Λ_c^* in a coupled-channel baryon-meson scattering*

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Abstract: Possible structures and properties of some excited states of Λ_c^+ dynamically generated in the coupled-channel P-wave meson-baryon scattering are studied by solving the Bethe-Salpeter(BS) equation in the framework of the Chiral Unitary Approach. It is shown that both $\Lambda_c^+(2765)$ and $\Lambda_c^+(2940)$ could be generated dynamically and could be compound states with multi-configuration molecular-like structures. The couplings of the generated states to various reaction channels are also calculated. Moreover, two highly excited states, $\Lambda_c^+(3024)$ and $\Lambda_c^+(3134)$, are predicted.

Key words: charmed baryon, multi-channel scattering, chiral unitary approach

PACS: 13.30.Eg, 12.40.Vv, 12.40.Yx **DOI:** 10.1088/1674-1137/35/1/004

1 Introduction

In recent years, many experimental groups have reported the observation of new hadron states, especially the excited states of charmed baryons, such as $\Lambda_c(2765)^+$, $\Lambda_c(2880)^+$, $\Lambda_c(2940)^+$, $\Sigma_c(2800)$, $\Xi_c(2815)^+$, $\Xi_c(2980)^+$, $\Xi_c(3077)^+$ and etc. [1–6]. Some of these states have a very narrow width and quite different properties to those predicted by the conventional quark model. It is apparent that an expanded mass spectra of charmed baryons would serve as a new resource for the understanding of the strong interaction and structure of hadron. This is because that the c quark, as an additional flavored quark, has a relatively larger mass than the light quark and, in particular, the strong decay of a narrow-width charmed baryon is OZI suppressed. Thus, the charmed baryon is an ideal laboratory for exploring the dynamics of a light quark in the heavy quark conditions circumstance and gaining detailed information about the strong interaction and structure of the hadron, even exotics. However, some characteristics and decay properties of observed charmed baryon excitations are still unknown or not clear, for example, the spin-parity (J^P) assignment and branch ratios of strong decays. Probing these natures both theoretically and experimentally is of great significance.

At this moment, we would focus here on some excited states of Λ_c. In 2007, the BaBar Collaboration claimed to have discovered a new narrow-width charmed baryon $\Lambda_c(2940)^+$ in the D^0p invariant mass spectrum by analyzing e^+e^- annihilation data [2]. The measured mass and width of the $\Lambda_c(2940)^+$ state are $M_{\Lambda_c(2940)^+} = 2939.8 \pm 1.3 \pm 1.0$ MeV and $\Gamma_{\Lambda_c(2940)^+} = 17.5 \pm 5.2 \pm 5.9$ MeV, respectively. As there is no structure around this mass in the D^+p invariant mass spectrum, the isospin of the state should be $I = 0$. Later, the Belle Collaboration confirmed this state in the $\Sigma_c(2455)^{0,++}\pi^{\pm}$ invariant mass spectrum in analyzing continuum e^+e^- annihilation data [3]. The reported mass and width of the $\Lambda_c(2940)^+$ state are $M_{\Lambda_c(2940)^+} = 2938.0 \pm 1.3^{+2.0}_{-4.0}$ MeV and $\Gamma_{\Lambda_c(2940)^+} = 13^{+8+27}_{-5-7}$ MeV, respectively. Up to now, the J^P value of this state is still not confirmed. Moreover, the nature of $\Lambda_c(2765)^+$ reported by the CLEO Collaboration in 2001 [6] is still not clear. The evidence of its existence is weak. Until now, no further experiment confirms its J^P value and whether it is a Λ_c^+ state or a Σ_c^+ state.

