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On charm baryon resonances and chiral symmetry

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Abstract

We study heavy-light baryon resonances with quantum numbers $J^P = \frac{1}{2}^-$ in terms of the non-linear chiral SU(3) Lagrangian. Within the χ–BS(3) approach a parameter-free leading order prediction is obtained for the scattering of Goldstone bosons off heavy-light baryon resonances with $J^P = \frac{1}{2}^+$. The three states $\Lambda_{c1}(2593)$, $\Lambda_{c0}(2880)$ and $\Xi_{c1}(2790)$ discovered by the CLEO Collaboration are recovered. We suggest the existence of resonance states that form an anti-quindecimplet, two sextet and two anti-triplet representations of the SU(3) group. In particular, narrow states with anomalous isospin (I) and strangeness (S) quantum numbers $(I, S) = (\frac{1}{2}, +1)$ are anticipated.

1. Introduction

In a recent work [1] it was demonstrated that chiral coupled-channel dynamics generates heavy-light meson resonances with quantum numbers $J^P = 0^+$ and $J^P = 1^+$. Such states were first predicted in [2,3] based on the chiral quark model and recently observed for the first time by the BaBar and CLEO Collaborations [4,5]. Due to the different dynamical assumptions of the chiral quark model versus the chiral coupled-channel approach the theoretical prediction are quite different (see [6,7] and references therein). Most spectacular is the prediction of $J^P = 0^+$ and $J^P = 1^+$ heavy-light meson states with negative strangeness. Whereas the chiral quark model implies an anti-triplet of open-charm or open-beauty states only, an additional sextet of states was predicted in [1]. These findings suggest that the chiral SU(3) symmetry plays a decisive role also in the physics of heavy-light

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baryons. Double-charm or double-beauty baryon states are completely analogous to the heavy-light mesons and all results from [1] can be straightforwardly taken over to this sector. Therefore in this paper we focus on baryons with one charm quark. The chiral quark model predicts $\frac{1}{2}^+$ and $\frac{3}{2}^+$ states that form anti-triplet and sextet representations of the SU(3) group, respectively [2,3,6,7]. The main goal of this work is to unravel the consequences of chiral coupled-channel dynamics for such states. Recently three $J^P = \frac{1}{2}^-$ states were observed by the CLEO Collaboration [8–10]. Open-charm baryons have been studied extensively in the literature [11–21]. For review articles on the heavy-quark effective theory approach we refer to [12,13].

We apply the χ–BS(3) approach developed originally for meson–baryon scattering [22–27] that is based on the *chiral* SU(3) Lagrangian and formulated in terms of the Bethe–Salpeter equation. The latter was applied recently successfully also to meson–meson scattering [1,28]. Using the chiral SU(3) Lagrangian with heavy-light baryon fields that transform non-linearly under the chiral SU(3) group a coupled-channel description of the meson–baryon scattering in the open charm sector is developed. The crucial importance of coupled-channel dynamics for the baryon-resonance formation in the u-, d-, s-sector of QCD was first anticipated in a series of works in the sixties [29–33]. Related works based on the chiral SU(3) Lagrangian are [34–41]. Our major result is the prediction that there exist open-charm states with $J^P = \frac{1}{2}^-$ quantum numbers forming one anti-quindecimplet, two anti-triplet and two sextet representations of the SU(3) group. This differs from the results implied by the linear realization of the chiral SU(3) symmetry leading to one anti-triplet and one sextet only. We recover the so far known $J^P = \frac{1}{2}^-$ resonance states established by the CLEO collaboration [8–10]. Most spectacular is the promise of new $J^P = \frac{1}{2}^-$ states with anomalous quantum numbers $(I, S) = (\frac{1}{2}, +1), (\frac{3}{2}, -1), (1, -2)$.

2. Chiral coupled-channel dynamics: the χ-BS(3) approach

The starting point to study the scattering of Goldstone bosons off heavy-light baryons is the chiral SU(3) Lagrangian. We identify the leading order interaction Lagrangian density [42–47] describing the interaction of Goldstone bosons with heavy-light baryons,

$$\begin{aligned}\mathcal{L}(x) = {}& \frac{i}{16 f^2}\,\mathrm{tr}\big(\overline{H}_{[\bar 3]}(x)\,\gamma^\mu\big[H_{[\bar 3]}(x),\big[\phi(x),\big(\partial_\mu\phi(x)\big)\big]_-\big]_+\big) \\ & + \frac{i}{16 f^2}\,\mathrm{tr}\big(\overline{H}_{[6]}(x)\gamma^\mu\big[H_{[6]}(x),\big[\phi(x),\big(\partial_\mu\phi(x)\big)\big]_-\big]_+\big), \qquad (1)\end{aligned}$$

