-light quark bound states studied in the local gauge coordinate [21, 22, 23]. Such a double vertex was also introduced in the coupling of short strings to quarks [24]. In the cumulant expansion formalism in terms of gluon correlators [21, 22], this means the quark-antiquark coupling will be dominated by four gluon correlators.

![Figure 1](image)

**FIG. 1:** The coupling of a quark to a string with a double gluon vertex

### III. THE DOUBLE VERTEX

**NON-PERTURBATIVE CONFINING INTERACTION**

In this section we construct the simplest possible coupling, that simulates that of a quark line with a scalar string using two vector couplings as building blocks. The most general coupling of this kind is presented in Fig. 1 and it is simply read,

$$\int \frac{d^4 k}{(2\pi)^4} (\Gamma V^a) S(k) (\Gamma V^b) G^{bc}(p-k)(\Gamma_S W^{cde}) G^{ad}(q-k) ,$$

where $\Gamma$ is the dirac structure of the fermion-gluon interaction and $V^a$ the usual color interaction $\lambda^a / 2$. We denote the quark propagator by $S(k)$ and the gluon propagator by $G^{ab}(k)$. Finally $\Gamma_S W^{cde}$ represents the coupling of the gluon pair to the string.

To get the coupling of the string to a light quark, we follow the coupling obtained in the heavy-light quark system, computed in the local coordinate gauge [22]. This model interpolates between the heavy-light meson in the local gauge and the effective QM. So, for the Dirac structure of the quark-gluon sub-vertices $\Gamma$ we have $\gamma^0$ matrices, which is also compatible with the Coulomb gauge.

In the color sector, as already stated, the coupling of the same sub-vertices has a $\lambda^a / 2$ structure, where $\lambda^a$ are the Gell-Mann matrices. The remaining sub-vertex includes the coupling of two color octets, see Fig. 1. The string is also a colored object; it contains a flux of color-electric field. For a scalar coupling, which is symmetric, we use the symmetric structure function $d^{abc}$ defined by

$$\{\lambda^a, \lambda^b\} = \frac{4}{3} \delta^{ab} + 2 d^{abc} \lambda^c .$$

This will result in a color contribution for the effective vertex of

$$d^{abc} \frac{\lambda^b}{2} \cdot \frac{\lambda^c}{2} = C \frac{\lambda^a}{2} , \quad C = \frac{5}{6} .$$

In QMs the string usually couples to the quark line with a $\lambda^a / 2$. In our case it couples with two $\lambda^a / 2$, one for
The energy loop integral can be easily calculated,\[ \int \frac{dk^0}{2\pi} S(k) \beta = (s_k \beta + c_k \mathbf{k} \cdot \mathbf{\alpha}) \beta. \]

Finally, using the Dirac delta function for the remaining integrals over the three momentum $k$, and summing in color indices, we get the following effective vertex,

\[ V_{\text{eff}} = C \frac{\lambda^c}{2} (s_k - c_k \mathbf{k} \cdot \gamma) \bigg|_{k = \frac{\pi}{2}}. \]

In the remainder of the paper we will always assume $k = (p + q)/2$.

Eq. (10) shows that the double vertex actually solves the problem of matching chiral symmetry breaking and scalar confinement. In the chiral limit of a vanishing quark mass, the effective vertex $V_{\text{eff}} \rightarrow -C \lambda^c/2 \mathbf{k} \cdot \gamma$ is proportional to the $\gamma^c$ and is therefore chiral invariant as it should be, whereas in the heavy quark limit, $V_{\text{eff}} \rightarrow C \lambda^c/2$ is simply a scalar vertex. The Gell-Mann matrix $\lambda^c$ provides the usual color vector coupling as expected in a QM. We anticipate that the dynamical generation of a quark mass will also generate a scalar coupling for light quarks, and this results in an effective vertex, including a chiral invariant vertex $\beta$ and the standard scalar vertex $\gamma$.

