Hadronic-loop induced mass shifts in scalar heavy-light mesons

Feng-Kun Guo\textsuperscript{1}, Siegfried Krewald\textsuperscript{1}, and Ulf-G. Meißner\textsuperscript{1,2}

\textsuperscript{1}Institut für Kernphysik (Theorie), Forschungszentrum Jülich, D-52425 Jülich, Germany
\textsuperscript{2}Helmholtz-Institut für Strahlen- und Kernphysik (Theorie) Universität Bonn, Nußallee 14-16, D-53115 Bonn, Germany

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We calculate the mass shifts of heavy-light scalar mesons due to hadronic loops under the assumption that these vanish for the groundstate heavy-light mesons. The results show that the masses calculated in quark models can be reduced significantly. We stress that the mass alone is not a signal for a molecular interpretation. Both the resulting mass and the width suggest the observed $D^*_s(2317)$ state could be a dressed $c\bar{q}$ state. We give further predictions for the bottom scalar mesons which can be used to test the dressing mechanism.

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I. INTRODUCTION

The constituent quark model has been very successful in describing hadron spectroscopy. In recent years, some newly observed hadrons attracted much interest from both the experimental and theoretical community since these hadrons do not fit to the quark model predictions \cite{1}. For instance, the mass of the observed charm-strange scalar meson $D^*_s(2317)$ \cite{2} is $2317.3\pm0.6$ MeV \cite{3}, while the predictions from most quark models spread from 2400 MeV to 2500 MeV \cite{4,5}. The high mass predicted in constituent quark models is obtained through an orbital angular momentum excitation. The result from QCD sum rules in heavy quark effective theory gives a mass range $2.42\pm0.13$ GeV for the $D^*_s(2317)$ state which is consistent with, but the central value is 100 MeV higher than, the experimental value \cite{6}. Due to the fact that the mass of the $D^*_s(2317)$ is just below the $DK$ threshold at 2.36 GeV, a $DK$ molecular interpretation was proposed by Barnes et al. \cite{7} and some others \cite{8}. We want

\footnotesize{\*Electronic address: \texttt{f.k.guo@fz-juelich.de}}
to remark that a $DK$ bound state can be dynamically generated with a mass consistent with the observed mass of the $D^*_s(2317)$ in the framework of the heavy chiral unitary approach \cite{9,10}. Other exotic explanations were also proposed, such as tetraquark state \cite{11}, and $D\pi$ atom \cite{12}. Besides these exotic explanations, some authors tried to modify the quark model predictions. In \cite{13}, one loop corrections to the spin-dependent one-gluon exchange potential was considered, and the predicted mass of the $D^*_s(2317)$ is higher than the experimental value by only about 20 MeV. Another kind of modification is the mixing of the $c\bar{s}$ with the $cq\bar{q}\bar{q}$ tetraquark \cite{14}, or considering the coupling of the $c\bar{s}$ to hadronic channels, such as $DK$ \cite{15}. The prediction for the $D^*_s(2317)$ from the QCD sum rules can also be lowered to $2.331 \pm 0.016$ GeV considering the $DK$ continuum explicitly \cite{16}. Note that it is possible to distinguish a hadronic bound state from elementary hadrons, as done for the deuteron \cite{17} and for the light scalars $a_0(980)$ and $f_0(980)$ \cite{18}. All the results from Refs. \cite{15,16} indicate the importance of the strongly coupled hadronic channels on determining the mass of a hadron. For the charm-non-strange sector, both the Belle and FOCUS collaborations reported a scalar meson with a large width \cite{19,20}. Although the reported masses by different collaborations are not consistent with each other, the measurements are considered as the same charm scalar meson by the Particle Data Group (PDG), and the PDG average value of the mass is $2352 \pm 50$ MeV. The structure of this state has not been clear yet. In this Letter, we revisit the mass shifts of the heavy scalar mesons induced by the strongly coupled hadronic loops. For instance, the $D^*_s(2317)$, the $1^3P_0\ c\bar{s}$ state in quark model, can couple to the $D^+K^0$, $D^0K^+$ and $D_s\eta$ loops, see Fig. 1. The coupling constants of the $D^*_s(2317)$ to the three channels can be related by SU(3) symmetry. We shall use three different coupling types to study the mass shifts, called Model I, II and III in the following. In Model I, the coupling of the scalar heavy meson ($S$) to the heavy pseudoscalar meson ($P$) and the Goldstone boson ($\phi$) is assumed to be a constant. In Model II, the coupling is derived in the framework of heavy meson chiral perturbation theory (HM$\chi$PT) which combines the chiral expansion with the heavy quark expansion \cite{21,22} (for a review, see Ref. \cite{23}). In Model III, a chiral effective coupling is constructed disregarding the heavy quark expansion. Of course, such corrections due to hadronic loops are always model-dependent and are, in general, difficult to quantify, for a recent discussion see \cite{24} (For further discussion of incorporating hadronic loops in the quark model, see \cite{25}). This is why we consider three different models and also need to assume that for the ground state meson $Q\bar{q}$ (with $Q = c, b$ and $q = u, d, s$), the shift
due to the hadronic loops vanishes as it was done e.g. in the calculation of the mass shifts of charmonia in Ref. [26].