with the Goldstone bosons field ϕ and massive baryon fields $H_{[\bar 3]}$ and $H_{[6]}$. The Weinberg–Tomozawa term (1) follows by gauging the kinetic term of the heavy-baryon fields with the chiral SU(3) group and expanding the resulting expression in powers of the Goldstone bosons fields. The parameter f in (1) characterizes the associated covariant derivative and is known from the weak decay process of the pions. We use $f = 90$ MeV through out this work. Since we will assume perfect isospin symmetry it is convenient to decompose the fields into their isospin multiplets. The fields can be written in terms of isospin multiplet fields like $K = (K^{(+)}, K^{(0)})^t$ and $\Xi_c = (\Xi_c^{(+)}, \Xi_c^{(0)})^t$,

$$\phi = \tau \cdot \pi(140) + \alpha^\dagger \cdot K(494) + K^\dagger(494) \cdot \alpha + \eta(547)\lambda_8,$$
$$H_{[\bar{3}]} = \frac{1}{\sqrt{2}}\alpha^\dagger \cdot \Xi_c(2470) - \frac{1}{\sqrt{2}}\Xi_c^t(2470) \cdot \alpha + i\tau_2 \Lambda_c(2284),$$
$$H_{[6]} = \frac{1}{\sqrt{2}}\alpha^\dagger \cdot \Xi_c'(2580) + \frac{1}{\sqrt{2}}\Xi_c^{\prime t}(2580) \cdot \alpha + \Sigma_c(2453) \cdot \tau i\tau_2 + \frac{\sqrt{2}}{3}\big(1 - \sqrt{3}\lambda_8\big)\Omega_c(2704),$$
$$\alpha^\dagger = \frac{1}{\sqrt{2}}(\lambda_4 + i\lambda_5, \lambda_6 + i\lambda_7), \quad \tau = (\lambda_1, \lambda_2, \lambda_3), \tag{2}$$

where the matrices λ_i are the standard Gell-Mann generators of the SU(3) algebra. The numbers in the brackets recall the approximate masses of the particles in units of MeV [48].

The scattering problem decouples into ten orthogonal channels specified by isospin (I) and strangeness (S) quantum numbers,

$$(I,S)_{[\bar{3}]} = \big(\big(\tfrac{1}{2}, +1\big), (0,0), (1,0), \big(\tfrac{1}{2}, -1\big), \big(\tfrac{3}{2}, -1\big), (0,-2), (1,-2)\big),$$
$$(I,S)_{[6]} - (I,S)_{[\bar{3}]} = \big(\big(\tfrac{3}{2}, +1\big), (2,0), \big(\tfrac{1}{2}, -3\big)\big), \tag{3}$$

where the scattering of Goldstone bosons off the anti-triplet leads to seven channels but the scattering off the sextet to additional three channels. At leading order the two sectors $(I,S)_{[\bar{3}]}$ and $(I,S)_{[6]}$ do not couple to each other. Only subleading terms in the chiral Lagrangian lead to processes like $\pi\Xi_c \to K\Omega_c$. In Table 1 the channels that contribute in a given sector (I,S) are listed. Heavy-light baryon resonances with quantum numbers $J^P = \frac{1}{2}^-$ manifest themselves as poles in the s-wave scattering amplitudes, $M^{(I,S)}_{[\bar{3}]}(\sqrt{s}\,)$, and, $M^{(I,S)}_{[6]}(\sqrt{s}\,)$, which in the χ–BS(3) approach [25,28] take the simple form

$$M^{(I,S)}\big(\sqrt{s}\,\big) = \big[1 - V^{(I,S)}\big(\sqrt{s}\,\big) J^{(I,S)}\big(\sqrt{s}\,\big)\big]^{-1} V^{(I,S)}\big(\sqrt{s}\,\big). \tag{4}$$

The effective interaction kernel $V^{(I,S)}(\sqrt{s}\,)$ in (4) is determined by the leading order chiral SU(3) Lagrangian (1),

$$V^{(I,S)}\big(\sqrt{s}\,\big) = \frac{C^{(I,S)}}{4f^2}\big(2\sqrt{s} - M - \overline{M}\,\big), \tag{5}$$

where M and $\overline{M}$ are the masses of initial and final baryon states. We use capital M for the masses of heavy-light baryons and small m for the masses of the Goldstone bosons. The matrix of coefficients $C^{(I,S)}$ that characterize the interaction strength in a given channel is given in Tables 2 and 3. The s-wave interaction kernels are identical for the two scattering problems considered here. The loop functions, diagonal in the coupled-channel space, are

