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

Feng-Kun Guo[1],* Siegfried Krewald[1], and Ulf-G. Meißner[1,2]

[1] *Institut für Kernphysik (Theorie), Forschungszentrum Jülich, D-52425 Jülich, Germany*

[2] *Helmholtz-Institut für Strahlen- und Kernphysik (Theorie) Universität Bonn, Nußallee 14-16, D-53115 Bonn, Germany*

(Dated: February 24, 2013)

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_0^* 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.

PACS numbers: 12.39.-x,12.40.Yx,14.40.Lb

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 [1]. For instance, the mass of the observed charm-strange scalar meson D_{s0}^* [2] is 2317.3 ± 0.6 MeV [3], while the predictions from most quark models spread from 2400 MeV to 2500 MeV [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 ± 0.13 GeV for the D_{s0}^* $(c\bar{s})$ state which is consistent with, but the central value is 100 MeV higher than, the experimental value [6]. Due to the fact that the mass of the $D_{s0}^*(2317)$ is just below the DK threshold at 2.36 GeV, a DK molecular interpretation was proposed by Barnes et al. [7] and some others [8]. We want

*Electronic address: 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 $D_{s0}^*(2317)$ in the framework of the heavy chiral unitary approach [9, 10]. Other exotic explanations were also proposed, such as tetraquark state [11], and $D\pi$ atom [12]. Besides these exotic explanations, some authors tried to modify the quark model predictions. In [13], one loop corrections to the spin-dependent one-gluon exchange potential was considered, and the predicted mass of the D_{s0}^* 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 $cq\bar{s}\bar{q}$ tetraquark [14], or considering the coupling of the $c\bar{s}$ to hadronic channels, such as DK [15]. The prediction for the D_{s0}^* from the QCD sum rules can also be lowered to 2.331 ± 0.016 GeV considering the DK continuum explicitly [16]. Note that it is possible to distinguish a hadronic bound state from elementary hadrons, as done for the deuteron [17] and for the light scalars $a_0(980)$ and $f_0(980)$ [18]. All the results from Refs. [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 [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 ± 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_{s0}^{*+}, 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_{s0}^{*+} 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 (ϕ) is assumed to be a constant. In Model II, the coupling is derived in the framework of heavy meson chiral perturbation theory (HMχPT) which combines the chiral expansion with the heavy quark expansion [21, 22] (for a review, see Ref. [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 [24] (For further discussion of incorporating hadronic loops in the quark model, see [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].

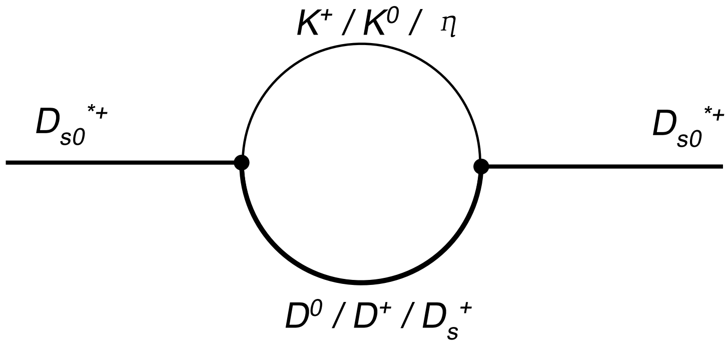

FIG. 1: The relevant hadronic loops coupled to the D_{s0}^*.

II. CHOICES OF COUPLING

A. Model I

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

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

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

$$G^{\mathrm{I}}(s) = \frac{1}{16\pi^2} \Big\{ R - 1 + \ln\frac{m_2^2}{\mu^2} + \frac{m_1^2 - m_2^2 + s}{2s} \ln\frac{m_1^2}{m_2^2} + \frac{\sigma}{2s} \big[\ln(s - m_1^2 + m_2^2 + \sigma) - \ln(-s + m_1^2 - m_2^2 + \sigma) + \ln(s + m_1^2 - m_2^2 + \sigma) - \ln(-s - m_1^2 + m_2^2 + \sigma) \big] \Big\}, \tag{2}$$

