The structure of $f_0(980)$ from charmed mesons decays

Ignacio Bediaga$^1$, Fernando S. Navarra$^2$ and Marina Nielsen$^2$

$^1$Centro Brasileiro de Pesquisas Físicas
Rua Xavier Sigaud 150, 22290-180 Rio de Janeiro, RJ, Brazil

$^2$Instituto de Física, Universidade de São Paulo
C.P. 66318, 05315-970 São Paulo, SP, Brazil

Abstract

We use the QCD sum rules to evaluate the form factors associated with the semileptonic decays of $D_s$ and $D$ mesons into $f_0(980)$. We consider the $f_0(980)$ meson as a quark-antiquark state with a mixture of strange and light components. The decay rates are evaluated in terms of the mixing angle. Using the same form factors to evaluate nonleptonic decays in the framework of the factorization approximation we conclude that the importance of the light quarks in $f_0(980)$ is not negligible.
Recent experimental data coming from charmed mesons decays have opened new possibilities to understand the spectroscopy of the controversial light scalar mesons, since they are abundantly produced in these decays \[1\]. Actually, with the firm identification of the scalars $\sigma$ and $\kappa$, the observed scalar states below 1.5 GeV are too numerous to be accommodated in a single $q\bar{q}$ multiplet. This proliferation of the scalar mesons is consistent with two nonets, one below 1 GeV region and another one near 1.5 GeV \[2\]. In this new scheme the light scalars (the isoscalars $\sigma(500)$, $f_0(980)$, the isodoublet $\kappa$ and the isovector $a_0(980)$) would form an SU(3) flavor nonet. In the naive quark model the flavor structure of these scalars would generically be:

$$
\sigma = \frac{\cos(\alpha)}{\sqrt{2}}(u\bar{u} + d\bar{d}) - \sin(\alpha)s\bar{s}, \quad f_0 = \cos(\alpha)s\bar{s} + \frac{\sin(\alpha)}{\sqrt{2}}(u\bar{u} + d\bar{d}),
$$  \hspace{1cm} (1)

$$
a_0^+ = \frac{1}{\sqrt{2}}(u\bar{u} - d\bar{d}), \quad a_0^0 = ud, \quad a_0^- = d\bar{u},
$$  \hspace{1cm} (2)

$$
\kappa^+ = u\bar{s}, \quad \kappa^0 = d\bar{s}, \quad \bar{\kappa}^0 = s\bar{d}, \quad \kappa^- = s\bar{u},
$$  \hspace{1cm} (3)

where we have already allowed a mixing between the isoscalars $s\bar{s}$ and $(u\bar{u} + d\bar{d})$. Although the predominant $s\bar{s}$ nature of the $f_0(980)$ has been supported by the radiative decay $\phi \rightarrow f_0(980)\gamma$, and by some theoretical calculations \[3, 4\], there is no fundamental theoretical reason to expect $\alpha = 0$, as we find in the vector meson sector. Actually, since instantons are supposed to be as important in the scalar sector as they are in the pseudoscalar sector \[5\], we would expect a mixing in the scalar sector similar to what we have in the pseudoscalar sector. In this sense, the measurements of $J/\psi \rightarrow f_0(980)\phi$ and $J/\psi \rightarrow f_0(980)\omega$ with similar branching ratios \[4\], indicating that $f_0(980)$ is not purely an $s\bar{s}$ state, can not be taken as a surprise. In ref. \[8\] these $J/\psi$ decays were used to estimate the mixing angle in Eq. (1), giving $\alpha = (34 \pm 6)^0$. A similar mixing angle, $35^0 \leq \alpha \leq 55^0$ \[8\], was found analysing the experimental results $D_s^+ \rightarrow f_0(980)\pi^+$ and $D_s^+ \rightarrow \phi\pi^+$ \[10\]. Using $f_0(980)$ as a pure $s\bar{s}$ state (i.e., using $\alpha = 0$) the authors of ref. \[11\], could not reproduce the experimental result of $D_s^+ \rightarrow f_0(980)\pi^+ / D_s^+ \rightarrow \phi\pi^+$ \[10\]. They concluded that there is room for a sizable light quark component in $f_0(980)$, corresponding to a mixing angle of about $\alpha \sim 40^0$.