$$J\big(\sqrt{s}\,\big) = \big(M + \big(M^2 + p_{\rm cm}^2\big)^{1/2}\big)\big(I\big(\sqrt{s}\,\big) - I(\mu)\big),$$
$$I\big(\sqrt{s}\,\big) = \frac{1}{16\pi^2}\Bigg(\frac{p_{\rm cm}}{\sqrt{s}}\Bigg(\ln\Bigg(1 - \frac{s - 2p_{\rm cm}\sqrt{s}}{m^2+M^2}\Bigg) - \ln\Bigg(1 - \frac{s + 2p_{\rm cm}\sqrt{s}}{m^2+M^2}\Bigg)\Bigg) + \Bigg(\frac{1}{2}\frac{m^2+M^2}{m^2-M^2} - \frac{m^2-M^2}{2s}\Bigg)\ln\Bigg(\frac{m^2}{M^2}\Bigg) + 1\Bigg) + I(0), \tag{6}$$

Table 1
The definition of coupled-channel states with $(I,S)_{[\bar{3}]}$ and $(I,S)_{[6]}$. Here $\sigma = (\sigma_1, \sigma_2, \sigma_3)$ are the isospin Pauli matrices. The isospin transition operator T connects isospin-$\frac{1}{2}$ and isospin-$\frac{3}{2}$ states. It is normalized by $T_i^\dagger T_j = \delta_{ij} - \sigma_i \sigma_j/3$. The matrix valued vector $S_{[n]}$ couples two isospin-1 states into a spin-2 state. It satisfies $\sum_{n=1}^{5} S^\dagger_{[n],ac} S_{[n],bd} = \frac{1}{2}\delta_{ab}\delta_{cd} + \frac{1}{2}\delta_{ad}\delta_{cb} - \frac{1}{3}\delta_{ac}\delta_{bd}$

$(I,S)_{[\bar{3}]}$			
$(\frac{1}{2},+1)$	$(0,0)$	$(1,0)$	$(\frac{1}{2},-1)$
$(\Lambda_c K)$	$\begin{pmatrix} (\Lambda_c \eta) \\ (\frac{1}{\sqrt{2}} K^t i\sigma_2 \Xi_c) \end{pmatrix}$	$\begin{pmatrix} (\Lambda_c \pi) \\ (\frac{1}{\sqrt{2}} K^t i\sigma_2 \sigma \Xi_c) \end{pmatrix}$	$\begin{pmatrix} (\frac{1}{\sqrt{3}} \pi \cdot \sigma \Xi_c) \\ (\Lambda_c i\sigma_2 \overline{K}^t) \\ (\eta \Xi_c) \end{pmatrix}$
$(\frac{3}{2},-1)$	$(0,-2)$	$(1,-2)$	
$(\pi \cdot T \Xi_c)$	$(\frac{1}{\sqrt{2}} \overline{K} \Xi_c)$	$(\frac{1}{\sqrt{2}} \overline{K} \sigma \Xi_c)$	

$(I,S)_{[6]}$			
$(\frac{1}{2},+1)$	$(\frac{3}{2},+1)$	$(0,0)$	$(1,0)$
$\frac{1}{\sqrt{3}}(\Sigma_c \sigma K)$	$(\Sigma_c \cdot T K)$	$\begin{pmatrix} \frac{1}{\sqrt{3}}(\Sigma_c \cdot \pi) \\ \frac{1}{\sqrt{2}}(K^t i\sigma_2 \Xi'_c) \end{pmatrix}$	$\begin{pmatrix} \frac{i}{\sqrt{2}}(\Sigma_c \times \pi) \\ (\eta \Sigma_c) \\ \frac{1}{\sqrt{2}}(K^t i\sigma_2 \sigma \Xi'_c) \end{pmatrix}$
$(2,0)$	$(\frac{1}{2},-1)$	$(\frac{3}{2},-1)$	$(0,-2)$
$(\Sigma_c \cdot S \cdot \pi)$	$\begin{pmatrix} \frac{1}{\sqrt{3}}(\pi \cdot \sigma \Sigma_c) \\ (\Sigma_c \eta) \\ (\Omega_c K) \end{pmatrix}$	$\begin{pmatrix} (\pi \cdot T \Xi'_c) \\ (\Sigma_c \cdot T i\sigma_2 \overline{K}^t) \end{pmatrix}$	$\begin{pmatrix} \frac{1}{\sqrt{2}}(\overline{K} \Xi'_c) \\ (\Omega_c \eta) \end{pmatrix}$
$(1,-2)$		$(\frac{1}{2},-3)$	
$\begin{pmatrix} \Omega_c \pi \\ \frac{1}{\sqrt{2}}(\overline{K} \sigma \Xi'_c) \end{pmatrix}$		$(\Omega_c i\sigma_2 \overline{K}^t)$	