where $R = -[2/(4-d) - \gamma_E + \ln(4\pi) + 1]$ will be set to zero in the calculations and γ_E is Euler's constant, μ 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

$$s - (\mathring{M}_{D_{s0}^*})^2 - g^2 \mathrm{Re}\left[2G^{\mathrm{I}}_{DK}(s) + \frac{2}{3} G^{\mathrm{I}}_{D_s\eta}(s) \right] = 0, \tag{3}$$

where $\mathring{M}_{D_{s0}^*}$ 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

$$s - (\mathring{M}_{D_0^*})^2 - g^2\mathrm{Re}\left[\frac{3}{2}G^{\mathrm{I}}_{D\pi}(s) + \frac{1}{6}G^{\mathrm{I}}_{D\eta}(s) + G^{\mathrm{I}}_{D_sK}(s)\right] = 0, \tag{4}$$

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], $g = 6.3\pm1.2$ GeV in the charm sector, $g = 21\pm7$ 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 HMχPT, the Lagrangian for the coupling of S to P and ϕ to leading order is [22, 23]

$$\begin{aligned}\mathcal{L} &= ih\langle S_b\gamma_\mu\gamma_5 A^\mu_{ba}\bar{H}_a\rangle + h.c. \\ &= \frac{i\sqrt{2}h}{f_\pi}\left(D_{0b}v^\mu\partial_\mu\Phi_{ba}P_a^\dagger - D^\nu_{1b}v^\mu\partial_\mu\Phi_{ba}P^{*\dagger}_{a\nu}\right) + \cdots,\end{aligned} \tag{5}$$

where the axial field

$$A^\mu_{ba} = \frac{i}{2}(\xi^\dagger\partial^\mu\xi - \xi\partial^\mu\xi^\dagger)_{ba} = -\frac{\partial^\mu\Phi_{ba}}{\sqrt{2}f_\pi} + \cdots \tag{6}$$

contains the Goldstone bosons

$$\Phi = \frac{1}{\sqrt{2}}\sum_{a=1}^{8}\lambda_a\phi_a = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{1}{\sqrt{3}}\eta \end{pmatrix}, \tag{7}$$

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

$$H_a = \frac{1+\not{v}}{2}\left(P_a^{*\mu}\gamma_\mu - P_a\gamma_5\right) \tag{8}$$

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 $\bar{H} = \gamma_0 H^\dagger \gamma_0$.

$$S_a = \frac{1+ \not{v}}{2}\left(D^\mu_{1a}\gamma_\mu\gamma_5 - D_{0a}\right), \tag{9}$$

represents the multiplet containing the scalar charm mesons, $D_0 = (D_0^{*0}, D_0^{*+}, D_{s0}^{*+})$, and axial charm mesons, $D_1 = (D_1^0, D_1^+, D_{s1}^+)$. 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 HMχPT is [21]

$$\frac{i}{2(v\cdot k + \frac{3}{4}\Delta)}, \tag{10}$$

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

$$\frac{i}{2(v\cdot k + \frac{3}{4}\Delta_s - \delta)}, \tag{11}$$

where $\Delta_s = M_{D^*} - M_D$ and $\delta = M_{D_s} - M_D$. In the actual calculations, we take $\delta = 0.1$ GeV which is an approximate value of $M_{D_s} - M_D$ 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_{s0}^*, and the SU(3) breaking mass difference can be taken as $\delta_S = 0.1$ GeV, the same as δ. The coupling of the D_0^{*+} to the D^0 and π^+ to leading order is

$$\begin{aligned} i\langle \pi^+(q) D^0(q')|\mathcal{L}|D_0^{*+}(p)\rangle &= -i\frac{\sqrt{2}h}{f_\pi}\sqrt{M_D \mathring{M}_{D_0^*}}\, v\cdot q \\ &= -i\frac{h}{\sqrt{2}f_\pi}\sqrt{M_D \mathring{M}_{D_0^*}}\frac{(\mathring{M}_{D_0^*})^2 - M_D^2 + m_\pi^2}{\mathring{M}_{D_0^*}}. \end{aligned} \tag{12}$$