![Diagram](image)

FIG. 1: The relevant hadronic loops coupled to the $D_{s0}^*$.  

II. CHOICES OF COUPLING

A. Model I

First, the coupling of $S$ to $P, \phi$ is taken as a constant. The loop that integral enters the dressed propagator is

\[
G^I(s) = i \int \frac{d^4q}{(2\pi)^4} \frac{1}{(q^2 - m_1^2 + i\epsilon)[(p-q)^2 - m_2^2 + i\epsilon]},
\]

where $s = p^2$. The analytic expression is given by [27, 28]

\[
G^I(s) = \frac{1}{16\pi^2} \left\{ R - 1 + \ln \frac{m_2^2}{\mu^2} + \frac{m_1^2 - m_2^2 + s}{2s} \ln \frac{m_1^2}{m_2^2} + \frac{\sigma}{2s} [\ln(s - m_1^2 + m_2^2 + \sigma) - \ln(-s + m_1^2 - m_2^2 + \sigma) + \ln(s + m_1^2 - m_2^2 + \sigma) - \ln(-s - m_1^2 + m_2^2 + \sigma)] \right\},
\]

where $R = -[2/(4 - d) - \gamma_E + \ln(4\pi) + 1]$ will be set to zero in the calculations and $\gamma_E$ is Euler’s constant, $\mu$ is the scale of dimensional regularization, and $\sigma = \sqrt{s - (m_1 + m_2)^2][s - (m_1 - m_2)^2]$. Taking into account the $D^0K^+$, $D^+K^0$ and $D_s^+\eta$ channels, the shifted mass of the $D_{s0}^{*+}$ is given by the solution of the equation

\[
s - (\tilde{M}_{D_{s0}})^2 - g^2Re \left[ 2G^I_{DK}(s) + \frac{2}{3}G^I_{D_s\eta}(s) \right] = 0,
\]

where $\tilde{M}_{D_{s0}}$ denotes the bare mass of the $D_{s0}^*$. Note that we only consider the lowest possible intermediate states. In principle, all states with quantum numbers allowed by conservation
laws can contribute \[25\]. For instance, besides the channels considered here, the \(D^*K^*, D_s\eta'\) and \(D_s\rho\) can contribute either. But their threshold are at least 500 MeV higher than that of the DK, thus their contributions are expected to be suppressed. The corresponding equation for the non-strange charm meson \(D_0^{*+}\) is

\[
s - (M_{D_0})^2 - g^2 \text{Re} \left[ \frac{1}{2} G_{D^0}(s) + \frac{1}{6} G_{D^0}(s) + G_{D_s K}(s) \right] = 0,
\]