The dependence in the relative momentum must comply with the linear confinement which is derived from the string picture,

\[ V_{\epsilon}(x) = \frac{16}{3c^2} V_{\epsilon}(x) = \frac{16}{3c^2} \sigma |x| e^{-\epsilon |x|}, \]

where $\sigma \sim 200$ MeV/Fm is the string constant and $C$ is the algebraic color factor defined in Eq. (3). The damping factor $\epsilon$ regularizes the Fourier transform,

\[ V_{\epsilon}(p) = -\frac{16}{3c^2} 8\pi \sigma \left( \frac{1}{(|p|^2 + \epsilon^2)^2} - \frac{4\epsilon^2}{(|p|^2 + \epsilon^2)^3} \right), \]

and in the limit $\epsilon \rightarrow 0$ we have

\[ -i V_0(p - q) = -i \left\{ \frac{8\pi \sigma}{|p - q|^2} \right\}. \]

**IV. MASS GAP EQUATION**

We solve the mass gap equation using the Schwinger-Dyson formalism,

\[ S^{-1} = S_0^{-1} - \Sigma, \]

where the dressed propagator is defined in eq. (6) and the free propagator has a similar definition with the quark bare mass $m_0$. We want to determine the constituent
The mass gap equation can be much simplified in the spin formalism. Some useful relations we will use for this purpose are

\[ u^{(1)}_{s}(p) v_{s'}(p) = 0 \cdot [\sigma(i\sigma_2)]_{ss'}, \]
\[ u^{(1)}_{s}(p) \beta v_{s'}(p) = c_p \hat{p} \cdot [\sigma(i\sigma_2)]_{ss'}, \]
\[ u^{(1)}_{s}(p) \alpha \gamma v_{s'}(p) = -\hat{q} \cdot \hat{p} \cdot \hat{q} \cdot \hat{p} \cdot [\sigma(i\sigma_2)]_{ss'}, \]
\[ u^{(1)}_{s}(p) \delta \gamma v_{s'}(p) = -(s_p \delta \hat{q} \cdot \hat{p} + (1-s_p)\hat{p} \cdot \hat{q} \cdot \hat{p} \cdot [\sigma(i\sigma_2)]_{ss'}. \]

With eqs. 10, 7 and 15, we arrive at the expected relation 26

\[ \pi_s(p) S^{-1}(p) v_{s'}(p) = 0, \]

implying that the propagator is diagonal in the particle-antiparticle projection. The mass gap equation becomes

\[ \pi_s(p) S_0^{-1}(p) v_{s'}(p) - \pi_s(p) \Sigma(p) v_{s'}(p) = 0. \]

The free propagator term for a zero bare mass, \( m_0 = 0 \), is simply

\[ u^{(1)}_{s}(p) \beta (-i) \not p v_{s'}(p) = i p c_p \hat{p} \cdot [\sigma(i\sigma_2)]_{ss'}. \]

The self-energy term is directly derived from the diagram of Fig. 8 using the double vertex loop \( V_{\text{eff}} \) already obtained in 10.

\[ \Sigma(p) = \int \frac{d^3q}{(2\pi)^3} (s_k - c_k \hat{k} \cdot \gamma) (s_q \beta + c_q \hat{q} \cdot \alpha) \beta \times \]
\[ \times (s_k - c_k \hat{k} \cdot \gamma) C^2 \frac{3}{16} (-i) V_\varepsilon(|p - q|) \]
\[ = \int \frac{d^3q}{(2\pi)^3} \left( \left( s_k^2 - c_k^2 \right) s_q \beta - 2 s_k c_k s_q \hat{k} \cdot \alpha + (s_k^2 + c_k^2) c_q \hat{q} \cdot \alpha + 2 s_k c_k c_q \hat{k} \cdot \hat{q} \cdot \beta - 2 c_k^2 c_q \hat{k} \cdot \hat{q} \cdot \hat{k} \cdot \alpha \right) \beta (-i) V_\varepsilon(|p - q|), \]