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

$$h = -\frac{\sqrt{2} f_\pi \mathring{M}_{D_0^*}}{\sqrt{M_D \mathring{M}_{D_0^*}}\left((\mathring{M}_{D_0^*})^2 - M_D^2 + m_\pi^2\right)} g. \tag{13}$$

Corresponding to $g = 6.3 \pm 1.2$ GeV, calculations via QCD sum rules give $h = -0.44 \pm 0.09$ for charm mesons [29]. 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 $p = \mathring{M}_S v + k$, where $\mathring{M}_S$ 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 $k' = p - q - M_P v = k + (\overset{\circ}{M}_S - M_P)v - q$. The loop integral is

$$G^{\mathrm{II}}(v \cdot k) = \frac{i}{2} \int \frac{d^4q}{(2\pi)^4} \frac{(v \cdot q)^2}{(q^2 - m^2 + i\varepsilon)[(v \cdot (k - q) + \overset{\circ}{M}_S - M_P + \Delta + i\varepsilon]}, \tag{14}$$

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

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

and

$$J(0;\omega) = \omega(R + \ln \frac{m^2}{\mu^2} - 1) + \begin{cases} 2\sqrt{\omega^2 - m^2}\cosh^{-1}(\frac{\omega}{m}) - 2\pi i \sqrt{\omega^2 - m^2}, & \omega > m \\ 2\sqrt{m^2 - \omega^2}\cos^{-1}(-\frac{\omega}{m}), & \omega^2 < m^2 \\ -2\sqrt{\omega^2 - m^2}\cosh^{-1}(-\frac{\omega}{m}), & \omega < -m \end{cases} \tag{16}$$

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

$$\frac{i}{2\left(v \cdot k + \frac{3}{4}\Delta_{Ss} - \delta_S\right) - \frac{2h^2}{f_\pi^2} \overset{\circ}{M}_{D_{s0}^*} \mathrm{Re}\left[2M_D G_{DK}^{\mathrm{II}}(v \cdot k) + \frac{2}{3} M_{D_s} G_{D_s\eta}^{\mathrm{II}}(v \cdot k)\right]}. \tag{17}$$

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

$$2\left(v \cdot p - \overset{\circ}{M}_{D_{s0}^*}\right) - \frac{2h^2}{f_\pi^2} \overset{\circ}{M}_{D_{s0}^*} \mathrm{Re}\left[2M_D G_{DK}^{\mathrm{II}}(v \cdot k) + \frac{2}{3} M_{D_s} G_{D_s\eta}^{\mathrm{II}}(v \cdot k)\right] = 0. \tag{18}$$

The corresponding propagator of the D_0^* is

$$\frac{i}{2\left(v \cdot k + \frac{3}{4}\Delta_S\right) - \frac{2h^2}{f_\pi^2} \overset{\circ}{M}_{D_0^*} \mathrm{Re}\left[\frac{3}{2} M_D G_{D\pi}^{\mathrm{II}}(v \cdot k) + \frac{1}{6} M_D G_{D\eta}^{\mathrm{II}}(v \cdot k) + M_{D_s} G_{D_s K}^{\mathrm{II}}(v \cdot k)\right]}. \tag{19}$$

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 χPT [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

$$\mathcal{L} = h' D_{0b} A^{\mu}_{ba} \partial_\mu P^\dagger_a + \text{h.c}\,. \tag{20}$$

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

$$\begin{aligned} i\langle \pi^+(q) D^0(q')|\mathcal{L}|D_0^{*+}(p)\rangle &= i\frac{h'}{\sqrt{2}f_\pi} q\cdot q' \\ &= i\frac{h'}{2\sqrt{2}f_\pi}\left((\overset{\circ}{M}_{D_0^*})^2 - M_D^2 - m_\pi^2\right). \end{aligned} \tag{21}$$

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

$$h' = \frac{2\sqrt{2}f_\pi}{(\overset{\circ}{M}_{D_0^*})^2 - M_D^2 - m_\pi^2} g. \tag{22}$$

