

Dynamically generated 0^+ heavy mesons in a heavy chiral unitary approach

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Abstract

In terms of the heavy chiral Lagrangian and the unitarized coupled-channel scattering amplitude, interaction between the heavy meson and the light pseudoscalar meson is studied. By looking for the pole of scattering matrix on an appropriate Riemann sheet, a DK bound state D_{s0}^* with the mass of 2.312 ± 0.041 GeV is found. This state can be associated as the narrow $D_{sJ}^*(2317)$ state found recently. In the same way, a $B\bar{K}$ bound state B_{s0}^* is found, and its mass of 5.725 ± 0.039 GeV is predicted. The spectra of D_0^* and B_0^* with $I = 1/2$ are further investigated. One broad and one narrow states are predicted in both charm and bottom sectors. The coupling constants and decay widths of the predicted states are also calculated.

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1. Introduction

The recently discovered narrow-width state $D_{sJ}^*(2317)$ [1] stimulates both experimental [2–6] and theoretical [7–36] interest. Many physicists surmised that this new state is a conventional $c\bar{s}$ state [7–17], and the others believed that it can be an exotic meson state, such as a four-quark state [18–23], a $D_s\pi$ quasi-bound state [24], a DK bound state [25–28], a mixed state of $c\bar{s}$ with DK [29,30] or with four-quark state [31,32], and etc. On the other hand, one proposed that $D_{sJ}^*(2317)$ with $J^P = 0^+$ could be the chiral partner of the ground state of D_s [8,12]. However, the author in Ref. [33] mentioned that the chiral doubler produced by using random phase approximation equations should be $(D_s(1968), D_s(2392))$ rather than $(D_s(1968), D_s(2317))$, although the scalar state $D_s^*(2392)$, as the scalar chiral partner of $D_s(1968)$ state, has not been found yet [33]. Up to now, the structure of $D_{sJ}^*(2317)$ is still indistinct and should carefully be studied. Moreover, the Belle Collaboration recently reported a broad 0^+ charmed meson with mass and width being $m_{D_0^{*0}} = 2308 \pm 60$ MeV and $\Gamma_{D_0^{*0}} = 276 \pm 99$ MeV, respectively [37], and the FOCUS Collaboration reported a broad 0^+ charmed meson with mass and width being $m_{D_0^{*0}} = 2407 \pm 56$ MeV and $\Gamma_{D_0^{*0}} = 240 \pm 114$ MeV, respectively [38]. Though they are consistent with each other within experimental errors, whether they are the same particle is still in dispute [22,39].

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On the other hand, it has been shown that the light scalar mesons σ , $f_0(980)$, $a_0(980)$ and κ can dynamically be generated through the S wave interaction between Goldstone bosons in the chiral unitary approach (ChUA) [40–46]. In such an approach, the amplitudes from the chiral perturbation theory (ChPT) are usually adopted as the kernels of the factorized coupled-channel Bethe–Salpeter (BS) equations. In this procedure, a Lagrangian in a specific expanded order, where the symmetries of ChPT should be preserved, is chosen at the beginning, and then the higher order corrections to the amplitudes are re-summed with the symmetries kept up to the order of the expansion considered. Namely, what the unitary CHPT does in the successive step is re-summing a string of infinite loop diagrams while the symmetries of ChPT are held [47–49]. Moreover, ChUA has been applied to study the S wave interaction between the lower lying vector meson and the Goldstone boson, and most of the known axial-vector mesons can also be generated dynamically [50]. Based on the valuable achievements mentioned above, extending ChUA to the heavy-light meson sector to study the S wave interaction between the heavy pseudoscalar meson and the Goldstone boson, and consequently the structures of possible heavy scalar mesons, would be extremely meaningful. In fact, similar work, called χ -BS(3) approach, has been done [26,27]. In such an approach, heavy–light meson resonances and open-charm meson resonances were predicted through checking speed plots together with the real and imaginary parts of the reduced scattering amplitudes. In our opinion, studying the poles on the appropriate Riemann sheet of the scattering amplitude would be a powerful procedure to reveal the properties of the generated states in a more accurate way. In this Letter, the S wave interaction between the heavy meson and the light pseudoscalar meson is studied by using the extended chiral unitary approach, called heavy chiral unitary approach. The poles that associate with the experimentally observed narrow $D_{sJ}^*(2317)$ and broad D_0^* in the $I = 0, S = 1$ and $I = \frac{1}{2}, S = 0$ channels, where I and S denote the isospin and the strangeness, respectively, are searched. The corresponding coupling constants and decay widths are also discussed.