where four channels, \(D^+\pi^0\), \(D^0\pi^+\), \(D^+\eta\) and \(D_s^+K^0\), are taken into account. The coupling constant \(g\) has been calculated by using light-cone QCD sum rules in \[29\], \(g = 6.3 \pm 1.2\) GeV in the charm sector, \(g = 21 \pm 7\) GeV in the bottom sector. A more recent analysis considering the \(D_{s0}^*(2317)\) state as a conventional \(c\bar{s}\) meson gives the coupling constant for \(D_{s0}^*DK\) as \(5.9^{+1.7}_{-1.6}\) GeV \[30\], which is consistent with that given in Ref. \[29\]. In this Letter, we study the mass shifts of bare \(c\bar{q}\) (and \(b\bar{q}\)) mesons induced by hadronic loops. The values of coupling constants given in Ref. \[29\] will be taken because the masses of the scalar heavy mesons used therein are consistent with the quark-model expectation (no fitting to the mass of the \(D_{s0}^*(2317)\) was performed since the state had not been discovered yet) and hence correspond to the bare masses.

### B. Model II

In the HM\(\chi\)PT, the Lagrangian for the coupling of \(S\) to \(P\) and \(\phi\) to leading order is \[22, 23\]

\[
\mathcal{L} = i\hbar \langle S_\mu \gamma_\mu \gamma_5 \Phi^\mu_a H_a \rangle + h.c.
\]

\[
= \frac{i\sqrt{2}\hbar}{f_\pi} \left( D_{\mu\nu} \partial_\mu \Phi^\nu_a P^{a\dagger}_\mu - D_{\mu\nu}^\dagger \partial_\mu \Phi^\nu_a P^{a\dagger}_\mu \right) + \cdots,
\]

where the axial field

\[
A^\mu_{ba} = i\left( \xi\dagger \partial^\mu \xi - \xi \partial^\mu \xi\dagger \right)_{ba} = \frac{\partial^\mu \Phi_{ba}}{\sqrt{2}f_\pi} + \cdots
\]

contains the Goldstone bosons

\[
\Phi = \frac{1}{\sqrt{2}} \sum_{a=1}^{8} \lambda_a \phi_a = \begin{pmatrix}
\frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{6}} \eta & \pi^+ & K^+
\pi^- & -\frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{6}} \eta & K^0
K^- & \bar{K}^0 & -\frac{1}{\sqrt{3}} \eta
\end{pmatrix},
\]

the subscripts \(a, b\) represent to the light quark flavor \(u, d, s\), and \(f_\pi = 92.4\) MeV is the pion decay constant.

\[
H_a = \frac{1 + \gamma_5}{2} (P^{a\mu} \gamma_\mu - P_{a\gamma_5})
\]
represents the multiplet containing the pseudoscalar charm mesons, \( P = (D^0, D^+, D^+_s) \), and vector charm mesons, \( P^* = (D^{*0}, D^{*+}, D^{*+}_s) \), and \( \bar{H} = \gamma_0 H^\dagger \gamma_0 \).

\[
S_\alpha = \frac{1 - y}{2} (D^\mu_\alpha \gamma_\mu \gamma_5 - D_0 \alpha),
\]

represents the multiplet containing the scalar charm mesons, \( D_0 = (D^{*0}_0, D^{*0}_+ , D^{*0}_s) \), and axial charm mesons, \( D_1 = (D^1_0, D^1_+, D^1_{s1}) \). These field operators in Eqs. (8) and (9) have dimension 3/2 since they contain a factor \( \sqrt{M} \), where \( M \) is the mass of the corresponding meson, in their definition. Let \( p \) denotes the momentum of a pseudoscalar charm meson, e.g. \( D \), and \( k = p - M_D v \) its residual momentum. The propagator of \( D \) in HM\( \chi \)PT is

\[
\frac{i}{2(v \cdot k + \frac{3}{2} \Delta)},
\]

where \( \Delta = M_{D^*} - M_D \). The propagator of the strange charm meson \( D_s \) is

\[
\frac{i}{2(v \cdot k + \frac{3}{2} \Delta_s - \delta)},
\]