For extracting the value of h', we take $\overset{\circ}{M}_{D_0^*} - M_D = 500$ 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

$$\frac{i}{s - (\overset{\circ}{M}_{D_{s0}^*})^2 - \frac{h'^2}{2f_\pi^2}\text{Re}\left[2G^{\text{III}}_{DK}(s) + \frac{2}{3}G^{\text{III}}_{D_s\eta}(s)\right]}, \tag{23}$$

and

$$\frac{i}{s - (\overset{\circ}{M}_{D_0^*})^2 - \frac{h'^2}{2f_\pi^2}\text{Re}\left[\frac{3}{2}G^{\text{III}}_{D\pi}(s) + \frac{1}{6}G^{\text{III}}_{D\eta}(s) + G^{\text{III}}_{D_sK}(s)\right]}, \tag{24}$$

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

$$G^{\text{III}}(s) = i\int\frac{d^4q}{(2\pi)^4}\frac{[q\cdot(p-q)]^2}{(q^2 - m_1^2 + i\epsilon)[(p-q)^2 - m_2^2 + i\epsilon]}. \tag{25}$$

The analytic expression can be worked out as

$$G^{\rm III}(s) = \frac{1}{16\pi^2}\left\{\left[m_1^4+m_1^2m_2^2+m_2^4-\frac{3}{4}(m_1^2+m_2^2)s+\frac{s^2}{4}\right]\left(R+\ln\frac{m_2^2}{\mu^2}\right)\right.$$
$$-\frac{(m_1^2+m_2^2-s)^2}{4}+\frac{(m_1^2-m_2^2+s)\left[\sigma^2+2m_1^2(m_1^2+m_2^2)\right]-2m_1^2\sigma^2}{8s}\ln\frac{m_1^2}{m_2^2}$$
$$+\frac{\sigma(m_1^2+m_2^2-s)^2}{8s}\left[\ln(s-m_1^2+m_2^2+\sigma)-\ln(-s+m_1^2-m_2^2+\sigma)\right.$$
$$\left.\left.+\ln(s+m_1^2-m_2^2+\sigma)-\ln(-s-m_1^2+m_2^2+\sigma)\right]\right\}\ . \tag{26}$$

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 can not 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 μ 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 $\ln(m_2^2/\mu^2)$ 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

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

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 μ on the choice of subtraction point is small. For instance, the values of μ determined when the subtraction point is chosen at $\sqrt{s}=M_D=1.87$ GeV and $\sqrt{s}=m_c=1.35$ 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 μ 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 $C_{\rm II} = m^2\omega/(16\pi^2)$ is proportional to $v \cdot p$, and changes its sign at just below M_D. In Fig. 2(b), we show $C_{\rm II}$ as a function of $v \cdot p$ where $M_P = M_D$ and $m = m_K$ are used. So the resulting value of μ depends strongly on the choice of subtraction point. The values of μ 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 μ might indicate that higher order contributions, which will change the behavior of $C_{\rm II}$, are important. To avoid this problem, a commonly used value $\mu = 1$ GeV in HMχPT is taken in Model II. Certainly, the μ dependence should be absorbed by counterterms at the next order. The same method can be used directly for the scalar bottom mesons