In this work we propose that experimental results of the semileptonic decays $D_s^+ \rightarrow f_0(980)\ell^+\nu_\ell$ and $D^+ \rightarrow f_0(980)\ell^+\nu_\ell$, can be used to get the minimum bias estimate of the importance of the light quark content in the scalar-isoscalar $f_0(980)$, in the quark-antiquark scenario. The hadronic part of the current of the $D_s^+$ decay, can produce the $f_0(980)$ only through the $s\bar{s}$ component, while the hadronic part of the $D^+$ decay can produce
the $f_0(980)$ only through $d\bar{d}$. To observe these decays it would be necessary a high statistics experiment and also a tagging to separate the semi-leptonic $D_s^+$ from the $D^+$ decays. These two conditions could be satisfied next year by the CLEO-C experiment [12].

In order to estimate theoretically these semileptonic $D_s^+$ and $D^+$ decays, we use the method of QCD sum rules [13] which has been successfully applied to several semileptonic decay processes. Since the semileptonic decay is supposed to occur on the quark level, the decay rate should depend crucially on the direct coupling of the resonances to the quark currents. The coupling of the $f_0(980)$ to the scalar current

$$j_s = \bar{s}s \cos(\alpha) + (\bar{u}u + \bar{d}d)\frac{\sin(\alpha)}{\sqrt{2}},$$

(4)

can be parametrized as

$$\langle 0|j_s|f_0(p)\rangle = \lambda_{f_0} = m_{f_0}\bar{f}_{f_0},$$

(5)

and can be determined by the QCD sum rule based on the two-point correlation function

$$\Pi(q^2) = i \int d^4xe^{iq\cdot x}\langle 0|T[j_s(x)j_s^\dagger(0)]|0\rangle.$$  

(6)

This same correlation function was studied in ref. [14] in the case $\alpha = 0$. Following the same procedure and considering a general mixing angle $\alpha$ we get the sum rule (up to order $m_s$):

$$\lambda_{f_0}^2e^{-m_{f_0}^2/M^2} = \cos^2(\alpha)\left[\frac{3}{8\pi^2}\int_0^{u_0} duue^{-u/M^2}
+ m_s e^{-m_s^2/M^2}\left(3\langle\bar{s}s\rangle + \frac{\langle\bar{g}_s\sigma.Gs\rangle}{M^2}\right)\right]
+ \sin^2(\alpha)\left(\frac{3}{8\pi^2}\int_0^{u_0} duue^{-u/M^2}\right).$$

(7)

In the numerical analysis of the sum rules, the values used for the strange quark mass and condensates are: $m_s = 0.14$ GeV, $\langle\bar{s}s\rangle = 0.8\langle\bar{q}q\rangle$, $\langle\bar{q}q\rangle = -(0.23)^3$ GeV$^3$, $\langle\bar{g}_s\sigma.Gs\rangle = m_0^2\langle\bar{s}s\rangle$ with $m_0^2 = 0.8$ GeV$^2$. We evaluate the 2-point sum rules in the same stability window found in [14]: $1.2 \leq M^2 \leq 2.0$ GeV$^2$. In Fig. 1 we show the different contributions to $\lambda_{f_0}^2$ as a function of the Borel mass using the continuum threshold $u_0 = 1.6$ GeV$^2$ and the mixing angle $\alpha = 37^0$. We see that $\lambda_{f_0}^2$ is very stable, as a function of the Borel mass, in the considered Borel range. The coupling is not very sensitive to changes in the values of the continuum threshold and mixing angle. For $u_0 = (1.6 \pm 0.1)$ GeV$^2$ and $0 \leq \alpha \leq 37^0$ we obtain

$$\lambda_{f_0} = (0.19 \pm 0.02) \text{GeV}^2.$$
We can proceed and evaluate the $f_0(980)$ mass from the above sum rule, taking the derivative of Eq. (7) with respect to $M^{-2}$ and dividing the resulting sum rule by Eq. (7). For $u_0 = 1.6$ GeV$^2$ and the mixing angle $\alpha = 37^0$, the mass is in a very good agreement with the experimental number and is very stable, as a function of the Borel mass. In the considered Borel range we obtain: $m_{f_0} = (0.98 \pm 0.01)$ GeV. For smaller (bigger) values of $\alpha$ we get a bigger (smaller) mass. The best agreement with the experimental result is obtained for $\alpha \sim 37^0$.