2. Coupled-channel heavy chiral unitary approach

In order to describe the interaction between the Goldstone boson and the heavy pseudoscalar boson, we employ a leading order heavy chiral Lagrangian [51–53]

$$\mathcal{L} = \frac{1}{4f_\pi^2} (\partial^\mu P[\Phi, \partial_\mu \Phi]P^\dagger - P[\Phi, \partial_\mu \Phi]\partial^\mu P^\dagger), \tag{1}$$

where $f_\pi = 92.4$ MeV is the pion decay constant, P represents the charmed mesons ($c\bar{u}, c\bar{d}, c\bar{s}$), namely (D^0, D^+, D_s^+) , and Φ denotes the octet Goldstone bosons and can be written in the form of 3×3 matrix

$$\Phi = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix}. \tag{2}$$

This Lagrangian is equivalent to the $SU(4)$ extrapolation of the ordinary meson–meson chiral Lagrangian, eliminating the exchanges of heavy vector mesons in the equivalent picture of vector meson exchange [54]. Obviously, the similar investigation in the bottom sector can be carried out by replacing P in Eq. (1) with the anti-bottom mesons ($b\bar{u}, b\bar{d}, b\bar{s}$), namely $(B^-, \bar{B}^0, \bar{B}_s)$.

We are interested in the heavy mesons in the $I = 0, S = 1$ and $I = \frac{1}{2}, S = 0$ channels that can be specified by their own isospins, respectively. In terms of Eq. (1), the amplitudes can easily be obtained by

$$V_{ij}^I(s, t, u) = \frac{C_{ij}^I}{4f_\pi^2}(s - u), \tag{3}$$

where i and j represent the initial state and the final state, respectively. In the $I = 0$ case, i (j) can be 1 and 2 which represent the coupled DK and $D_s\eta$ channels in the charmed sector, respectively, and $B\bar{K}$ and $B_s\eta$ channels in the bottom sector, respectively. In the $I = \frac{1}{2}$ case, i (j) can take 1, 2 and 3 which denote the coupled $D\pi$, $D\eta$ and $D_s\bar{K}$ channels in the charmed sector, respectively, and $B\pi$, $B\eta$ and B_sK channels in the bottom sector, respectively. The coefficients C_{ij}^I are listed in Table 1.

The tree level amplitudes can be projected to the S wave by using

$$V_{ij}^{I,l=0}(s) = \frac{1}{2} \int_{-1}^1 d\cos\theta V_{ij}^I(s, t(s, \cos\theta), u(s, \cos\theta)), \tag{4}$$

Table 1
Coefficients C_{ij}^I in Eq. (3)

C_{11}^0	C_{12}^0	C_{22}^0	$C_{11}^{1/2}$	$C_{12}^{1/2}$	$C_{22}^{1/2}$	$C_{13}^{1/2}$	$C_{23}^{1/2}$	$C_{33}^{1/2}$
-2	$\sqrt{3}$	0	-2	0	0	$-\frac{\sqrt{6}}{2}$	$-\frac{\sqrt{6}}{2}$	-1

and

$$-u(s, \cos \theta) = s - m_2^2 - m_4^2 - 2\sqrt{\left[m_1^2 + \frac{\lambda(s, m_1^2, m_2^2)}{4s}\right] \left[m_3^2 + \frac{\lambda(s, m_3^2, m_4^2)}{4s}\right]} + \frac{1}{2s}\sqrt{\lambda(s, m_1^2, m_2^2)\lambda(s, m_3^2, m_4^2)} \cos \theta, \quad (5)$$

where $\lambda(s, m_i^2, m_j^2) = [s - (m_i + m_j)^2][s - (m_i - m_j)^2]$ and the on-shell condition for the Mandelstam variables, $s + t + u = \sum_{i=1}^4 m_i^2$, is applied.