where we used the properties of the \( \beta \) and \( \alpha \) matrices. The result for the mass gap self-energy term is

\[ \pi_s(p) \Sigma(p) v_{s'}(p) = \]
\[ = \int \frac{d^3q}{(2\pi)^3} (-i) V_\varepsilon(|p - q|) \left[ (s_k^2 - c_k^2) s_q (-c_p \hat{p}) - 2 s_k c_k s_q (-\hat{k} + (1-s_p)\hat{k} \cdot \hat{p}) + c_q (-\hat{q} + (1-s_p)\hat{q} \cdot \hat{p} \cdot \hat{p}) + 2 s_k c_k c_q \hat{k} \cdot \hat{q} (-c_p \hat{p}) - 2 c_k^2 c_q \hat{k} \cdot \hat{q} (-\hat{k} + (1-s_p)\hat{k} \cdot \hat{p} \cdot \hat{p}) \right] [\sigma(i\sigma_2)]. \]

As we can see from 18 and 20 both terms of the mass gap equation are proportional to \( \hat{p} \cdot \sigma(i\sigma_2) \). Since the Pauli matrices \( \sigma \) are linearly independent, we can substitute \( \sigma(i\sigma_2) \) by \( \hat{p} \) and still have a mass gap condition. With this simplification the mass gap equation reduces to,

\[ i p s_p - \int \frac{d^3q}{(2\pi)^3} \left[ (c_k^2 - s_k^2) s_q c_p + 2 s_k c_k s_q \hat{k} \cdot \hat{p} - c_q s_p \hat{q} \cdot \hat{p} - 2 s_k c_k c_q \hat{k} \cdot \hat{q} \right] + 2 c_k^2 c_q s_p \hat{k} \cdot \hat{q} \cdot \hat{k} \cdot \hat{p} \right] i V_\varepsilon(|p - q|) = 0. \]

Notice that if we take the integrand and set \( q = p \) we will get \( 0 \times V_\varepsilon(0) \). In section V we will deal numerically with this IR behavior.

V. NUMERICAL SOLUTION OF THE MASS GAP EQUATION

The mass gap equation is a difficult non-linear integral equation, that does not converge with the usual iterative methods. We developed a method to solve the mass gap equation with a differential equation, using a convergence parameter \( \lambda \). This parameter is the radius of a sphere centered in \( u = p - q = 0 \) it allows us to separate the integral into an integral inside the sphere and another outside of it,

\[ \int \frac{d^3u}{(2\pi)^3} f(p, u) V_\varepsilon(u) = \int_{u=0} \frac{d^3u}{(2\pi)^3} f(p, u) V_\varepsilon(u) + \int_{R^3-0} \frac{d^3u}{(2\pi)^3} f(p, u) V_\varepsilon(u). \]

In our model \( f(p, u) \) is the function dependent on the chiral angle presented in eq. 21.

Let us first focus on the integral inside the sphere, where we have \( u < \lambda \). Eventually we will take the limit where \( \lambda \to 0 \) and this term will vanish. But for now we will expand the function \( f \) around \( u = 0 \) and take only
the first non-vanishing term,

\[
\left. \int_{0}^{\frac{d^3u}{(2\pi)^3}} f(p, u, \omega) V_\varepsilon(u) \right|_{\varepsilon = 0} \approx \int_{0}^{\infty} \frac{d^3u}{(2\pi)^3} \frac{\partial^2 f}{\partial u^2} \bigg|_{u = 0} V_\varepsilon(u)
\]

where \( \omega \) is the cosine of the angle between \( p \) and \( u \). For our particular model we have

\[
\int_{-1}^{1} d\omega \left. \frac{\partial^2 f}{\partial u^2} \right|_{u = 0} = \frac{\sin(2\varphi_p) - 2p \cos(2\varphi_p) \varphi'_p + p^2 \varphi''_p}{3p^2} \tag{24}
\]