The coupling given in Eq. (8) will be used in the analysis of the semileptonic decays of $D_s$ and $D$ into $f_0(980)$. The $D_I \to f_0(980)\ell\nu_\ell$ form factors are defined through the matrix elements

$$\langle f_0(p')|\bar{s}\gamma_\mu(1-\gamma_5)c|D_s(p)\rangle = \cos(\alpha) \left(f_+^{D_s}(t)(p+p')_\mu + f_-^{D_s}(t)q_\mu\right),$$

and

$$\langle f_0(p')|\bar{d}\gamma_\mu(1-\gamma_5)c|D(p)\rangle = \frac{\sin(\alpha)}{\sqrt{2}} \left(f_+^{D}(t)(p+p')_\mu + f_-^{D}(t)q_\mu\right),$$

with $t = q^2$ and $q = p - p'$. Since in the decay rate the form factor $f_-(t)$ is multiplied by the difference of the lepton masses, its contribution is negligible for both $e$ and $\mu$ decays. Therefore, the differential semileptonic decay rates are given by

$$\frac{d\Gamma(t)}{dt} = \cos^2(\alpha) \frac{G_F^2|V_{cs}|^2}{192\pi^3m_{D_s}^3}\lambda^{3/2}(m_{D_s}^2, m_{f_0}^2, t)(f_+^{D_s}(t))^2,$$
for $D_s^+ \to f_0(980)\ell^+\nu_\ell$, and
\[
\frac{d\Gamma(t)}{dt} = \frac{\sin^2(\alpha)}{2} \frac{G_F^2|V_{cd}|^2}{192\pi^3m_D^3} \lambda^{3/2}(m_D^2, m_F^2, t)(f_D^P(t))^2,
\]
(12)
for $D^+ \to f_0(980)\ell^+\nu_\ell$. In Eqs. (11) and (12), $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$, $G_F$ is the Fermi coupling constant and $V_{cs}$ and $V_{cd}$ are the Cabibbo-Kobayashi-Maskawa transition elements.

Using the QCD sum rule technique [13], the form factors in Eqs. (9) and (10) can be evaluated from the time ordered product of the interpolating fields for $D_I$ and $f_0$, and the weak current $j_\mu^W = \bar{q}_I \gamma_\mu (1 - \gamma_5)c$ (where $q_I$ is $s$ or $d$ for $D_I$ being $D_s$ or $D$ respectively):
\[
T_\mu(p, p') = i^2 \int d^4xd^4y \langle 0|T[j_s(x)j_\mu^W(y)j_D^I(0)]|0\rangle e^{i(p'.x + q.y)},
\]
(13)
where the $D_I^+$ meson in the initial state is interpolated by the pseudoscalar current
\[
j_{D_I}(x) = \bar{q}_I(x)i\gamma_5c(x),
\]
(14)
and the $f_0(980)$ is interpolated by the scalar current given in Eq. (11).

In order to evaluate the phenomenological side we insert intermediate states for $D_I$ and $f_0$, we use the definitions in Eqs. (9) and (10), and obtain the following relations
\[
T_\mu^{\text{phen}}(p, p') = \frac{m_D^2 f_D s}{m_c + m_s} \lambda_0 \sin(\alpha) \frac{f_D^P(t)(p + p')_\mu + f_D^S(t)_\mu}{(m_D^2 - p^2)(m_0^2 - p'^2)}
\]
+ contributions of higher resonances,
(15)
for $D_s^+ \to f_0(980)$ and
\[
T_\mu^{\text{phen}}(p, p') = \frac{m_D^2 f_D s}{m_c + m_q} \lambda_0 \sin(\alpha) \frac{f_D^P(t)(p + p')_\mu + f_D^S(t)_\mu}{\sqrt{2}(m_D^2 - p^2)(m_0^2 - p'^2)}
\]
+ contributions of higher resonances,
(16)
for $D^+ \to f_0(980)$.