In ChUA, under the on-shell approximation, the full scattering amplitude can be converted into an algebraic BS equation [40]

$$T = (1 - VG)^{-1}V, \quad (6)$$

where V is a matrix whose elements are the S wave projections of the tree diagram amplitudes and G is a diagonal matrix with the element being a two-meson loop integral

$$G_{ii}(s) = i \int \frac{d^4q}{(2\pi)^4} \frac{1}{q^2 - m_1^2 + i\varepsilon} \frac{1}{(p_1 + p_2 - q)^2 - m_2^2 + i\varepsilon}, \quad (7)$$

where p_1 and p_2 are the four-momenta of the two initial particles, respectively, and m_1 and m_2 are the masses of the particles appearing in the loop. It was shown that the scattering matrix derived in such a way satisfies the unitary relation [41,42,46].

The loop integral can usually be calculated in the center-of-mass frame by using a three-momentum cut-off parameter q_{\max} [40]. However, in this method, an artificial singularity of the loop function might be produced [46], and the applicability of the method is limited. The better way to remove the singularity of the loop integral is using the dispersion relation where a subtraction constant is employed. Then, the analytic expression of $G_{ii}(s)$ can be expressed by [42]

$$G_{ii}(s) = \frac{1}{16\pi^2} \left\{ a(\mu) + \log \frac{m_1^2}{\mu^2} + \frac{\Delta - s}{2s} \log \frac{m_1^2}{m_2^2} + \frac{\sigma}{2s} [\log(s - \Delta + \sigma) + \log(s + \Delta + \sigma) - \log(-s + \Delta + \sigma) - \log(-s - \Delta + \sigma)] \right\}, \quad (8)$$

where $a(\mu)$ is the subtraction constant, μ denotes the regularization scale, $\sigma = [-(s - (m_1 + m_2)^2)(s - (m_1 - m_2)^2)]^{1/2}$ and $\Delta = m_1^2 - m_2^2$. This result is independent of μ , because the change in G_{ii} , caused by a variation of μ , is cancelled by the corresponding change of the subtraction constant $a(\mu)$.

3. Poles on appropriate Riemann sheets

The physical states are closely associated with the poles of the scattering amplitude on the appropriate Riemann sheet of the energy plane. For instance, considering only one channel, a bound state is associated with a pole below the threshold value in the real axis of the energy plane, and the three-momentum of the scattered meson in the center of mass frame of the two mesons system can be written as $p_{\text{cm}} = i|p_{\text{cm}}|$. A resonance should be related with a pole on the second Riemann sheet, namely, $\text{Im } p_{\text{cm}} < 0$. In the coupled channel case, the situation is somewhat complicated. Detailed relation can be found in Ref. [55].

Before searching for poles of the scattering amplitude, the range of subtraction constant values in the dispersion relation method should firstly be estimated. It can be done by comparing the calculated value of loop integration in the dispersion relation method with the one obtained in the cut-off method, although there might be an artificial singularity problem in the cut-off method [46]. The cut-off momentum can approximately be chosen as

$$q_{\max} \sim \sqrt{\Lambda_\chi^2 - m_\phi^2}, \quad (9)$$

where m_ϕ is the mass of the Goldstone boson and Λ_χ denotes the chiral symmetry breaking scale which is about 1 GeV. The resultant q_{\max} for $\phi = \pi, K$ and η are all in the region of 0.8–0.9 GeV. Thus, it is reasonable to pick up a value of q_{\max} in the region of 0.8 ± 0.2 GeV. Then, we adjust the renormalization scale μ or the subtraction constant $a(\mu)$ to match the calculated value of the loop integral in the dispersion relation method with the one obtained in the cut-off method at $\sqrt{s} = m_D(m_B) + m_K$ in a specific q_{\max} value case, say $q_{\max} = 0.6, 0.8$ and 1.0 GeV, respectively. The resultant loop integration curves versus s in two different methods are very close in the region around and below the matching point \sqrt{s} . The corresponding values of $a(\mu)$ and q_{\max} are tabulated in Table 2. With the estimated $a(\mu)$ value, the full scattering amplitude can be calculated.