In what concerns the integral outside the sphere \( u > \lambda \) we can take from the beginning \( \varepsilon = 0 \) since this integral is already regulated by \( \lambda \). In this case we have

\[
\int_{R^3 - \varepsilon}^{\infty} \frac{d^3u}{(2\pi)^3} f(p, u) V_0(u) = \\
= \int_{0}^{\infty} \left. d\omega \frac{1}{u^2} f(p, u, \omega) \right|_{u = 0} = \int_{1}^{\infty} du \int_{1}^{\infty} d\omega \frac{1}{u^2} f(p, u, \omega).
\]

Placing the two terms in the mass gap equation we finally get the equation that we can iterate to find the solution,

\[
3p s_p + \frac{\sigma}{\pi} \lambda \left( \varphi''_p - \frac{2}{p} \cos(2\varphi_p) \varphi'_p + \frac{1}{p^2} \sin(2\varphi_p) \right) + 6 \int_{1}^{\infty} du \int_{1}^{\infty} d\omega \frac{1}{u^2} f(p, u, \omega) = 0 \tag{26}
\]

![Fig. 4: Testing the convergence of the numerical method with the quark condensate \( \langle \bar{\psi}\psi \rangle \).](image)

**VI. RESULTS AND CONCLUSION**

In this paper we build a QM for the coupling of quark to a scalar string. The quark confining interaction has a single parameter \( \sigma \). This QM matches the apparently conflicting vector coupling of QCD with a scalar confinement. Our model can be interpreted as a double vertex that couples the quark to the string, or alternatively as a \( \gamma_5 S(k) \gamma_0 \) vertex. Either way this vertex decomposes in the sum of a scalar vertex 1 and a chiral invariant \( k \cdot \gamma \) vertex weighed by simple functions of the dynamical quark mass. In the chiral limit the scalar vertex vanishes, while in the heavy quark limit the confining potential is essentially scalar. Our results for the weighing factors of the scalar vertex and the remaining chiral invariant vertex are shown in Fig. 6.

We solve the mass gap equation for the dynamical generation of the quark mass with the \( S \backslash \)SB, and we in-
deed generate the constituent quark mass. We show that S\(\chi\)SB not only generates a quark mass, but generates also a scalar vertex for the confinement. The results are encouraging because the quark condensate indeed increases when compared with the simpler one vertex vector confining potential.

It is clear that the next step of this work, will consist in adding the shorter range one gluon exchange potential to the confining potential. The resulting model will have two parameters, one for the short range potential, and another one for the confining potential. These parameters will be determined in the fit of the hadron spectrum. In what concerns the mass gap equation, we expect that this will further enhance the quark condensate, possibly up to the expected \((-230 \, \text{MeV})^3\).