In the above equations we have used the standard definition of the couplings of $D_I$ with the corresponding current:
\[
\langle 0|j_{D_I}|D_I\rangle = \frac{m_D^2 f_D}{m_c + m_q}.
\]
(17)

The three-point function Eq.(13) can be evaluated by perturbative QCD if the external momenta are in the deep Euclidean region
\[
p \ll (m_c + m_s)^2, \quad p'^2 \ll 4m_s^2, \quad t \ll (m_c + m_s)^2.
\]
(18)
In order to approach the not-so-deep-Euclidean region and to get more information on the nearest physical singularities, nonperturbative power corrections are added to the perturbative contribution. In practice, only the first few condensates contribute significantly, the most important ones being the 3-dimension, quark condensate, and the 5-dimension, mixed (quark-gluon) condensate. For the invariant structure, \((p + p')_\mu\), we can write
\[
T^{\text{theor}}(p^2; p'^2, t) = \frac{-1}{4\pi^2} \int_{(m_c + m_{f_0})^2}^{\infty} ds \int_0^{\infty} du \frac{\rho_+ (s, u, t)}{(s - p^2)(u - p'^2)}
+ T_+^{D=3}\langle \bar{q}_1 q_1 \rangle + T_+^{D=5}\langle \bar{q}_1 g_s \sigma G q_1 \rangle + \cdots .
\]

The perturbative contribution is contained in the double discontinuity \(\rho_+\).

In order to suppress the condensates of higher dimension and at the same time reduce the influence of higher resonances, the series in Eq. (19) is double Borel improved. Furthermore, we make the usual assumption that the contributions of higher resonances are well approximated by the perturbative expression with appropriate continuum thresholds \(s_0\) and \(u_0\). By equating the Borel transforms of the phenomenological expression for the \((p + p')_\mu\) invariant structure in Eqs. (15) and (16), and that of the “theoretical expression”, Eq. (19), we obtain the sum rules for the form factor \(f_+^{D_i}(t)\). The sum rule for \(f_+^{D_i}(t)\) (at the order \(m_s\)) is given by:
\[
f_+^{D_i}(t) C e^{-m_{D_i}^2/M^2} e^{-m_{f_0}^2/M^2} = \frac{-1}{4\pi^2} \int_{(m_c + m_s)^2}^{s_0} ds \int_0^{u_0} du \left[ e^{-s/M^2} e^{-u/M^2} \rho_+ (s, u, t) \right]
+ \frac{2}{2} e^{-m_{D_i}^2/M^2} \left[ -m_c + 2m_s + \frac{m_c^2 m_s}{2M^2} \right] + m_c^2 \frac{\langle \bar{s} s \rangle}{s/s} e^{-m_{D_i}^2/M^2} \left[ \frac{m_c^2 (m_c - m_s)}{8M^4} - \frac{2m_c - m_s}{6M^2} \right]
\]
\[
+ m_c^2 \left[ 4m_c - 3m_s \right] - 2t (m_c - m_s) \right) - \frac{m_c^2 - 2m_s}{6M^2} + \frac{m_c^2 m_s - 2t (m_c - m_s)}{24M^2 M^2} \right] .
\]

where \(C = \frac{m_c^2 f_{D_i} f_{D_s}}{m_c + m_s} \lambda \) and
\[
\rho_+ (s, u, t) = \frac{3}{\lambda^{3/2}(s, u, t)} \left\{ u \left[ 2m_c m_s (2m_c^2 - s - t + u) + m_c^2 (s - t + u) + s(-s + t + u) \right]
- (2m_c^2 - s - t + u)(su + m_c m_s (s - t + u)) \right\} \Theta(s - s_M),
\]
with \(s_M = m_c^2 + \frac{m_c^2 u}{m_c^2 - t}\). The sum rule for \(f_+^{D}(t)\) can be obtained from Eq. (20) by just neglecting the \(m_s\) terms, changing \(\langle \bar{s} s \rangle\) by \(\langle \bar{q} q \rangle\) and \(D_s\) by \(D\).