The theoretical interpretations for these states

Received 3 October 2010
* Supported by National Natural Science Foundation of China (10775147, 10975038, 11035006)
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have been intensively discussed. The results depend on the particular model applied. For instance, $\Lambda_c(2940)^+$ can be regarded as a conventional three-quark state with J^P being either $3/2^+$ or $5/2^-$ in a relativized quark model [7], or a $D^{*0}p$ molecular state with $J^P=1/2^-$ in a meson exchange model [8], or a conventional D-wave three-quark state with J^P being either $1/2^+$ or $3/2^+$ in studying its decay in the $^3P_0$ model [9], or a first orbital excitation of the light-diquark in $\Lambda_c(2286)^+$ or a $2S$ excitation of $\Sigma_c(2455)^+$ with J^P being $3/2^+$ in a c-quark-light-diquark model [10], or a D^*N molecular state with J^P being $1/2^+$ [11] and etc. As to the $\Lambda_c(2765)^+$ state, it can be considered as a conventional $1S$ three-quark state with J^P being either $1/2^+$, or a $2S$ state with $I(J^P)$ being $0(1/2^+)$ in a c-quark-light-diquark model [10] or in the heavy quark effective theory [12].

In this paper, we re-study these Λ_c excitations by solving a coupled-channel Bethe-Salpeter(BS) equation in the P-wave pseudoscalar meson-baryon scattering. In Sec. 2, we briefly introduce the formulism. We present the result and discussion in Sec. 3 and, finally, in Sec. 4 we give a short summary.

2 Brief formalism

The hadron-hadron scattering is an important process for extracting information about the inner structure of the hadron and the nature of strong interaction. The Chiral Unitary Approach (ChUA) [13, 14], which combines chiral dynamics with the unitarity of scattering processes, is an ideal framework with which this problem is treated. Extending the ChUA to a hadronic system with heavy flavor, we reach a so-called Heavy Chiral Unitary Approach (HChUA) [15, 16]. In terms of the HChUA, one can calculate the unitarized scattering amplitude for a coupled-channel scattering by solving the coupled-channel BS equation

$$T=[\mathrm{I}-VG]^{-1}V \tag{1}$$

and consequently obtain information about the bound state or resonance of the hadronic system studied. In Eq. (1), V stands for the kernel of the interaction, which can adopt the amplitude in the chiral perturbation theory in the lowest order, and G is a diagonal matrix with the element being a two-meson loop integral. In past years, the ChUA and the HChUA have widely been applied to study the hadronic system and have been successful in explaining the meson-meson and meson-baryon interactions, and dynamically generated mesons and baryons have been achieved [17, 18].

In this paper, we study the excited states of Λ_c with $J^P=1/2^+$ or $3/2^+$ as dynamically generated intermediate states in the scattering process with the HChUA. For this purpose, we choose a pseudoscalar meson-baryon scattering process with $I=0$, $S=0$, $C=+1$. There are eight reaction channels, $\Sigma_c\pi$, ND, $\Lambda_c\eta$, Ξ_cK, $\Xi_c'K$, ΛD_s, $\Lambda_c\eta'$ and $\Lambda_c\eta_c$, in this case. The thresholds for each channel are tabulated in Table 1.

Table 1. Thresholds for each channels (in MeV).

$\Sigma_c\pi$	ND	$\Lambda_c\eta$	Ξ_cK	$\Xi_c'K$	ΛD_s	$\Lambda_c\eta'$	$\Lambda_c\eta_c$
2592	2806	2834	2965	3072	3084	3244	5266

According to the meson exchange theory and the vector meson dominant model, a pseudoscalar meson can interact with a baryon via a vector meson exchange. The basic Feynman diagram is shown in Fig. 1.

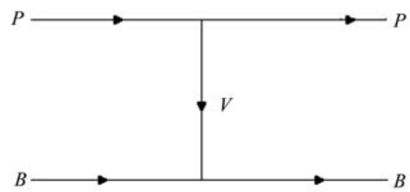

Fig. 1. Feynman diagram for pseudoscalar meson-baryon interaction via the vector meson exchange.