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[1] J. Bjorken, “Why Hadron Physics”, document available at the APS Topical Group on Hadronic Physics, [http://fafnir.phyast.pitt.edu/topical/bj.ps](http://fafnir.phyast.pitt.edu/topical/bj.ps).
[2] A. P. Szczepaniak and E. S. Swanson, Phys. Rev. D 55, 3087 (1997) [arXiv:hep-ph/9611310]; A. P. Szczepaniak and E. S. Swanson, Phys. Rev. D 65, 025012 (2002) [arXiv:hep-ph/0107078].
[3] J. P. Lagae, Phys. Lett. B 240, 451 (1990).
[4] A. Barchielli, N. Brambilla and G. M. Prosperi, Nuovo Cim. A 103, 59 (1990).
[5] N. Brambilla, A. Pineda, J. Soto and A. Vairo, Phys. Rev. D 60, 091502 (1999) [arXiv:hep-ph/9903355].
[6] P. Bicudo, Phys. Rev. C 67, 035201 (2003).
[7] S. Adler and A. Davis, Nucl. Phys. B 244, 469 (1984); S. L. Adler, Prog. Theor. Phys. Suppl. 86, 12 (1986).
[8] M. S. Bhagwat, M. A. Pichowsky, C. D. Roberts and P. C. Tandy, Phys. Rev. C 68, 015203 (2003) [arXiv:nucl-th/0304003].
[9] R. Alkofer and L. von Smekal, Phys. Rept. 353, 281 (2001) [arXiv:hep-ph/0007355].
[10] A. B. Henriques, B. H. Kellett and R. G. Moorhouse, Phys. Lett. B 64, 85 (1976); A. B. Henriques, B. H. Kellett and R. G. Moorhouse, Annals Phys. 93, 125 (1975); J. Dias de Deus, A. B. Henriques and J. M. Pulido, Z. Phys. C 7, 157 (1981); A. B. Henriques, Z. Phys. C 11, 31 (1981).
[11] L.-H. Chan, Phys. Lett. 71B, 422 (1977).
[12] N. Isgur and G. Karl, Phys. Rev. D 18, 4187 (1978); N. Isgur, Phys. Rev. D 62, 014025 (2000) [arXiv:hep-ph/9910272].
[13] W. Kwong, J. L. Rosner and C. Quigg, Ann. Rev. Nucl. Part. Sci. 37, 325 (1987).
[14] C. Michael, Phys. Rev. Lett. 56, 1219 (1986).
[15] D. Gromes, Z. Phys. C 26, 401 (1984); W. Lucha, F. F. Schöberl and D. Gromes, Phys. Rept. 200 (1991) 127.
[16] G. S. Bali, K. Schilling and A. Wachter, Phys. Rev. D 56, 2566 (1997) [arXiv:hep-lat/9703019].
[17] T. J. Allen, M. G. Olsson and S. Vesli, Phys. Rev. D 62, 094021 (2000) [arXiv:hep-ph/0001227].
[18] N. Isgur and J. Paton, Phys. Rev. D 31, 2910 (1985); N. Isgur and J. Paton, Phys. Lett. B 124, 247 (1983).
[19] F. E. Low, Phys. Rev. D 12, 163 (1975).
[20] S. Nussinov, Phys. Rev. Lett. 34, 1286 (1975).
[21] Y. A. Simonov, Phys. Atom. Nucl. 60, 2069 (1997) [Yad. Fiz. 60, 2252 (1997)] [arXiv:hep-ph/9704301]; Y. A. Simonov, Z. Phys. C 53, 419 (1992); Y. A. Simonov, Phys. Atom. Nucl. 63, 94 (2000) [Yad. Fiz. 63, 106 (2000)]; Y. S. Kalashnikova, A. V. Nefediev and Y. A. Simonov, Phys. Rev. D 64, 014037 (2001) [arXiv:hep-ph/0103274]; S. M. Fedorov and Y. A. Simonov, [arXiv:hep-ph/0306216]
[22] P. Bicudo, N. Brambilla, E. Ribeiro and A. Vairo, Phys. Lett. B 442, 349 (1998) [arXiv:hep-ph/9807460].
[23] A. V. Nefediev, [arXiv:hep-ph/0308274]
[24] F. V. Gubarev, M. I. Polikarpov and V. I. Zakharov, [arXiv:hep-th/9812030]
[25] H. G. Dosch and Y. A. Simonov, Phys. Lett. B 205, 339 (1988).
[26] P. J. Bicudo, Phys. Rev. C 60, 035209 (1999).
[27] P. J. Bicudo and J. E. Ribeiro, Phys. Rev. D 42, 1611 (1990).
[28] P. J. Bicudo, J. E. Ribeiro and A. V. Nefediev, Phys. Rev. D 65, 085026 (2002) [arXiv:hep-ph/0201173].
[29] A. Le Yaouanc, L. Oliver, O. Pene and J.-C. Raynal, Phys. Rev. D 29, 1233 (1984); Phys. Rev. D 31, 137 (1985). P. J. Bicudo and J. E. Ribeiro, Phys. Rev. D 42, 1611 (1990).