The poles of the scattering matrix in the $I = 0, S = 1$ channel in both the charmed sector and bottom sector are searched for first. It is shown that on the first Riemann sheet of the energy plane, there is only one pole located on the real axis below the lowest strong decay threshold, $m_D + m_K = 2.367$ GeV, in the charmed sector and only one pole on the real axis below the lowest strong decay threshold, $m_B + m_K = 5.773$ GeV, in the bottom sector as well. The resultant pole positions with different $a(\mu)$, which correspond to the $q_{\max} = 0.6, 0.8$ and 1.0 GeV cases, are tabulated in Table 3, respectively. These poles are apparently associated with the

Table 2

The values of $a(\mu)$ from matching. We use $\mu = m_D$ for the charm sector, and $\mu = m_B$ for the bottom sector, respectively

q_{\max} (GeV)	0.6	0.8	1.0
$a(m_D)$	-0.373	-0.630	-0.864
$a(m_B)$	0.0232	-0.0856	-0.187

Table 3

Poles in the $(I, S) = (0, 1)$ channel

q_{\max} (GeV)	0.6	0.8	1.0
D_{s0}^* (GeV)	2.353	2.317	2.270
B_{s0}^* (GeV)	5.764	5.729	5.661

Table 4

Mass of B_{s0}^* predicted in different models

	Our result	[8]	[33]	[16]	[26]
$m_{B_{s0}^*}$ (GeV)	5.725 ± 0.039	5.728 ± 0.035	5.71 ± 0.03	5.627	5.643

Table 5

Poles in $(I, S) = (\frac{1}{2}, 0)$ channel

q_{\max} (GeV)	0.6	0.8	1.0
D_0^* (GeV)	2.115 - i 0.147 2.488 - i 0.039	2.099 - i 0.100 2.445 - i 0.049	2.079 - i 0.067 2.429 - i 0.002
B_0^* (GeV)	5.564 - i 0.160 5.864 - i 0.027	5.534 - i 0.110 5.827 - i 0.026	5.507 - i 0.074 5.821 - i 0.019

DK bound state and the $B\bar{K}$ bound state, respectively. Due to the existence of the \bar{s} quark, these bound states should be scalar heavy mesons, namely D_{s0}^* and \bar{B}_{s0}^* , respectively. More specifically, when $a(m_D) = -0.630$, corresponding to $q_{\max} = 0.8$ GeV, the mass of the DK state, namely D_{s0}^* , is about 2317 MeV, which is almost the same as the measured value of D_{sJ}^* (2317). Taking into account the uncertainty of subtraction constant, the mass of the D_{s0}^* (0, 1) state in our model is 2.312 ± 0.041 GeV. Also due to the uncertainty of $a(m_B)$, the predicted mass of the $B\bar{K}$ bound state, namely B_{s0}^* (0, 1) state, is 5.725 ± 0.039 GeV. This mass is consistent with the mass predicted in Refs. [8,33], but larger than that in Refs. [16,26]. For comparison, we list the mass of B_{s0}^* predicted in different models in Table 4.

In the $I = \frac{1}{2}, S = 0$ case, the poles are located on nonphysical Riemann sheets. Usually, if $\text{Im } p_{\text{cm}}$ is negative for all the channels open for a certain energy, the width obtained would correspond more closely with the physical one. We search for poles in this particular sheet.

There are two poles in either charmed sector or bottom sector. The width of the lower pole is broad and the width of the higher one is narrow. The obtained poles are listed in Table 5. In either the charmed or bottom sector, the lower pole is located on the second Riemann sheet ($\text{Im } p_{\text{cm}1} < 0, \text{Im } p_{\text{cm}2} > 0, \text{Im } p_{\text{cm}3} > 0$, where $p_{\text{cm}i}$ denotes the momentum of one of the interacting mesons in the i -th channel in the center of mass system). This pole should be associated with a $D\pi$ ($B\pi$) resonance in the charmed (bottom) sector. Consequently, this state should easily decay into $D\pi$ ($B\pi$) in the charmed (bottom) sector.