Table 2
The coefficients $C^{(I,S)}$ that characterize the interaction of Goldstone bosons with the heavy baryon fields $H_{[\bar{3}]}$ as introduced in (3). The ordering of the states is introduced in Table 1

$(I,S)_{[\bar{3}]}$	$(\frac{1}{2},+1)$	$(0,0)$	$(1,0)$	$(\frac{1}{2},-1)$	$(\frac{3}{2},-1)$	$(0,-2)$	$(1,-2)$
11	-1	0	0	2	-1	1	-1
12	–	$-\sqrt{3}$	1	$\sqrt{\frac{3}{2}}$	–	–	–
22	–	2	0	1	–	–	–
13	–	–	–	0	–	–	–
23	–	–	–	$-\sqrt{\frac{3}{2}}$	–	–	–
33	–	–	–	0	–	–	–

where $\sqrt{s} = \sqrt{M^2 + p_{\rm cm}^2} + \sqrt{m^2 + p_{\rm cm}^2}$. A crucial ingredient of the χ–BS(3) scheme is its approximate crossing symmetry guaranteed by a proper choice of the subtraction scales $\mu^{(I,S)}_{[\bar{3}]}$ and $\mu^{(I,S)}_{[6]}$. Referring to the detailed discussions in [25–28] we define

Table 3
The coefficients $C^{(I,S)}$ that characterize the interaction of Goldstone bosons with heavy baryon fields $H_{[6]}$ as introduced in (3). The ordering of the states is introduced in Table 1

$(I,S)_{[6]}$	11	12	22	13	23	33	14	24	34	44
$(\frac{1}{2},+1)$	1	–	–	–	–	–	–	–	–	–
$(\frac{3}{2},+1)$	−1	–	–	–	–	–	–	–	–	–
$(0,0)$	4	$\sqrt{3}$	2	–	–	–	–	–	–	–
$(1,0)$	2	0	0	$\sqrt{2}$	$-\sqrt{3}$	0	–	–	–	–
$(2,0)$	0	–	–	–	–	–	–	–	–	–
$(\frac{1}{2},-1)$	2	$-\frac{1}{\sqrt{2}}$	3	0	$-\frac{3}{\sqrt{2}}$	0	$\sqrt{3}$	0	$\sqrt{3}$	2
$(\frac{3}{2},-1)$	−1	$\sqrt{2}$	0	–	–	–	–	–	–	–
$(0,-2)$	1	$-\sqrt{6}$	0	–	–	–	–	–	–	–
$(1,-2)$	0	$-\sqrt{2}$	−1	–	–	–	–	–	–	–
$(\frac{1}{2},-3)$	−2	–	–	–	–	–	–	–	–	–

$$\mu^{(I,0)}_{[\bar{3}]} = M_{\Lambda_c(2284)}, \qquad \mu^{(I,\pm1)}_{[\bar{3}]} = M_{\Xi_c(2470)}, \qquad \mu^{(I,-2)}_{[\bar{3}]} = M_{\Lambda_c(2284)},$$
$$\mu^{(I,0)}_{[6]} = M_{\Sigma_c(2453)}, \qquad \mu^{(I,\pm1)}_{[6]} = M_{\Xi'_c(2580)}, \qquad \mu^{(I,-2)}_{[6]} = M_{\Omega_c(2704)},$$
$$\mu^{(I,-3)}_{[6]} = M_{\Xi'_c(2580)}. \tag{7}$$

Given the subtraction scales (7) the leading order calculation presented in this work is parameter free. Of course chiral correction terms do lead to further so far unknown parameters which need to be adjusted to data. Within the χ–BS(3) approach such correction terms enter the effective interaction kernel V rather than leading to subtraction scales different from (7). In particular the leading correction effects are determined by the counter terms of chiral order Q^2. The effect of altering the subtraction scales away from their optimal values (7) can be compensated for by incorporating counter terms in the chiral Lagrangian that carry order Q^3. Our scheme has the advantage that once the parameters describing subleading effects are determined in a subset of sectors one has immediate predictions for all sectors (I,S). In order to estimate the size of correction terms it is nevertheless useful to vary the subtraction scales around their optimal values. With (4)–(7) the brief exposition of the χ–BS(3) approach as applied to heavy-light baryon resonances is completed.