The $SU(4)$ pseudoscalar meson fields considered can be re-combined to follow multiplet fields [16]

$$\begin{aligned}\Phi_{[9]} &= \tau\cdot\pi(139)+\alpha^\dagger\cdot K(494)+K^\dagger(494)\cdot\alpha \\ &\quad +\eta(547)\lambda_8+\sqrt{\frac{2}{3}}\eta', \end{aligned} \tag{2}$$

$$\begin{aligned}\Phi_{[\bar{3}]} &= \frac{1}{\sqrt{2}}\alpha^\dagger\cdot D(1867)-\frac{1}{\sqrt{2}}D^t(1867)\cdot\alpha \\ &\quad +i\tau_2 D^{(s)}(1969),\end{aligned} \tag{3}$$

$$\Phi_{[1]} = \eta_c(2980). \tag{4}$$

In these equations, $\Phi_{[9]}$, $\Phi_{[\bar{3}]}$ and $\Phi_{[1]}$ denote the nonet light pseudoscalar fields, anti-symmetric triplet pseudoscalar fields with one charm flavor and the singlet pseudoscalar field with hidden charm flavor, respectively. It should be noted that the triplet pseudoscalar field with one charm flavor, $\Phi_{[3]}$, is not explicitly given because it will not appear in our considered problem. The matrices τ and α can be expressed by Gell-Mann matrices $\lambda_i(i=1,2,\cdots,8)$,

$$\tau=(\tau_1,\tau_2,\tau_3), \quad \alpha^\dagger=\frac{1}{\sqrt{2}}(\lambda_4+i\lambda_5,\lambda_6+i\lambda_7). \tag{5}$$

Similarly, the considered $SU(4)$ vector fields can be re-combined into following multiplet fields [16]

$$V_\mu^{[9]} = \tau\cdot\rho_\mu(770)+\alpha^\dagger\cdot \mathrm{K}_\mu(894)+\mathrm{K}_\mu^\dagger(894)\cdot\alpha + \left(\frac{2}{3}+\frac{1}{\sqrt{3}}\lambda_8\right)\omega_\mu(783) + \left(\frac{\sqrt{2}}{3}+\sqrt{\frac{2}{3}}\lambda_8\right)\phi_\mu(1020), \quad (6)$$

$$V_\mu^{[\bar{3}]} = \frac{1}{\sqrt{2}}\alpha^\dagger\cdot \mathrm{D}_\mu(2008)-\frac{1}{\sqrt{2}}\mathrm{D}_\mu^{\mathrm{t}}(2008)\cdot\alpha + \mathrm{i}\tau_2\mathrm{D}_\mu^{(\mathrm{s})}(2112) \quad (7)$$

with $V_\mu^{[9]}$ and $V_\mu^{[\bar{3}]}$ being the nonet and anti-triplet vector fields, respectively. The unused singlet vector field $V_\mu^{[1]}$ is not listed here. The $SU(4)$ baryon fields can also be re-arranged in the following multiplet fields:

$$\sqrt{2}B_{[8]} = \tau\cdot\Sigma(1193)+\alpha^\dagger\cdot\mathrm{N}(939)+\Xi^{\mathrm{t}}(1318)\cdot\alpha + \Lambda(1116)\lambda_8, \quad (8)$$

$$\sqrt{2}B_{[6]} = \frac{1}{\sqrt{2}}\alpha^\dagger\cdot\Xi_{\mathrm{c}}(2576)+\frac{1}{\sqrt{2}}\Xi_{\mathrm{c}}^{\mathrm{t}}(2576)\cdot\alpha + \Sigma_{\mathrm{c}}(2452)\cdot(\mathrm{i}\tau\tau_2) + \frac{\sqrt{2}}{3}(1-\sqrt{3}\lambda_8)\Omega_{\mathrm{c}}(2698), \quad (9)$$

$$\sqrt{2}B_{[\bar{3}]} = \frac{1}{\sqrt{2}}\alpha^\dagger\cdot\Xi_{\mathrm{c}}(2469)-\frac{1}{\sqrt{2}}\Xi_{\mathrm{c}}^{\mathrm{t}}(2469)\cdot\alpha + \mathrm{i}\tau_2\Lambda_{\mathrm{c}}(2285), \quad (10)$$

where $B_{[8]}$, $B_{[6]}$ and $B_{[\bar{3}]}$ represent the octet, sextet and anti-triplet baryon fields, respectively. The unused triplet field $B_{[3]}$ is not listed.