where \( \Delta_s = M_{D^*} - M_{D_s} \) and \( \delta = M_{D_{s0}} - M_{D_s} \). In the actual calculations, we take \( \delta = 0.1 \) GeV which is an approximate value of \( M_{D_{s0}} - M_{D_0} \) and \( M_{D^*_s} - M_{D^*} \). The propagator of the scalar charm mesons are similar with proper mass differences, i.e. \( \Delta_S = M_{D^*_1} - M_{D^*_0} \) for \( D^*_0 \), \( \Delta_{Ss} = M_{D_{s1}} - M_{D_{s0}} \) for \( D^*_s \), and the SU(3) breaking mass difference can be taken as \( \delta_S = 0.1 \) GeV, the same as \( \delta \). The coupling of the \( D^*_0 \) to the \( D^0 \) and \( \pi^+ \) to leading order is

\[
i\langle \pi^+(q) D^0(q')|\mathcal{L}|D^*_0(p) \rangle = -i \frac{\sqrt{2} h}{f_\pi} \sqrt{M_D \bar{M}_{D^*_0} v \cdot q}
\]

\[
= -i \frac{h}{\sqrt{2} f_\pi} \sqrt{M_D \bar{M}_{D^*_0}} \frac{(\bar{M}_{D^*_0})^2 - M_{D^*_0}^2 + m^2_\pi}{M_{D^*_0}}.
\]

The last equality holds for on-shell \( D^0 \) and \( \pi^+ \) mesons. The relation between \( h \) and \( g \) is

\[
h = -\frac{\sqrt{2} f_\pi \bar{M}_{D^*_0}}{\sqrt{M_D \bar{M}_{D^*_0}} \left( (\bar{M}_{D^*_0})^2 - M_{D^*_0}^2 + m^2_\pi \right)} g.
\]

Corresponding to \( q = 6.3 \pm 1.2 \) GeV, calculations via QCD sum rules give \( h = -0.44 \pm 0.09 \) for charm mesons. In Ref. [29], the mass difference between the scalar and the pseudoscalar heavy meson is taken to be 500 MeV. For bottom mesons, we have \( h = -0.52 \pm 0.18 \) [29]. The momentum of a scalar charm meson is \( p = \bar{M}_S v + k \), where \( \bar{M}_S \) is the bare mass of
the scalar charm meson and $k$ is the residual momentum. The momentum of the Goldstone boson in the loop is denoted by $q$. Then the residual momentum of the pseudoscalar charm meson, whose mass is $M_P$, in the loop should be $k' = p - q - M_P v = k + (\hat{M}_S - M_P) v - q$. The loop integral is

$$G^{\Pi}(v \cdot k) = \frac{i}{2} \int \frac{d^4q}{(2\pi)^4} \frac{(v \cdot q)^2}{(q^2 - m^2 + i\epsilon)[(v \cdot (k - q) + \hat{M}_S - M_P + \Delta + i\epsilon]}, \quad (14)$$

where $m$ is the mass of the Goldstone boson in the loop. The loop integral can be worked out as $[31, 32]

$$G^{\Pi}(v \cdot k) = \frac{m}{16\pi^2} J(0; \omega) \quad (15)$$

and

$$J(0; \omega) = \omega (R + \ln \frac{m^2}{\mu^2} - 1)$$

$$+ \begin{cases} 2\sqrt{\omega^2 - m^2} \cosh^{-1}(\frac{\omega}{m}) - 2\pi i \sqrt{\omega^2 - m^2}, & \omega > m \\ 2\sqrt{m^2 - \omega^2} \cos^{-1}(\frac{\omega}{m}), & \omega^2 < m^2 \\ -2\sqrt{\omega^2 - m^2} \cosh^{-1}(\frac{\omega}{m}), & \omega < -m \end{cases} \quad (16)$$

where $\omega = v \cdot k + \hat{M}_S - M_P + 3\Delta/4$. The dressed propagator for $D^{*+}_{s0}$ is

$$i \frac{2 (v \cdot k + \frac{3}{4} \Delta_{Ss} - \delta_S) - \frac{2\hbar^2}{f_{\pi}^2} \hat{M}_{D^{*}_{s0}} \text{Re}[2M_D G_{DK}^{\Pi}(v \cdot k) + \frac{2}{3} M_{D_s} G_{D_s\eta}^{\Pi}(v \cdot k)]]}{2 (v \cdot k + \frac{3}{4} \Delta_{Ss} - \delta_S) - \frac{2\hbar^2}{f_{\pi}^2} \hat{M}_{D^{*}_{s0}} \text{Re}[2M_D G_{DK}^{\Pi}(v \cdot k) + \frac{2}{3} M_{D_s} G_{D_s\eta}^{\Pi}(v \cdot k)]]} \quad (17)$$