3. Results

To study the formation of heavy-light baryon resonances we generate speed plots as suggested by Höhler [49]. The speed $\mathrm{Speed}^{(I,S)}_{ab}(\sqrt{s})$ of a given channel $a\,b$ is introduced by [49,50],

$$t^{(I,S)}_{ab}\big(\sqrt{s}\,\big) = \frac{1}{8\pi\sqrt{s}}\big((M_a+E_a)\,p^{(a)}_{\rm cm}(M_b+E_b)\,p^{(b)}_{\rm cm}\big)^{1/2} M^{(I,S)}_{ab}\big(\sqrt{s}\,\big),$$
$$\mathrm{Speed}^{(I,S)}_{ab}\big(\sqrt{s}\,\big) = \left|\sum_c \left[\frac{d}{d\sqrt{s}}\,t^{(I,S)}_{ac}\big(\sqrt{s}\,\big)\right]\big(\delta_{cb} + 2i\,t^{(I,S)}_{cb}\big(\sqrt{s}\,\big)\big)^{\dagger}\right|, \tag{8}$$

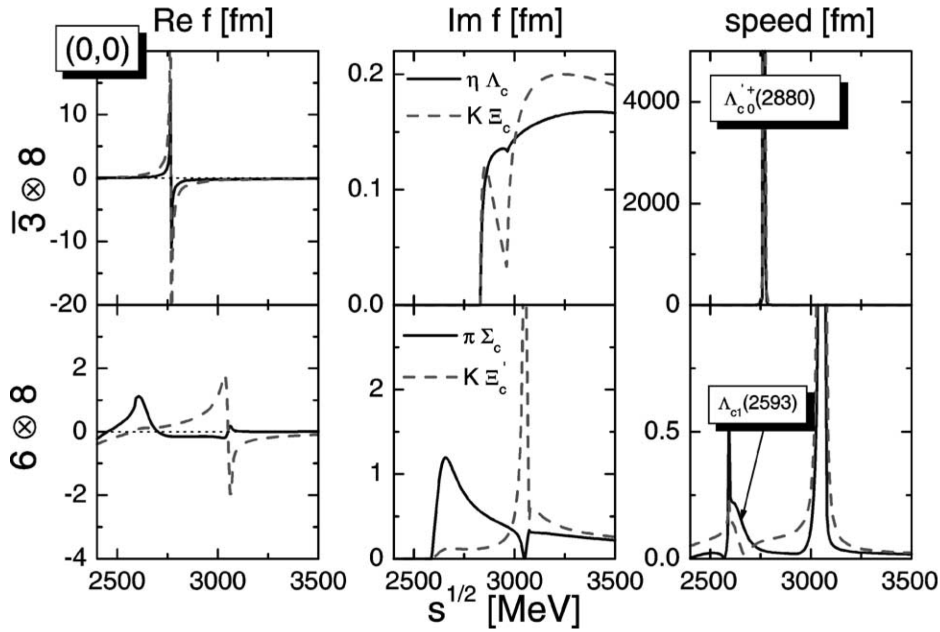

Fig. 1. Open-charm baryon resonances with $J^P = \frac{1}{2}^-$ and $(I, S) = (0, 0)$ as seen in the scattering of Goldstone bosons off anti-triplet $(\Lambda_c(2284), \Xi_c(2470))$ and sextet $(\Sigma_c(2453), \Xi_c'(2580), \Omega_c(2704))$ baryons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, f_{ab}, with $t_{ab} = f_{ab}(p_{\rm cm}^{(a)} p_{\rm cm}^{(b)})^{1/2}$ (see (8)).

where $E_a = (M_a^2 + (p_{\rm cm}^{(a)})^2)^{1/2}$. If a resonance is formed in a scattering process its Speed($\sqrt{s}$) will show a maximum at the resonance mass (see, e.g., [28]).

In order to explore the SU(3) multiplet structure of the resonance states we first study the 'heavy' SU(3) limit [1,27,28] with $m_{\pi,K,\eta} = 500$ MeV. For the anti-triplet $J^P = \frac{1}{2}^+$ states we use a somewhat arbitrary common mass $M = 2400$ MeV. In this case we obtain an anti-triplet of mass 2778 MeV with poles in the $(0, 0)$, $(\frac{1}{2}, -1)$ amplitudes and a sextet of mass 2900 MeV with poles in the $(1, 0)$, $(\frac{1}{2}, -1)$, $(0, -2)$ amplitudes. The result is reasonably stable against small variations of the optimal subtraction scale. Lowering the latter by 200 MeV reduces the anti-triplet and sextet masses by 40 MeV and 5 MeV only. Our finding reflects that the Weinberg–Tomozawa interaction (1),