In our case, interactive Lagrangians for the pseudoscalar - pseudoscalar - vector and baryon - vector - baryon vertices can be written by referring to Ref. [16] as

$$\begin{aligned}\mathcal{L}_{\mathrm{PPV}} = & \frac{\mathrm{i}}{4}g\mathrm{Tr}\{2[(\partial_\mu\Phi_{[\bar{3}]})\Phi_{[\bar{3}]}^\dagger-\Phi_{[\bar{3}]}(\partial_\mu\Phi_{[\bar{3}]}^\dagger)]V_{[9]}^\mu \\ & -[(\partial_\mu\Phi_{[\bar{3}]})\Phi_{[\bar{3}]}^\dagger-\Phi_{[\bar{3}]}(\partial_\mu\Phi_{[\bar{3}]}^\dagger)]\mathrm{Tr}(V_{[9]}^\mu) \\ & +\frac{(m_{[9]}^{(V)})^2}{2g^2f^2}[(\partial_\mu\Phi_{[9]})\Phi_{[9]}-\Phi_{[9]}(\partial_\mu\Phi_{[9]})]V_{[9]}^\mu \\ & +2[(\partial_\mu\Phi_{[\bar{3}]}^\dagger)\Phi_{[9]}V_{[\bar{3}]}^\mu-\Phi_{[9]}(\partial_\mu\Phi_{[\bar{3}]})V_{[\bar{3}]}^{\dagger\mu}] \\ & +2[(\partial_\mu\Phi_{[9]})\Phi_{[\bar{3}]}V_{[\bar{3}]}^{\dagger\mu}-\Phi_{[\bar{3}]}^\dagger(\partial_\mu\Phi_{[9]})V_{[\bar{3}]}^\mu] \\ & +\sqrt{2}[(\partial_\mu\Phi_{[1]})\Phi_{[\bar{3}]}V_{[\bar{3}]}^{\dagger\mu}-\Phi_{[\bar{3}]}^\dagger(\partial_\mu\Phi_{[1]})V_{[\bar{3}]}^\mu] \\ & +\sqrt{2}[(\partial_\mu\Phi_{[\bar{3}]}^\dagger)\Phi_{[1]}V_{[\bar{3}]}^\mu-\Phi_{[1]}(\partial_\mu\Phi_{[\bar{3}]})V_{[\bar{3}]}^{\dagger\mu}] \\ & +[\Phi_{[\bar{3}]}^\dagger V_{[\bar{3}]}^\mu-\Phi_{[\bar{3}]}V_{[\bar{3}]}^{\dagger\mu}]\mathrm{Tr}(\partial_\mu\Phi_{[9]})\} \\ & +[(\partial_\mu\Phi_{[\bar{3}]})V_{[\bar{3}]}^{\dagger\mu}-(\partial_\mu\Phi_{[\bar{3}]}^\dagger)V_{[\bar{3}]}^\mu]\mathrm{Tr}(\Phi_{[9]})\}, \quad (11)\end{aligned}$$

and

$$\begin{aligned}\mathcal{L}_{\mathrm{BBV}} = & \frac{1}{2}g\mathrm{Tr}\{2\bar{B}_{[\bar{3}]}\gamma_\mu V_{[9]}^\mu B_{[\bar{3}]}+2\bar{B}_{[6]}\gamma_\mu V_{[9]}^\mu B_{[6]} \\ & +\bar{B}_{[8]}\gamma_\mu[V_{[9]}^\mu,B_{[8]}]+\bar{B}_{[8]}\gamma_\mu B_{[8]}\mathrm{Tr}(V_{[9]}^\mu) \\ & +\sqrt{2}[\bar{B}_{[8]}\gamma_\mu B_{[6]}V_{[\bar{3}]}^{\dagger\mu}+\bar{B}_{[6]}\gamma_\mu B_{[8]}V_{[\bar{3}]}^\mu] \\ & -\sqrt{6}[\bar{B}_{[8]}\gamma_\mu B_{[\bar{3}]}V_{[\bar{3}]}^{\dagger\mu}+\bar{B}_{[\bar{3}]}\gamma_\mu B_{[8]}V_{[\bar{3}]}^\mu]\}, \quad (12)\end{aligned}$$

where $f\approx 92$ MeV is the decay constant of a pion, $g\approx 6.6$ is the dimensionless vector coupling strength and $m_{[9]}^{(V)}$ denotes the mass of the light nonet vector meson.