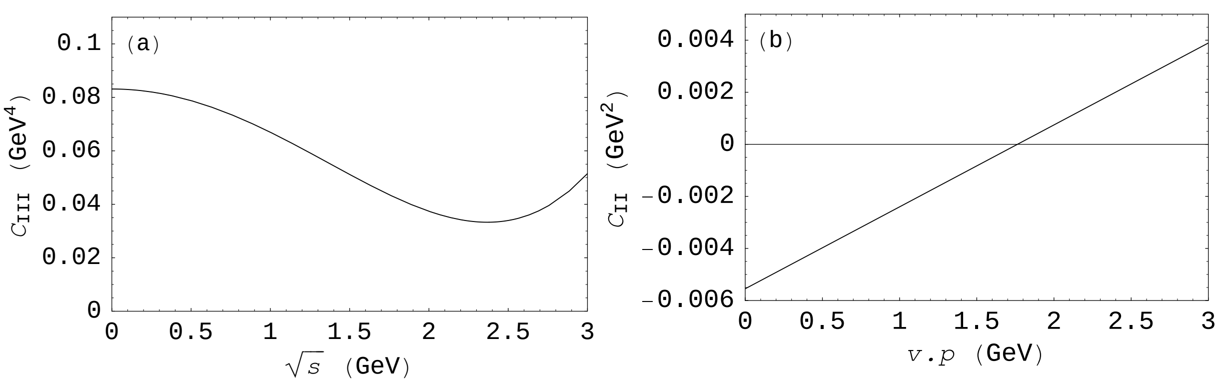

FIG. 2: The coefficients of the logarithm $\ln(m_2^2/\mu^2)$ in the loop functions in Model III (a) and Model II (b). $m_1 = M_D$ and $m_2 = m_K$ are taken.

TABLE I: The values of the renormalization scale μ 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],[1] i.e. $\mathring{M}_{D^*_{s0}} = 2.48$ GeV, $\mathring{M}_{D^*_0} = 2.40$ GeV, $\mathring{M}_{B^*_{s0}} = 5.83$ GeV and $\mathring{M}_{B^*_0} = 5.76$ GeV.

The resulting masses of all the four mesons are listed in Table II. The results from Model I

TABLE II: The resulted masses of the lowest scalar heavy mesons from dressing in different models.

	Bare Mass	Model I	Model II	Model III	$\Delta M_{\rm III}$ (GeV)
$M_{D^*_{s0}}$ (GeV)	2.48	2.33-2.39	2.26-2.35	2.36-2.40	0.08-0.12
$M_{D^*_0}$ (GeV)	2.40	2.30-2.35	2.39-2.40	2.30-2.34	0.06-0.10
$M_{B^*_{s0}}$ (GeV)	5.83	5.58-5.72	5.73-5.73	5.62-5.70	0.13-0.21
$M_{B^*_0}$ (GeV)	5.76	5.49-5.67	5.85-5.86	5.55-5.64	0.12-0.21

are consistent with those in Model III, which show that the bare mass of the D^*_{s0} can be pulled down significantly to the region close to the mass of the $D^*_{s0}(2317)$ 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 π, K and η are Goldstone bosons. Model II has a large μ dependence, which makes it not preferable for a phenomenological analysis. The Lagrangian for Model III is constructed from chiral symmetry, and the μ dependence is really small (for instance, the resulting mass of the D^*_{s0} would change to 2.37-2.41 GeV if we choose $m_c = 1.35$ 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 ± 50 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

$$\Gamma_{D^*_0} = \frac{h'^2}{2f_\pi^2 M_{D^*_0}} \frac{3}{2} {\rm Im} G^{\rm III}_{D\pi}(M^2_{D^*_0}). \tag{27}$$

[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 II, we obtain

$$\Gamma_{D_0^*} = \ 99 - 167 \text{ MeV} . \tag{28}$$

The result is roughly consistent with the experimental width for the D_0^*, which was reported as $276 \pm 21 \pm 63$ MeV by the Belle Collaboration [19] and $240 \pm 55 \pm 59$ 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 II. 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 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 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

$$\Gamma_{B_0^*} = \ 62 - 100 \text{ MeV} . \tag{29}$$

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 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 [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 μ 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 [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.

Acknowledgments

We would like to thank Ch. Hanhart for useful discussions. The work is partially supported by the Helmholtz Association through funds provided to the virtual institute "Spin and strong QCD"(VH-VI-231) and by the EU Integrated Infrastructure Initiative Hadron Physics Project under contract number RII3-CT-2004-506078.

[1] E. S. Swanson, Phys. Rept. **429** (2006) 243 [arXiv:hep-ph/0601110].

[2] B. Aubert *et al.* [BABAR Collaboration], Phys. Rev. Lett. **90** (2003) 242001 [arXiv:hep-ex/0304021].

[3] W. M. Yao *et al.* [Particle Data Group], J. Phys. G **33** (2006) 1.

[4] S. Godfrey and N. Isgur, Phys. Rev. D **32** (1985) 189; S. Godfrey and R. Kokoski, Phys. Rev. D **43** (1991) 1679.