For the mass difference $\Delta_{Ss}$ the physical values is taken in the propagator which is correct to the order considered. Thus substituting $v \cdot k$ by $v \cdot p - \hat{M}_{D^{*}_{s0}}$, the shifted mass can be given the value of $v \cdot p$ which is the solution of the following equation $[31]$

$$2 \left( v \cdot p - \hat{M}_{D^{*}_{s0}} \right) - \frac{2\hbar^2}{f_{\pi}^2} \hat{M}_{D^{*}_{s0}} \text{Re} \left[ 2M_D G_{DK}^{\Pi}(v \cdot k) + \frac{2}{3} M_{D_s} G_{D_s\eta}^{\Pi}(v \cdot k) \right] = 0 \quad (18)$$

The corresponding propagator of the $D^{*0}_{0}$ is

$$i \frac{2 (v \cdot k + \frac{3}{4} \Delta_S) - \frac{2\hbar^2}{f_{\pi}^2} \hat{M}_{D^0} \text{Re}[\frac{3}{2} M_D G_{D\pi}^{\Pi}(v \cdot k) + \frac{1}{6} M_D G_{\eta}^{\Pi}(v \cdot k) + M_{D_s} G_{D_sK}^{\Pi}(v \cdot k)]]}{2 (v \cdot k + \frac{3}{4} \Delta_S) - \frac{2\hbar^2}{f_{\pi}^2} \hat{M}_{D^0} \text{Re}[\frac{3}{2} M_D G_{D\pi}^{\Pi}(v \cdot k) + \frac{1}{6} M_D G_{\eta}^{\Pi}(v \cdot k) + M_{D_s} G_{D_sK}^{\Pi}(v \cdot k)]]} \quad (19)$$

From Eq. (15), the mass shift vanishes in the chiral limit, similar to the mass shift of the nucleon due to $N\pi$ loop in heavy baryon $\chi$PT $[31]$, in contrast to that in Model I and III, see below.
C. Model III

Disregarding the heavy quark expansion, we can directly construct an effective chiral Lagrangian describing the coupling of a scalar charm meson with a pseudoscalar charm meson and a Goldstone boson. The Lagrangian is

\[ \mathcal{L} = h' D_{0b} A_{ba}^\mu \partial_\mu P_a^1 + \text{h.c.} \]  

(20)

Note that the field operators of the heavy mesons in Eq. (20) have dimension 1, different from those in Eqs. (8,9). This Lagrangian drives the coupling to be of the type

\[ i \langle \pi^+(q) D^0(q') | \mathcal{L} | D_0^{*+}(p) \rangle = i \frac{h'}{\sqrt{2} f_\pi} q \cdot q' = i \frac{h'}{2 \sqrt{2} f_\pi} \left( (\tilde{M}_{D_0^*})^2 - M_D^2 - m_\pi^2 \right). \]

(21)

The second equality is fulfilled only for on-shell $D^0$ and $\pi^+$ mesons. The relation between $h'$ and $g$ is

\[ h' = \frac{2 \sqrt{2} f_\pi}{(\tilde{M}_{D_0^*})^2 - M_D^2 - m_\pi^2} g. \]

(22)