From Lagrangians (11) and (12), we can easily write the scattering amplitude from the ith channel to the jth channel in the lowest order

$$V_{ij}(s,t,u) = C_{ij}\sqrt{\frac{E_i+M_i}{2M_i}}\sqrt{\frac{E_j+M_j}{2M_j}}\left\{2\sqrt{s}-M_j -M_i+\frac{(M_j-M_i)(m_j^2-m_i^2)}{m_{\mathrm{V}}^2}\right\}, \quad (13)$$

where s, t, and u are Mandelstam variables, and E_i, M_i and m_i denote the energy of the baryon and the masses of the baryon and meson in the ith channel, respectively. The coefficients C_{ij} are tabulated in Table 2 with

$$C_{44}=C_{55}=\frac{3m_\rho^2}{t-m_\rho^2}-\frac{m_\omega^2}{t-m_\omega^2}+\frac{2m_\phi^2}{t-m_\phi^2}.$$

Substituting $V_{ij}(s,t,u)$, as an integration kernel, into the BS Eq. (1) and using the expression of the L-wave amplitude

$$T_L(s)=\frac{1}{2}\int_{-1}^{1}\mathrm{d}\cos\theta P_L(\cos\theta)T(s,t,u), \quad (14)$$

where $P_L(\cos\theta)$ is the Legendre polynomial, to make the P-wave projection we obtain the total amplitude T of the pseudoscalar-baryon P-wave scattering, in which contributions from all orders are included. The matrix element of the diagonal matrix G in Eq. (1) is a loop integral composed of a meson propagator and a baryon propagator

$$G_{ii}=\mathrm{i}\int\frac{\mathrm{d}^4q}{(2\pi)^4}\frac{2M_i}{(k+p-q)^2-M_i^2+\mathrm{i}\epsilon}\frac{1}{q^2-m_{\mathrm{i}}^2+\mathrm{i}\epsilon} \quad (15)$$

with $p(k)$ being a four-momentum of the baryon (meson) in the ith reaction channel. To eliminate the divergence of the loop integration, a re-normalization procedure is usually performed with the so-called three-momentum cut-off method. The resultant analytical expression for G_{ii} reads [18]

$$G_{ii}(s) = \frac{2M_i}{16\pi^2 s}\left\{\sigma\left(\arctan\frac{s+\Delta}{\sigma\lambda_1}+\arctan\frac{s-\Delta}{\sigma\lambda_2}\right)\right.$$
$$-\left[(s+\Delta)\ln\left(\frac{q_{\max}}{M_i}(1+\lambda_1)\right)\right.$$
$$\left.\left.+(s-\Delta)\ln\left(\frac{q_{\max}}{m_i}(1+\lambda_2)\right)\right]\right\}, \qquad (16)$$

where $q_{\max}$ is the cut-off three-momentum and

$$\sigma = \{-[s-(M_i+m_i)^2][s-(M_i-m_i)^2]\}^{1/2},$$
$$\Delta = M_i^2-m_i^2,$$
$$\lambda_1 = \sqrt{1+\frac{M_i^2}{q_{\max}^2}},$$
$$\lambda_2 = \sqrt{1+\frac{m_i^2}{q_{\max}^2}}.$$

Table 2. Coefficients C_{ij} for the lowest order scattering amplitude $V_{ij}(s,t,u)$.