The decay constant \(f_{D_i}\) defined in Eq. (17), and appearing in the constant \(C\), can also be determined by sum rules obtained from the appropriate two-point functions [11, 15].

The value of the charm quark and meson masses are: \(m_c = 1.3\) GeV, \(m_{f_0} = 0.98\) GeV, \(m_{D_s} = 1.97\) GeV and \(m_D = 1.87\) GeV. For the \(D_s\) and \(D\) decay constants we use \(f_{D_s} = \)
FIG. 2: Various contributions to the OPE of the form factor $f_{+}^{D^*(t)}$, at zero momentum transfer, as a function of the Borel parameter $M^2$. Solid curve: total contribution; long-dashed: perturbative; dashed: quark condensate; dot-dashed mixed condensate contribution.

$(0.22 \pm 0.02)$ GeV and $f_D = (0.17 \pm 0.02)$ GeV. For the continuum thresholds we take the values discussed in refs. $[11, 14, 15]$: $s_0 = (7.7 \pm 1.1)$ GeV$^2$ for $D_s$, $s_0 = (6.0 \pm 0.2)$ GeV$^2$ for $D$ and $u_0 = (1.6 \pm 0.1)$ GeV.

We evaluate our sum rules in the range $4.0 \leq M^2 \leq 8.0$ GeV$^2$, at a fixed ratio $M^2/M^2 = (m_{D_I}^2 - m_c^2)/m_{f_0}^2$, which is compatible with the Borel ranges used for the two-point functions in refs. $[11, 14]$. In Fig. 2 we show the different contributions to the form factor $f_{+}^{D^*}$ at zero momentum transfer, from the sum rule in Eq. (20), as a function of the Borel variable $M^2$. We see that the perturbative contribution is the largest one, and that the mixed condensate contribution is negligible. A similar behavior is also obtained for $f_{+}^{D}(t)$. Varying the continuum thresholds $s_0$ and $u_0$, and the couplings $\lambda_{f_0}$, $f_{D_s}$ and $f_D$ in the ranges given above, we get for the form factors at $t = 0$:

$$f_{+}^{D_I}(0) = \begin{cases} 
0.50 \pm 0.13 & \text{for } D_s \\
0.53 \pm 0.15 & \text{for } D
\end{cases}$$

Our values for $f_{+}^{D^*}(0)$ are compatible with the values found in refs. $[8, 17]$, where the non-leptonic decay of the $D^+_s$ meson were studied.

The $t$ dependence of the form factors evaluated at $M^2 = 7$ GeV$^2$ in the range $-0.5 \leq t \leq 0$ GeV$^2$ can be fitted by a linear expression

$$f_{+}^{D_I}(t) = f_{+}^{D_I}(0) + At.$$
This expression is consistent, in the considered range of momentum transfer, with a monopole expression \( f^{D^+}(t) = \frac{f^{D^+}(0)}{M_P^2} \), with the mass of the pole \( M_P = (1.70 \pm 0.05) \) GeV for \( D_s \) and \( M_P = (1.95 \pm 0.05) \) GeV for \( D \). It is interesting to notice that \( M_P \) is compatible with the values found for the decays \( D_s \to \eta \) [19] and \( D \to \kappa \) [18].

In the limits of the variables and the continuum thresholds discussed above we obtain for the semileptonic decay widths

\[
\begin{align*}
\Gamma(D^+_s \to f_0(980)\ell^+\nu_\ell) &= \cos^2(\alpha)(8.1 \pm 4.1) \text{ GeV,} \\
\Gamma(D^+ \to f_0(980)\ell^+\nu_\ell) &= \sin^2(\alpha)(1.5 \pm 0.8) \text{ GeV,}
\end{align*}
\]  

(24)

where we have used \( V_{cs} = 0.975 \) and \( V_{cd} = 0.22 \). The ratio between the two decay widths given in Eq. (24) gives us a direct information about the mixing angle:

\[
\frac{\Gamma(D^+_s \to f_0(980)\ell^+\nu_\ell)}{\Gamma(D^+ \to f_0(980)\ell^+\nu_\ell)} = \frac{56 \pm 2}{\tan^2(\alpha)}.
\]  