For extracting the value of $h'$, we take $\tilde{M}_{D_0^*} - M_D = 500$ MeV following Ref. [29]. Then we have $h' = 0.78 \pm 0.15$ for charm mesons. Similarly, the coupling constant for bottom mesons can be obtained as $h' = 1.00 \pm 0.33$. The dressed propagators of the $D^{*+}_{s0}$ and $D^{*+}_0$ are

\[ i \left( s - (\tilde{M}_{D^{*}_{s0}})^2 - \frac{h'^2}{2 f_\pi^2} \text{Re} \left[ 2 G^{\text{III}}_{DK}(s) + \frac{2}{3} G^{\text{III}}_{D\eta}(s) \right] \right), \]

(23)

and

\[ i \left( s - (\tilde{M}_{D^{*}_0})^2 - \frac{h'^2}{2 f_\pi^2} \text{Re} \left[ \frac{3}{2} G^{\text{III}}_{D\pi}(s) + \frac{1}{6} G^{\text{III}}_{D\eta}(s) + G^{\text{III}}_{D\eta_K}(s) \right] \right), \]

(24)

respectively. The physical masses of the charm scalar mesons can be obtained by setting the denominators of the propagators to zero. The loop integral in the dressed propagators is

\[ G^{\text{III}}(s) = i \int \frac{d^4q}{(2\pi)^4} \frac{[q \cdot (p-q)]^2}{(q^2 - m_1^2 + i\epsilon)(p-q)^2 - m_2^2 + i\epsilon)}. \]

(25)
The analytic expression can be worked out as

\[
G^{III}(s) = \frac{1}{16\pi^2} \left\{ \left[ m_1^4 + m_1^2 m_2^2 + m_2^4 - \frac{3}{4}(m_1^2 + m_2^2)s + \frac{s^2}{4} \right] \left( R + \ln \frac{m_2^2}{\mu^2} \right) \right.
\]

\[
- \frac{(m_1^2 + m_2^2 - s)^2}{4} + \frac{m_1^2 - m_2^2 + s}{8s} \left[ \sigma^2 + 2(m_1^2 + m_2^2) \right] - 2m_1^2\sigma^2 \ln \frac{m_1^2}{m_2^2} \right.
\]

\[
+ \frac{\sigma(m_1^2 + m_2^2 - s)^2}{8s} \left[ \ln(s - m_1^2 + m_2^2 + \sigma) - \ln(-s + m_1^2 - m_2^2 + \sigma) \right]
\]

\[
+ \ln(s + m_1^2 - m_2^2 + \sigma) - \ln(-s + m_1^2 + m_2^2 + \sigma) \right\} .
\]

(26)

III. RESULTS

In general, the effect of a hadronic loop coupling to a bare state is to pull its bare mass down to the physical one (if the physical mass is above the threshold of the channel, the imaginary part of the loop will contribute to the width of the state). One can expect that the mass of a charm meson cannot be pulled down to below the mass of the charm quark \( m_c \). So we can assume that at some point, called the subtraction point, above \( m_c \), the contribution from any hadronic loop vanishes. On the other hand, the subtraction point cannot be too high since a state close to the threshold of a strong decay channel would have a strong coupling to the channel, and hence its mass would be affected. A similar idea has been taken to study the mass shifts of charmonia \[26\]. After choosing a specific subtraction point, the renormalization scale \( \mu \) considered as a parameter can be determined from this assumption. In the loop function in Model I, Eq. (2), the coefficient of the chiral logarithm \( \ln \left( m_2^2/\mu^2 \right) \) is a constant \( 1/(16\pi^2) \); in the loop functions in Model II and III, Eq. (15) and Eq. (26), the coefficients of the same logarithm are momentum-dependent. In Model III, the coefficient

\[
C_{III} = \frac{1}{16\pi^2} \left[ m_1^4 + m_1^2 m_2^2 + m_2^4 - \frac{3}{4}(m_1^2 + m_2^2)s + \frac{1}{4}s^2 \right]
\]