[5] J. Zeng, J. W. Van Orden and W. Roberts, Phys. Rev. D **52** (1995) 5229; [arXiv:hep-ph/9412269]; S. N. Gupta and J. M. Johnson, Phys. Rev. D **51** (1995) 168 [arXiv:hep-ph/9409432]; D. Ebert, V. O. Galkin and R. N. Faustov, Phys. Rev. D **57** (1998) 5663 [Erratum-ibid. D **59** (1999) 019902] [arXiv:hep-ph/9712318]; T. A. Lahde, C. J. Nyfalt and D. O. Riska, Nucl. Phys. A **674** (2000) 141 [arXiv:hep-ph/9908485];

T. Matsuki and T. Morii, Phys. Rev. D **56** (1997) 5646 [Austral. J. Phys. **50** (1997) 163] [arXiv:hep-ph/9702366]; M. Di Pierro and E. Eichten, Phys. Rev. D **64** (2001) 114004 [arXiv:hep-ph/0104208]. J. Vijande, F. Fernandez and A. Valcarce, J. Phys. G **31** (2005) 481 [arXiv:hep-ph/0411299]; T.-J. Kühn, *Elektroschwache Zerfälle schwerer Mesonen im relativistischen Quarkmodell*, Diploma thesis, Universität Bonn, 2006.

[6] Y. B. Dai, C. S. Huang, C. Liu and S. L. Zhu, Phys. Rev. D **68** (2003) 114011 [arXiv:hep-ph/0306274].

[7] T. Barnes, F. E. Close and H. J. Lipkin, Phys. Rev. D **68** (2003) 054006 [arXiv:hep-ph/0305025].

[8] Y. J. Zhang, H. C. Chiang, P. N. Shen and B. S. Zou, Phys. Rev. D **74** (2006) 014013 [arXiv:hep-ph/0604271]. A. Faessler, T. Gutsche, V. E. Lyubovitskij and Y. L. Ma, Phys. Rev. D **76** (2007) 014005 [arXiv:0705.0254 [hep-ph]].

[9] E. E. Kolomeitsev and M. F. M. Lutz, Phys. Lett. B **582** (2004) 39 [arXiv:hep-ph/0307133]; J. Hofmann and M. F. M. Lutz, Nucl. Phys. A **733** (2004) 142 [arXiv:hep-ph/0308263]; M. F. M. Lutz and M. Soyeur, arXiv:0710.1545 [hep-ph]; D. Gamermann, E. Oset, D. Strottman and M. J. Vicente Vacas, Phys. Rev. D **76** (2007) 074016 [arXiv:hep-ph/0612179].

[10] F. K. Guo, P. N. Shen, H. C. Chiang and R. G. Ping, Phys. Lett. B **641** (2006) 278 [arXiv:hep-ph/0603072].

[11] H. Y. Cheng and W. S. Hou, Phys. Lett. B **566** (2003) 193 [arXiv:hep-ph/0305038]; K. Terasaki, Phys. Rev. D **68** (2003) 011501 [arXiv:hep-ph/0305213]; Y. Q. Chen and X. Q. Li, Phys. Rev. Lett. **93** (2004) 232001 [arXiv:hep-ph/0407062]; L. Maiani, F. Piccinini, A. D. Polosa and V. Riquer, Phys. Rev. D **71** (2005) 014028 [arXiv:hep-ph/0412098]; U. Dmitrasinovic, Phys. Rev. Lett. **94** (2005) 162002; H. Kim and Y. Oh, Phys. Rev. D **72** (2005) 074012 [arXiv:hep-ph/0508251]; M. Nielsen, R. D. Matheus, F. S. Navarra, M. E. Bracco and A. Lozea, Nucl. Phys. Proc. Suppl. **161** (2006) 193 [arXiv:hep-ph/0509131]; Z. G. Wang and S. L. Wan, Nucl. Phys. A **778** (2006) 22 [arXiv:hep-ph/0602080].

[12] A. P. Szczepaniak, Phys. Lett. B **567** (2003) 23 [arXiv:hep-ph/0305060].