(25)

For \( \alpha \sim 37^0 \), from Eq. (25) we can conclude that the semileptonic decay width of \( D_s \) into \( f_0(980) \) would be around one hundred times larger than the semileptonic decay width of \( D \) into \( f_0(980) \). Since in the semileptonic decays there are no complications due to strong interactions, we believe that an experimental measurement of the ratio in Eq. (25) is the cleanest way to evaluate the mixing angle \( \alpha \).

It is interesting to notice that if the current in Eq. (4) were used in ref. [11] for \( f_0(980) \), instead of a pure \( \bar{s}s \) current, the ratio calculated there would change to

\[
R = \frac{\Gamma(D^+_s \to f_0(980)\pi^+)}{\Gamma(D^+_s \to f_0(980)\phi^+)} = \cos^2(\alpha)(0.44 \pm 0.18) .
\]  

(26)

Therefore, for \( \alpha \sim 37^0 \), which is compatible with the mixing angle found in refs. [8, 9], the above ratio would be reduced to \( R = (0.28 \pm 0.11) \) in agreement with the experimental result \( R^{exp} = (0.210 \pm 0.069) \) [10].

At this point we might conclude by saying that \( f_0(980) \) is well described by the relation in Eq. (1), with the mixing angle being \( \alpha \sim 35^0 \), and by making the prediction in Eq. (25). However, there are two experimental facts which do not fit in this picture and that might even be in contradiction with Eq. (1):

- A mixing angle in \( f_0(980) \) also implies that \( \sigma \) would have a strange component (see Eq. (1)). Therefore, the decay \( D^+_s \to \sigma \pi^+ \) should also occur. As a matter of fact,
the ratio $\Gamma(D_s^+ \to \sigma \pi^+)/\Gamma(D_s^+ \to f_0(980)\pi^+)$ is proportional to $\tan^2(\alpha)$. However, the E791 Collaboration did not observe any contribution from $D_s^+ \to \sigma \pi^+$ to the $D_s^+ \to \pi^+\pi^+\pi^-$ decay [10] pointing towards $\alpha \sim 0$.

- In the framework of generalized factorization, the amplitude for the $D^+ \to f_0(980)\pi^+$ decay is given by (neglecting the annihilation term)

$$A(D^+ \to f_0(980)\pi^+) = \frac{G_F}{\sqrt{2}} V_{cd} V_{ud}^* \left( c_1(\mu) + \frac{c_2(\mu)}{3} \right) f_\pi (m_D^2 - m_{f_0}^2) \frac{\sin(\alpha)}{\sqrt{2}} f^D_f(0), \quad (27)$$

since $\langle 0 | V_\mu | f_0 \rangle = 0$ due to charge conjugation invariance and conservation of vector current. In Eq. (27) $c_i$ is the Wilson coefficient entering the effective weak Hamiltonian. Therefore, using the results in Eq. (22) we get

$$\frac{\Gamma(D_s^+ \to f_0(980)\pi^+)}{\Gamma(D^+ \to f_0(980)\pi^+)} = 46 \pm 2 \frac{\tan^2(\alpha)}{\tan^2(\alpha)}. \quad (28)$$

From the E791 Collaboration, the experimental value of this quantity is [10]:

$$\left( \frac{\Gamma(D_s^+ \to f_0(980)\pi^+)}{\Gamma(D^+ \to f_0(980)\pi^+)} \right)^{exp} = 13.3 \pm 0.4, \quad (29)$$

which would lead to $\alpha \sim 62^\circ$.

It is important to remember that in the analysis of the nonleptonic decays we are using the factorization approximation. Therefore, the apparent inconsistency between the data and the mixing angle can still be due to this approximation. This is why the measurement of the ratio in Eq. (28) would bring important information about this puzzle. If the light scalars could not be seen in the semileptonic decays, this would clearly indicate that their structure is more complicated than simple quark-antiquark states. One possibility is that they are four-quark states, as suggested in refs. [2, 20, 21].