is always positive, and it changes slowly with respect to \( \sqrt{s} \) below 3 GeV for the \( DK \) loop, see Fig. 2(a). So in Model I and III, the dependence of \( \mu \) on the choice of subtraction point is small. For instance, the values of \( \mu \) determined when the subtraction point is chosen at \( \sqrt{s} = M_D = 1.87 \text{ GeV} \) and \( \sqrt{s} = m_c = 1.35 \text{ GeV} \) are listed in Table I. One can see when the subtraction point is changed from \( M_D \) to \( m_c \), the resulting values of \( \mu \) in the \( DK \) loop changes slightly, and all the values are not far away from \( M_D \). In the following, we shall
take \( M_D \) as the subtraction point for Model I and III. However, in Model II, the coefficient \( C_{II} = m_2^2 \omega/(16\pi^2) \) is proportional to \( v \cdot p \), and changes its sign at just below \( M_D \). In Fig. 2(b), we show \( C_{II} \) as a function of \( v \cdot p \) where \( M_P = M_D \) and \( m = m_K \) are used. So the resulting value of \( \mu \) depends strongly on the choice of subtraction point. The values of \( \mu \) in Model II are also given in Table I. It seems that these values are unphysical. Because we only take the leading order in heavy quark expansion to describe the coupling in Model II, the unphysical value of \( \mu \) might indicate that higher order contributions, which will change the behavior of \( C_{II} \), are important. To avoid this problem, a commonly used value \( \mu = 1 \) GeV in HM\( \chi \)PT is taken in Model II. Certainly, the \( \mu \) dependence should be absorbed by counterterms at the next order. The same method can be used directly for the scalar bottom mesons.

![Graph](image)

**FIG. 2:** The coefficients of the logarithm \( \ln (m_2^2/\mu^2) \) in the loop functions in Model III (a) and Model II (b). \( m_1 = M_D \) and \( m_2 = m_K \) are taken.

**TABLE I:** The values of the renormalization scale \( \mu \) in the \(DK\) loop determined when the subtraction point being chosen at \( M_D \) and \( m_c \). All units are in GeV.

| Subtraction point | Model I | Model II | Model III |
|------------------|---------|----------|-----------|
| \( M_D = 1.87 \) GeV | 1.58    | 1095.3   | 1.85      |
| \( m_c = 1.35 \) GeV | 1.84    | 0.20     | 1.82      |

with no more free parameters. In Model I and III, we use the same assumption that the contribution from any hadronic loop vanish at the mass of the lowest heavy flavor meson. That is, in the bottom case the subtraction point is chosen to be \( M_B = 5.28 \) GeV. The bare masses of the scalar heavy mesons are taken from the popular Godfrey-Isgur quark model.
which can describe the meson spectroscopy, especially the low-lying states very well [4], i.e. \( M_{D^*_0} = 2.48 \text{ GeV}, M_{D^*_0} = 2.40 \text{ GeV}, M_{B^*_0} = 5.83 \text{ GeV} \) and \( M_{B^*_0} = 5.76 \text{ GeV} \).

The resulting masses of all the four mesons are listed in Table II. The results from Model I are consistent with those in Model III, which show that the bare mass of the \( D^*_0 \) can be pulled down significantly to the region close to the mass of the \( D^*_0(2317) \) state. If we choose a bare mass from another quark model, the obtained mass can even be consistent with the experimental value. In Model I, a constant coupling is taken which would violate Goldstone’s theorem because the \( \pi, K \) and \( \eta \) are Goldstone bosons. Model II has a large \( \mu \) dependence, which makes it not preferable for a phenomenological analysis. The Lagrangian for Model III is constructed from chiral symmetry, and the \( \mu \) dependence is really small (for instance, the resulting mass of the \( D^*_0 \) would change to 2.37-2.41 GeV if we choose \( m_c = 1.35 \text{ GeV} \) as the subtraction point, and the change is no more than 10 MeV). Therefore, we choose the results in Model III to give further predictions. The absolute mass shifts in Model III are listed in the last column in Table II. The mass of the \( D^*_0 \) is consistent with the experimental data \( 2352 \pm 50 \text{ MeV} \) [3]. One strong decay channel \( D\pi \) is open for the state, and the decay width of this channel should give the dominant contribution to the width of the \( D^*_0 \). Then the width of the \( D^*_0 \) can be obtained from