[13] O. Lakhina and E. S. Swanson, Phys. Lett. B **650** (2007) 159 [arXiv:hep-ph/0608011].

[14] T. E. Browder, S. Pakvasa and A. A. Petrov, Phys. Lett. B **578** (2004) 365 [arXiv:hep-ph/0307054]; J. Vijande, F. Fernandez and A. Valcarce, Phys. Rev. D **73** (2006)

034002 [Erratum-ibid. D **74** (2006) 059903] [arXiv:hep-ph/0601143].

[15] E. van Beveren and G. Rupp, Phys. Rev. Lett. **91** (2003) 012003 [arXiv:hep-ph/0305035]; Phys. Rev. Lett. **97** (2006) 202001 [arXiv:hep-ph/0606110]; D. S. Hwang and D. W. Kim, Phys. Lett. B **601** (2004) 137 [arXiv:hep-ph/0408154]; Yu. A. Simonov and J. A. Tjon, Phys. Rev. D **70** (2004) 114013 [arXiv:hep-ph/0409361].

[16] Yuan-Ben Dai, Xin-Qiang Li, Shi-Lin Zhu, Ya-Bing Zuo, arXiv:hep-ph/0610327.

[17] S. Weinberg, Phys. Rev. **137** (1965) B672.

[18] V. Baru, J. Haidenbauer, C. Hanhart, Yu. Kalashnikova and A. E. Kudryavtsev, Phys. Lett. B **586** (2004) 53 [arXiv:hep-ph/0308129].

[19] K. Abe *et al.* [Belle Collaboration], Phys. Rev. D **69** (2004) 112002 [arXiv:hep-ex/0307021].

[20] J. M. Link *et al.* [FOCUS Collaboration], Phys. Lett. B **586** (2004) 11 [arXiv:hep-ex/0312060].

[21] M. B. Wise, Phys. Rev. D **45** (1992) 2188; T. M. Yan, H. Y. Cheng, C. Y. Cheung, G. L. Lin, Y. C. Lin and H. L. Yu, Phys. Rev. D **46** (1992) 1148 [Erratum-ibid. D **55** (1997) 5851]; G. Burdman and J. F. Donoghue, Phys. Lett. B **280** (1992) 287.

[22] U. Kilian, J. G. Körner and D. Pirjol, Phys. Lett. B **288** (1992) 360.

[23] R. Casalbuoni, A. Deandrea, N. Di Bartolomeo, R. Gatto, F. Feruglio and G. Nardulli, Phys. Rept. **281** (1997) 145 [arXiv:hep-ph/9605342].

[24] S. Capstick *et al.*, arXiv:0711.1982 [hep-ph].

[25] S. Ono and N. A. Törnqvist, Z. Phys. C **23** (1984) 59; P. Geiger and N. Isgur, Phys. Rev. D **41** (1990) 1595; T. Barnes and E. S. Swanson, arXiv:0711.2080 [hep-ph].

[26] M. R. Pennington and D. J. Wilson, Phys. Rev. D **76** (2007) 077502 [arXiv:0704.3384 [hep-ph]].

[27] J. A. Oller and E. Oset, Phys. Rev. D **60** (1999) 074023 [arXiv:hep-ph/9809337].

[28] J. A. Oller and Ulf-G. Meißner, Phys. Lett. B **500** (2001) 263 [arXiv:hep-ph/0011146].

[29] P. Colangelo, F. De Fazio, G. Nardulli, N. Di Bartolomeo and R. Gatto, Phys. Rev. D **52** (1995) 6422 [arXiv:hep-ph/9506207].

[30] Z. G. Wang and S. L. Wan, Phys. Rev. D **74** (2006) 014017 [arXiv:hep-ph/0606002].

[31] V. Bernard, N. Kaiser and Ulf-G. Meißner, Int. J. Mod. Phys. E **4** (1995) 193 [arXiv:hep-ph/9501384].

[32] S. Scherer, Adv. Nucl. Phys. **27** (2003) 277 [arXiv:hep-ph/0210398].

[33] J. L. Rosner, Phys. Rev. D **74**, 076006 (2006) [arXiv:hep-ph/0608102].