The higher pole in either charmed or bottom sector is found on the third Riemann sheet ($\text{Im } p_{\text{cm}1} < 0, \text{Im } p_{\text{cm}2} < 0, \text{Im } p_{\text{cm}3} > 0$) when $a(\mu)$ corresponds to $q_{\max} = 0.6$ GeV or 0.8 GeV, or on the second Riemann sheet when $a(\mu)$ corresponds to $q_{\max} = 1.0$ GeV. The pole should be associated with an unstable $D_s\bar{K}$ ($B_s K$) bound state in the charmed (bottom) sector due to its narrow width. It should be mentioned that the situation for the higher pole in the later case, namely $a(\mu)$ corresponding to $q_{\max} = 1.0$ GeV, is somewhat complicated. Besides a pole on the second Riemann sheet, $\text{pole}_{\text{II}} = 2.429 - i0.002$ GeV shown in Table 5, there is a shadow pole, $\text{pole}_{\text{III}} = 2.397 - i0.043$ GeV, on the third Riemann sheet. Note that $\text{Re}(\text{pole}_{\text{II}}) > m_D + m_\eta$ and $\text{Re}(\text{pole}_{\text{III}}) < m_D + m_\eta$. A sketch plot for the paths of these two poles to the physical region in the energy plane is shown in Fig. 1. From this cartoon, one sees that pole_{II} corresponds more closely with the physical one. Therefore, we choose $\text{pole}_{\text{II}} = 2.429 - i0.002$ GeV as the result. Similar complexity appears at $\text{pole}_{\text{III}} = 2.488 - i0.039$ GeV in Table 5, due to the existence of $\text{pole}_{\text{V}} = 2.048 - i0.020$ GeV. With the same reason, we disregard pole_{V} .

Considering the deviations of the data caused by the uncertainty of $a(\mu)$ in Table 5, we predict the mass and the width of the broad D_0^* ($\frac{1}{2}, 0$) state as 2.097 ± 0.018 GeV and 0.213 ± 0.080 GeV, respectively, and the mass and the width of the narrow D_0^* ($\frac{1}{2}, 0$) state as 2.448 ± 0.030 GeV and 0.051 ± 0.047 GeV, respectively. In the same way, we forecast the mass and the width of

than that to the $D_s \bar{K}$ ($B_s K$) channel and the $D\eta$ ($B\eta$) channel, and the coupling of the higher state to the $D_s \bar{K}$ ($B_s K$) channel is stronger than that to the $D\eta$ ($B\eta$) channel and the $D\pi$ ($B\pi$) channel. These are consistent with the pole analysis interpreting the lower one being the $D\pi$ ($B\pi$) resonance and the higher one being the unstable bound state of $D_s \bar{K}$ ($B_s K$).

The decay widths of generated states are further evaluated. We first study the states in the $(0, 1)$ channel. The D_{s0}^* state cannot decay into either DK or $D_s\eta$, because the mass of the state is lower than the threshold of the DK channel. Moreover, the $D_{s0}^{*+}(2317) \rightarrow D_s^+\pi^0$ decay violates the isospin symmetry. Thus, the decay width of $D_{s0}^{*+}(2317)$ should be very small. This decay can only occur through π^0 - η mixing [57]. According to Dashen's theorem [58], the π^0 - η transition matrix should be

$$t_{\pi\eta} = \langle \pi^0 | \mathcal{H} | \eta \rangle = -0.003 \text{ GeV}, \quad (12)$$

and the decay width reads

$$\Gamma = \frac{p_{\text{cm}}}{8\pi M^2} \left| \frac{g_2 t_{\pi\eta}}{m_{\pi^0}^2 - m_\eta^2} \right|^2, \quad (13)$$

where M is the mass of the initial state, g_2 represents the coupling of $D_{s0}^*(2317)$ to $D\eta$, and p_{cm} denotes the three-momentum in the center of mass frame and can be written as

$$p_{\text{cm}} = \frac{1}{2M} \sqrt{(M^2 - (m_{D^+} + m_{\pi^0})^2)(M^2 - (m_{D^+} - m_{\pi^0})^2)}. \quad (14)$$