$$\bar{3} \otimes 8 = \bar{3} \oplus 6 \oplus \overline{15} \tag{9}$$

is attractive in the anti-triplet and sextet channels but repulsive in the anti-quindecimplet channel. Similarly, using a common mass for the sextet $J^P = \frac{1}{2}^+$ states of 2500 MeV we obtain an anti-triplet of mass 2807 MeV with poles in the $(0, 0)$, $(\frac{1}{2}, -1)$ amplitudes, a sextet of mass 2875 MeV with poles in the $(1, 0)$, $(\frac{1}{2}, -1)$, $(0, -2)$ amplitudes and an anti-quindecimplet of mass 3000 MeV with poles in the $((\frac{1}{2}, 1), (0, 0), (1, 0), (\frac{1}{2}, -1), (\frac{3}{2}, -1), (1, -2))$ amplitudes. In this case the Weinberg–Tomozawa interaction (1),

$$6 \otimes 8 = \bar{3} \oplus 6 \oplus \overline{15} \oplus 24 \tag{10}$$

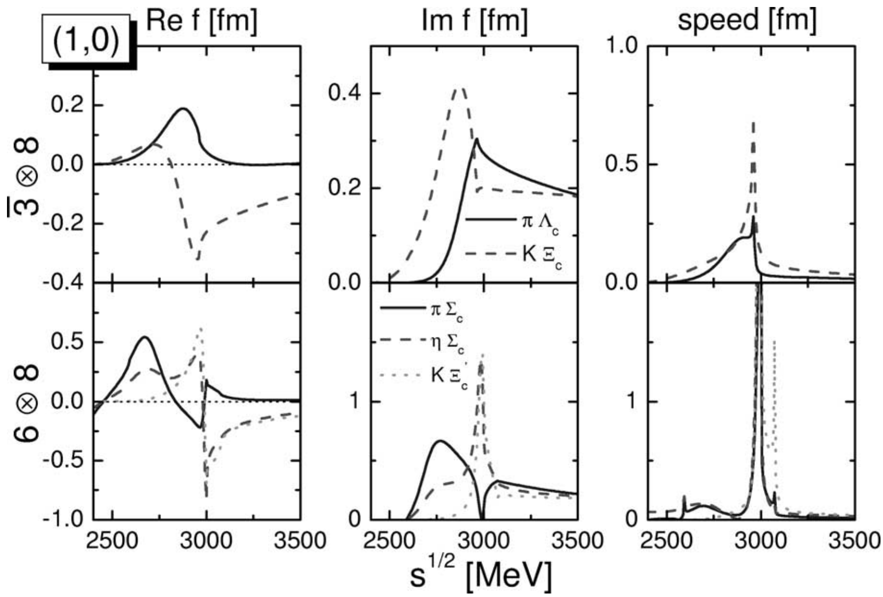

Fig. 2. Open-charm baryon resonances with $J^P = \frac{1}{2}^-$ and $(I, S) = (1, 0)$ as seen in the scattering of Goldstone bosons off anti-triplet $(\Lambda_c(2284), \Xi_c(2470))$ and sextet $(\Sigma_c(2453), \Xi_c'(2580), \Omega_c(2704))$ baryons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, f_{ab}, with $t_{ab} = f_{ab}(p_{\rm cm}^{(a)}\, p_{\rm cm}^{(b)})^{1/2}$ (see (8)).

predicts attraction in the anti-triplet, sextet and anti-quindecimplet channels but repulsion in the 24-plet channel.

In Figs. 1–4 the spectrum as it is predicted by the χ–BS(3) approach in terms of physical masses and the pion-decay constant $f = 90$ MeV is shown. In the upper (lower) panels the figures show the amplitudes and speed plots describing the scattering of Goldstone bosons off the anti-triplet (sextet) $J^P = \frac{1}{2}^+$ open-charm baryons. At leading order in the chiral expansion channels involving the anti-triplet and sextet open charm baryons decouple. We predict a bound state of mass 2767 MeV in the $(0,0)_{[\bar{3}]}$-sector (see Fig. 1). This state should be identified with the $\Lambda_c(2880)$ recently detected by the CLEO collaboration [10] via its decay into the $\pi\,\Sigma_c(2453) \to \Lambda_c\pi\pi$ channel. The narrow width of the observed state of smaller than 8 MeV [10] appears consistent with a suppressed coupling of that state to the $\pi\,\Sigma_c(2453)$ channel as predicted by chiral symmetry. In the lower panel of the figure the $(0,0)_{[6]}$-sector is presented. A resonance at about 2650 MeV that couples strongly to the $\pi\,\Sigma_c(2453)$ channel is predicted. The properties of this state are close to the ones of the $\Lambda_c(2593)$ resonance [8]. Given the fact that our computation is parameter-free this is a remarkable result. However, here we obtain a decay width which is significantly larger than the empirical width of about 4 MeV [8]. Chiral correction terms that couple the states seen in the lower and upper panels of Fig. 1 are expected to decrease this width. Level–level repulsion of the two observed states should lower the mass of the lighter state but push up the mass of the heavier state. A clear prediction of chiral-coupled channel dynamics