To summarize, we have presented a QCD sum rule study of the $D_s^+$ and $D$ semileptonic decays to $f_0(980)$, considered as a mixture of the scalars $\bar{u}u$ and $\bar{d}d$. We have evaluated the $t$ dependence of the form factors $f^D_{+}(t)$ in the region $-0.5 \leq t \leq 0$ GeV$^2$. The $t$ dependence of the form factors could be fitted by a linear form compatible, in the studied range, with a monopole form, and extrapolated to the full kinematical region.

The form factors were used to evaluate the decay widths of the decays $D_s^+ \to f_0(980)\ell^+\nu_\ell$ and $D^+ \to f_0(980)\ell^+\bar{\nu}_\ell$ as a function of the mixing angle. Experimental data about these
decays would provide a direct estimate of the mixing angle. Using the same form factors to evaluate nonleptonic decays in the framework of the factorization approximation, it was not possible to explain all available experimental data with a fixed mixing angle. However, data seem to suggest that there is a sizable non-strange component in the $f_0(980)$ meson.

Acknowledgments: This work has been supported by CNPq and FAPESP (Brazil).

[1] E791 Collaboration, E.M. Aitala et al. Phys. Rev. Lett. 86, 770 (2001); ibid. 89, 121801 (2002); CLEO Collaboration, H. Muramatsu et al., Phys. Rev. Lett. 89 (2002) 251802; Ning Wu, “BES R measurements and $J/\Psi$ decays”, hep-ex/0104050 BELLE Collaboration, Abe, K. et al., hep-ex/0308043 contributed paper to the XXI International Symposium on Lepton and Photon Interactions at High Energies, Fermilab Aug 11-16, 2003.
[2] F.E. Close and N.A. Törnqvist, J. Phys. G28, R249 (2002).
[3] N.A. Törnqvist, Phys. Rev. Lett. 49, 624 (1982); Z. Phys. C68, 647 (1995).
[4] E. van Beveren, G. Rupp and M.D. Scadron, Phys. Lett. B495, 300 (2000).
[5] T. Schaefer and E. Shuryak, Rev. Mod. Phys. 70, 323 (1998).
[6] Particle Data Group, K.Hagiwara et al., Phys. Rev. D66, 010001 (2002).
[7] N.N. Achasov, Phys. Atom Nucl. 65, 573 (2002).
[8] H.-Y. Cheng, Phys. Rev. D67, 034024 (2003).
[9] V.V. Anisovich, L.G. Dakhno, V.A. Nikonov, hep-ph/0302137
[10] E791 Collaboration, E.M. Aitala et al. Phys. Rev. Lett. 86, 765 (2001).
[11] I. Bediaga and M. Nielsen, Phys. Rev. D68, 036001 (2003).
[12] G. Viehhauser for CLEO-C Collaboration Nucl.Instrum.Meth. A462, 146 (2001).
[13] M.A. Shifman, A.I. and Vainshtein and V.I. Zakharov, Nucl. Phys., B147, 385 (1979).
[14] F. de Fazio and M.R. Pennington, Phys. Lett. B521, 15 (2001).
[15] P. Colangelo and A. Khodjamirian, QCD sum rules: A modern perspective, in At the Frontier of Particle Physics, ed. M. Shifman, Singapore 2001, hep-ph/0010175
[16] Alpha Collaboration, A. Jütner and J. Rolf, hep-lat/0302016
[17] A. Deandrea et al., Phys. Lett. B502, 79 (2001).
[18] H.G. Dosch, E.M. Ferreira, F.S. Navarra and M. Nielsen, \textit{Phys. Rev.}, \textbf{D65}, 114002 (2002).

[19] P. Colangelo and F. de Fazio, \textit{Phys. Lett.} \textbf{B520}, 78 (2001).

[20] R.L. Jaffe, \textit{Phys. Rev.} \textbf{D15}, 267, 281 (1977); \textbf{D17}, 1444 (1978).

[21] M. Alford and R.L. Jaffe, \textit{Nucl. Phys.} \textbf{B578}, 367 (2000).