	$\Sigma_c\pi$	ND	$\Lambda_c\eta$	Ξ_cK	$\Xi_c'K$	ΛD_s	$\Lambda_c\eta'$	$\Lambda_c\eta_c$
$\Sigma_c\pi$	$\frac{1}{f^2}\frac{-m_\rho^2}{t-m_\rho^2}$	$\frac{\sqrt{6}g^2}{4}\frac{-1}{t-m_{D^*}^2}$	0	0	$\frac{\sqrt{3}}{4f^2}\frac{-m_{K^*}^2}{t-m_{K^*}^2}$	0	0	0
ND		$\frac{3g^2}{4}\left(\frac{-1}{t-m_\rho^2}+\frac{-1}{t-m_\omega^2}\right)$	$\frac{\sqrt{2}g^2}{4}\frac{1}{t-m_{D^*}^2}$	0	0	$\frac{\sqrt{3}g^2}{2}\frac{1}{t-m_{K^*}^2}$	$\frac{g^2}{2}\frac{1}{t-m_{D^*}^2}$	$\frac{\sqrt{3}g^2}{2}\frac{-1}{t-m_{D^*}^2}$
$\Lambda_c\eta$			0	$\frac{\sqrt{3}}{4f^2}\frac{-m_{K^*}^2}{t-m_{K^*}^2}$	0	$\frac{g^2}{\sqrt{6}}\frac{1}{t-m_{D_s^*}^2}$	0	0
Ξ_cK				$\frac{-1}{8f^2}C_{44}$	0	$\frac{\sqrt{2}g^2}{4}\frac{1}{t-m_{D^*}^2}$	0	0
$\Xi_c'K$					$\frac{-1}{8f^2}C_{55}$	$\frac{\sqrt{6}g^2}{4}\frac{1}{t-m_{D^*}^2}$	0	0
ΛD_s						$\frac{g^2}{2}\frac{-1}{t-m_\phi^2}$	$\frac{\sqrt{3}g^2}{6}\frac{-1}{t-m_{D_s^*}^2}$	$\frac{g^2}{2}\frac{1}{t-m_{D_s^*}^2}$
$\Lambda_c\eta'$							0	0
$\Lambda_c\eta_c$								0

In the numerical calculation, $q_{\max}$ is a unique adjustable parameter, which is taken to be 0.8 and 1.0 GeV for comparison.

3 Numerical results and discussion

The dynamically generated state is closely related to the singularity of the scattering amplitude. The position of the singular point (namely, the pole) of the amplitude T can be determined by

$$\det[I-VG]=0. \qquad (17)$$

In order to discover the dynamically generated state in the scattering process, we perform an analytic continuation for the scattering amplitude in the complex energy plane. In the case of N reaction channels, there are N thresholds along the positive real energy axis, namely $s_i^{\rm th}=E_i^2(i=1,2,...,N)$, and consequently $2^{\rm N}$ Riemann sheets. Each Riemann sheet can be defined by the sign of Im p_i with p_i being the momentum of the scattered particle in the ith channel in the center of mass frame

$$p_i=\sqrt{\frac{[s-(M_i+m_i)^2][s-(M_i-m_i)^2]}{4s}}. \qquad (18)$$

We list first few important Riemann sheets in the following:

RS1: Im $p_i>0$, $(i=1,2,\cdots,8)$
RS2: Im $p_1<0$, Im $p_j>0$, $(j=2,\cdots,8)$
RS3: Im $p_i<0$, Im $p_j>0$, $(i=1,2;\ j=3,\cdots,8)$
⋮
RS9: Im $p_i<0$, $(i=1,2,\cdots,8)$

3.1 Poles on Riemann sheets

The resultant poles of T on various Riemann sheets are tabulated in Table 3, where the cut-off parameter is taken to be $q_{\max}$ = 0.8 and 1.0 GeV, respectively. From the data in the table, it can be shown that the real and imaginary parts of the pole found on each Riemann sheet are not sensitive to the adopted value of the cut-off parameter $q_{\max}$. Therefore, in the following discussion, we only consider the result with $q_{\max}=0.8$ GeV.

In a Riemann sheet with proper energy, a physical state corresponds to a pole of the scattering amplitude. Because the first Riemann sheet is physical one, only the pole located along the real axis, namely the imaginary part of the pole is zero or very close to zero, is meaningful. This kind of pole corresponds to a stable state or a bound state. The other Riemann Sheets are all non-physical. We show the masses and widths that correspond to the obtained poles in the pseudoscalar-baryon P-wave scattering with $I=0$, $C=+1$, $S=0$ in Table 3, poles which do not have physical meaning are not listed.

Table 3. Masses and widths of pole states(GeV).