\[
\Gamma_{D^*_0} = \frac{\hbar^2}{2f_{\pi}^2 M_{D^*_0}} \frac{3}{2} \text{Im} G^{\text{III}}_{D\pi}(M_{D^*_0}^2). \tag{27}
\]

\[1\] If the parameters of a quark model, e.g. the constituent quark masses and strong coupling constant, were determined by fitting to the whole hadron mass spectrum, loop effects would be incorporated to some extent although in an unclear way. But in reality, the quark model parameters were fitted to mostly the low-lying states. Especially in the Godfrey-Isgur quark model the physical spectrum used in fit does not contain the heavy scalar mesons studied here.
Using the mass $M_{D^*_0}$ from Table III we obtain

$$\Gamma_{D^*_0} = 99 - 167 \text{ MeV}.$$  \hspace{1cm} (28)

The result is roughly consistent with the experimental width for the $D^*_0$, which was reported as $276 \pm 21 \pm 63 \text{ MeV}$ by the Belle Collaboration [19] and $240 \pm 55 \pm 59 \text{ MeV}$ by the FOCUS Collaboration [20]. Both the mass and the width suggest the observed $D^*_0$ state can be the dressed $c\bar{q}$ state. Note that in Ref. [10], two molecular states were predicted, and both of them are not consistent with the data. Certainly, the uncertainty of the data is large so far, and more precise measurements are highly desirable. Another noticeable result is that the mass shifts in the bottom sector are about twice of those in the charm sector, as can be seen from the last column in Table III. Based on heavy quark symmetry, and assuming the $D^*_{s0}(2317)$ is a $DK$ bound state, the mass of the $B\bar{K}$ bound state was predicted as $5733 \text{ MeV}$ [33] which was confirmed by a dynamical calculation [10], larger than the mass region obtained here by dressing the $b\bar{q}$ state. That means if a $B^*_{s0}$ state with a mass which is much smaller than $5733 \text{ MeV}$ were found, it would probably be a dressed $b\bar{q}$ state rather than a $B\bar{K}$ bound state, or the $D^*_{s0}(2317)$ would not be a $DK$ bound state. Similar to that of the $D^*_0$, the width of the $B^*_0$ can be estimated as

$$\Gamma_{B^*_0} = 62 - 100 \text{ MeV}.$$  \hspace{1cm} (29)

IV. SUMMARY

The $D^*_{s0}(2317)$ is considered as a $DK$ molecular state by many authors because in the quark model its mass simply comes out too high. However, hadronic loops can pull down its mass. To make the calculation of such an effect quantitative, we assume that the hadronic-loop induced mass shifts of a hadron vanish at some point. The point is chosen as $\sqrt{s} = m_D$ for the charm sector and $m_B$ for the bottom sector. Then we calculate the mass shifts of the heavy scalar mesons by using three different types of coupling. The input bare masses are taken from the Godfrey-Isgur quark model [4]. The mass of the $D^*_{s0}$ state is lowered significantly, and it can even be pulled down to $2317 \text{ MeV}$ if we adjust the bare mass in the predicted region from different quark models. That means the simple argument against the quark model for too high masses is not valid. What the quark model predicts is just the bare masses of hadrons. In order to compare with the experimental spectroscopy, the
bare hadron masses need to be dressed \cite{24}. Note, however, that such a dressing is always
model-dependent and must be considered in different approaches, as done here. Because
chiral symmetry is fulfilled and the $\mu$ dependence is small, we choose Model III to give
further predictions. The results of both mass and width show that the observed $D_0^*$ could
be a dressed $c\bar{q}$ state. We also give predictions for the dressed bottom scalar mesons. The
mass of the $B_{s0}^*$ is smaller than that of the $B\bar{K}$ bound state which was obtained assuming
the $D_{s0}^*(2317)$ be a $DK$ bound state \cite{33}. Precise experimental data from $B$ factories are
highly desirable to test the dressing mechanism. We would like to stress that a consistent
treatment of the mass and width is required. The difference from the molecular state will
presumably be revealed in the decay pattern with into various channels, this is a further
step to be investigated.

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