Then, the partial decay width of the $D_{s0}^{*+}(2317) \rightarrow D_s^+\pi^0$ process can be obtained as

$$\Gamma(D_{s0}^{*+}(2317) \rightarrow D_s^+\pi^0) = 8.69 \text{ keV}. \quad (15)$$

This value is compatible with that in Refs. [35,36]. Similarly, the partial decay width of the isospin violated decay $B_{s0}^*(5729) \rightarrow B_s^0\pi^0$ can be evaluated as

$$\Gamma(B_{s0}^*(5729) \rightarrow B_s^0\pi^0) = 1.54 \text{ keV}. \quad (16)$$

We then study the states in the $(\frac{1}{2}, 0)$ channel. For the higher state, two strong decay channels are opened. The fraction ratio of the decay widths for these two decay channels can be calculated by utilizing the coupling constants given in Table 7. Let Γ_1 and Γ_2 denote the partial decay widths with the final states being $D(B)\pi$ and $D(B)\eta$, respectively. The ratio $\Gamma_1/(\Gamma_1 + \Gamma_2)$ can be written by

$$R \equiv \frac{\Gamma_1}{\Gamma_1 + \Gamma_2} = \frac{|g_1|^2 p_{\text{cm}1}}{|g_1|^2 p_{\text{cm}1} + |g_2|^2 p_{\text{cm}2}}. \quad (17)$$

For higher D_0^* and B_0^* states, we have

$$R(D_0^*) = 0.446, \quad R(B_0^*) = 0.829. \quad (18)$$

It is shown that in the bottom sector, the higher narrow state is easier to decay into $B\pi$ than into $B\eta$, but in the charmed sector, the higher narrow state can decay into $D\pi$ and $D\eta$ in almost the same weight.

5. Conclusion

Based on the heavy chiral unitary approach, the S wave interaction between the pseudoscalar heavy meson and the Goldstone boson is studied. By calculating full scattering amplitudes via an algebraic BS equation, the poles on some appropriate Riemann sheets are found. These poles can be associated with bound states or resonances. With a reasonably estimated single parameter $a(\mu)$ in the loop integration, a pole on the real axis on the first Riemann sheet, which is associated with the bound state, in the two-coupled-channel calculation in the $(0, 1)$ channel is found. Because the mass of the pole in the charmed sector is about 2.312 ± 0.041 GeV, this state should be a 0^+ DK bound state and can be regarded as the recently observed $D_{sJ}^*(2317)$. Meanwhile, a 0^+ state B_{s0}^* , which should be a $B\bar{K}$ bound state, is predicted. Its mass is about 5.725 ± 0.039 GeV. In the $I = \frac{1}{2}$, $S = 0$ case, three-coupled-channel calculations are performed in both charmed and bottom sectors. In the charm sector, a broad pole structure, which is associated with a resonance, is found at about $(2.097 \pm 0.018 - i0.107 \pm 0.040)$ GeV. Besides, a narrow pole structure, which can be interpreted as a quasi-bound state of $D_s \bar{K}$, at about $(2.448 \pm 0.030 - i0.026 \pm 0.024)$ GeV is also found. In the bottom sector, one broad and one narrow poles are found at about $(5.536 \pm 0.029 - i0.117 \pm 0.043)$ GeV and $(5.842 \pm 0.022 - i0.018 \pm 0.010)$ GeV, respectively. The coupling constants of the generated states to the relevant coupled channels are calculated. They are consistent with the results in the pole structure analysis. In the $(0, 1)$ channel, the width of the isospin violated decays $D_{s0}^{*+}(2317) \rightarrow D_s^+\pi^0$ and $B_{s0}^*(5729) \rightarrow B_s^0\pi^0$ are calculated. They are about 8.69 and 1.54 keV, respectively. Finally in the $(\frac{1}{2}, 0)$ channel, the decay ratio $\Gamma_1/(\Gamma_1 + \Gamma_2)$ for the higher narrow state is also estimated.

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