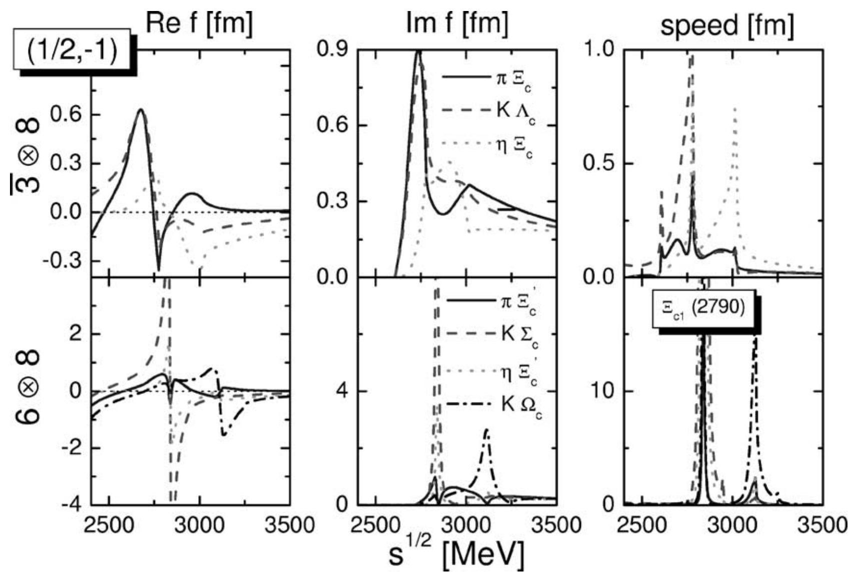

Fig. 3. Open-charm baryon resonances with $J^P = \frac{1}{2}^-$ and $(I, S) = (\frac{1}{2}, -1)$ as seen in the scattering of Goldstone bosons off anti-triplet $(\Lambda_c(2284), \Xi_c(2470))$ and sextet $(\Sigma_c(2453), \Xi_c'(2580), \Omega_c(2704))$ baryons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, f_{ab}, with $t_{ab} = f_{ab}(p_{\rm cm}^{(a)} p_{\rm cm}^{(b)})^{1/2}$ (see (8)).

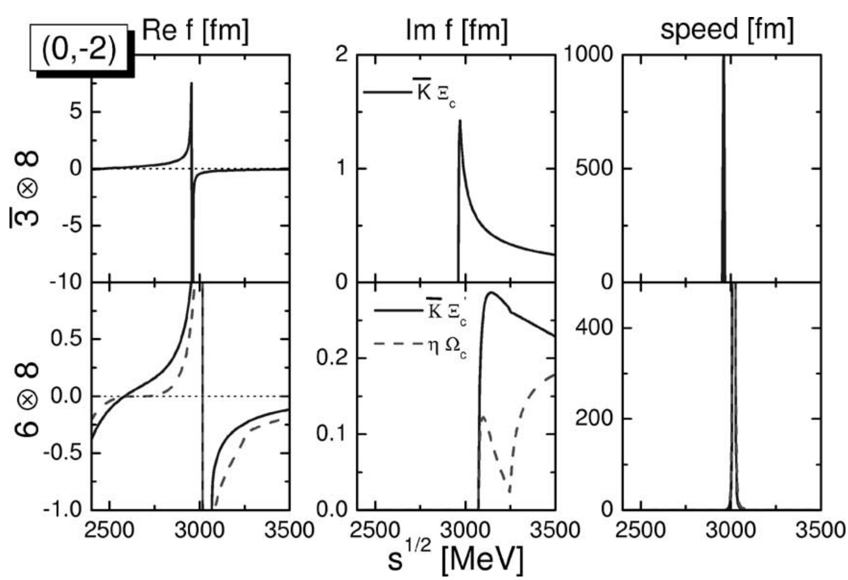

Fig. 4. Open-charm baryon resonances with $J^P = \frac{1}{2}^-$ and $(I, S) = (0, -2)$ as seen in the scattering of Goldstone bosons off anti-triplet $(\Lambda_c(2284), \Xi_c(2470))$ and sextet $(\Sigma_c(2453), \Xi_c'(2580), \Omega_c(2704))$ baryons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, f_{ab}, with $t_{ab} = f_{ab}(p_{\rm cm}^{(a)} p_{\rm cm}^{(b)})^{1/2}$ (see (8)).

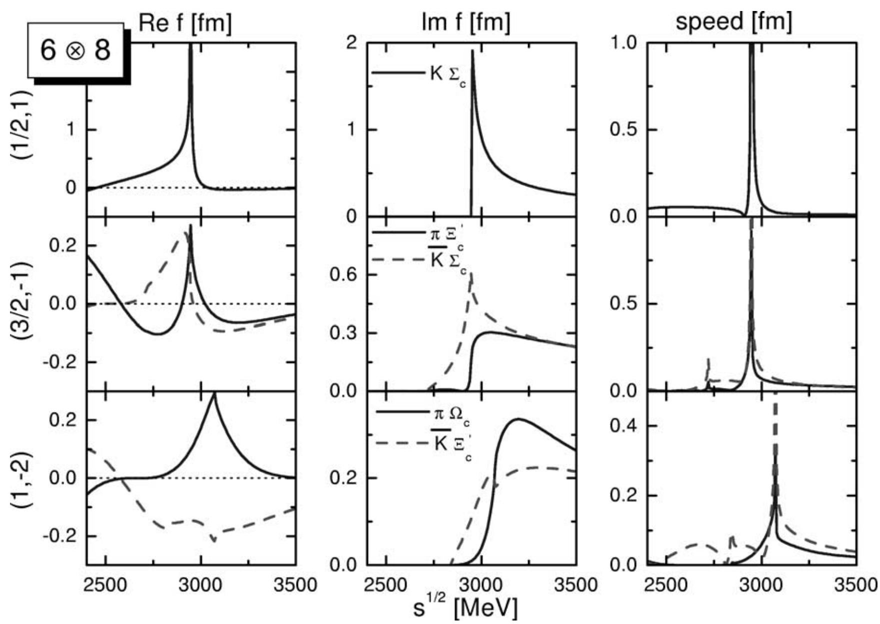

Fig. 5. Open-charm baryon resonances with $(I, S) = (\frac{1}{2}, 1)$, $(\frac{3}{2}, -1)$, $(1, -2)$ and $J^P = \frac{1}{2}^-$ as seen in the scattering of Goldstone bosons off sextet $(\Sigma_c(2453), \Xi_c(2580), \Omega_c(2704))$ baryons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, f_{ab}, with $t_{ab} = f_{ab}(p_{\rm cm}^{(a)} p_{\rm cm}^{(b)})^{1/2}$ (see (8)).

is an additional narrow resonance state in the $(0,0)_{[6]}$-sector at 3050 MeV as part of the anti-quindecimplet discussed above.

In Fig. 2 our predictions for excitations of the Σ_c baryon with $J^P = \frac{1}{2}^-$ are displayed. Two broad sextet states at about 2800 MeV that are eventually expected to mix and one narrow state at 2985 MeV part of the anti-quindecimplet are obtained. So far none of these states has been observed. In particular the latter state which couples strongly to the $\eta\Sigma_c$ and $K\Xi_c'$ channels may warrant a dedicated search for. In Fig. 3 the spectrum of the Ξ_c baryons with $J^P = \frac{1}{2}^-$ is shown. All together we expect five states. In the $(\frac{1}{2}, 1)_{[\bar{3}]}$-sector a clear resonance that couples strongly to the $\pi\Xi_c$ and $K\Lambda_c$ channels is seen at about 2750 MeV. At somewhat higher masses a weak resonance signal is suggested in the $\eta\Xi_c$ channel. In the lower panel of Fig. 3 two narrow states at 2830 MeV and 3120 MeV are shown. The lower state couples strongly to the $K\,\Sigma_c$ and $\eta\Xi_c$ channels and due to its small width should be identified with the $\Xi_c(2790)$ resonance [9]. The latter state has an empirical width smaller than 15 MeV [9]. The upper state has a strong coupling constant to the $K\,\Omega_c$ channel. One would expect the width of these states to increase somewhat once the subleading decay channels are incorporated. We do not find any clear signal of a third state in the lower panel of Fig. 3. We turn to Fig. 4 which probes with $(0, -2)$ exclusively the sextet states. Two distinct narrow states at 2959 MeV and 3016 MeV are anticipated. Finally in Fig. 5 our results for the channels $(\frac{1}{2}, 1)$, $(\frac{3}{2}, -1)$, $(1, -2)$ are displayed. Only members of the anti-quindecimplet manifest themselves as states in these channels. A narrow structure at 2947 MeV is predicted in the $(\frac{1}{2}, 1)$ channel. This state has a mass that is very close to the $K\Sigma_c$-threshold. Therefore in the speed plot of Fig. 5 it is difficult to separate the resonance signal from the square-root singularity induced by the opening of the corresponding channel. In the remaining channels $(\frac{3}{2}, -1)$, $(1, -2)$ no clear resonance signal is seen. The bound-state obtained in the 'heavy' SU(3) limit become rather broad resonances at around 3000 MeV.

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