$q_{\max}$	0.8		1.0	
	M_{Λ_c}	$\Gamma_{\Lambda_c}/2$	M_{Λ_c}	$\Gamma_{\Lambda_c}/2$
RS2	2.76115	0.05409	2.74809	0.06331
RS4	2.93443	0.00988	2.93604	0.01030
RS5	3.02787	0.04668	3.02401	0.0469256
RS7	3.13011	0.05521	3.13428	0.06816

On the first Riemann sheet, all the appeared poles are not located along the real axis or near the real axis, its imaginary part is large and, consequently, have no physical meaning. Similarly, all the appeared poles on the third, sixth and eighth Riemann sheets have no physical meaning. Therefore, these poles are not listed in Table 3. We now consider the pole on the second Riemann sheet: (1) If the real part of the pole is less than the first threshold of $M_{\pi\Sigma_c}=2.592$ GeV, no channel is opened. Additionally, if the imaginary part of the pole is not zero (or close to zero), this pole does not correspond to a state in the physical space; (2) If the real part of the pole is less than the second threshold of $M_{DN}=2.806$ GeV and greater than the first threshold, the $\pi\Sigma_c$ channel is open. If the imaginary part of the pole has a certain size simultaneously, this pole might relate to a resonance of Λ_c and can decay into π and Σ_c; (3) If the real part of the pole is greater than the second threshold of $M_{DN}=2.806$ GeV, the pole does not correlate to the state in this physical space. In our calculation, we have only one pole whose real part is between the first and second thresholds and whose imaginary part has a certain size. This pole might correspond to $\Lambda_c(2765)^+$ with $J^P=1/2^+$ or $3/2^+$, $M^{(1)}_{\Lambda_c}=2.761$ GeV and $\Gamma^{(1)}_{\Lambda_c}=0.108$ GeV. The structure of this state is possibly a unstably bound ND molecular-like state with a sizable configuration of $\Sigma_c\pi$, because it is close to the threshold of the $\Sigma_c\pi$ channel and has a certain size of width. On the fourth Riemann sheet, there is a pole which is fairly close to the Ξ_cK threshold and has a narrow width. This pole may closely related to $\Lambda_c(2940)^+$ with $J^P=1/2^+$ or $3/2^+$. Its mass and width are about $M^{(2)}_{\Lambda_c}=2.934$ GeV and $\Gamma^{(2)}_{\Lambda_c}=0.019$ GeV, respectively. These results are compatible with the recent data in Refs. [2, 3]. The structure of this state is possibly an almost bound Ξ_cK molecular-like state with a certain amount of $\Lambda_c\eta$, ND and $\Sigma_c\pi$ components. Moreover, we also find a pole on the fifth Riemann sheet and label the corresponding state as $\Lambda_c(3028)^+$. This state could be a unstable Ξ'_cK bound state with sizable components of Ξ_cK, $\Lambda_c\eta$, ND and $\Sigma_c\pi$. Its quantum number is $J^P=1/2^+$ or $3/2^+$, and the mass and width are around $M^{(3)}_{\Lambda_c}=3.028$ GeV and $\Gamma^{(3)}_{\Lambda_c}=0.093$ GeV, respectively. Finally, we have a pole on the seventh Riemann sheet. The corresponding state is labeled as $\Lambda_c(3130)^+$. The structure of the state can be regarded as a wide-width ΛD_s resonance with a certain amount of Ξ'_cK, Ξ_cK, $\Lambda_c\eta$, ND and $\Sigma_c\pi$ components. Its J^P value could be $1/2^+$ or $3/2^+$. Its mass and width are about $M^{(4)}_{\Lambda_c}=3.130$ GeV and $\Gamma^{(4)}_{\Lambda_c}=0.110$ GeV, respectively.

3.2 Couplings between Λ_c^+'s and channels

To study the decay property of the dynamically generated states and, consequently, extract the effective coupling constants between the dynamically generated states and various reaction channels, we make a Laurent expansion for the scattering amplitude T_{ij} near the pole [13]

$$T_{ij}=\frac{g_ig_j}{s-s_{\text{pole}}}+\gamma_0+\gamma_1(s-s_{\text{pole}})+\cdots, \tag{19}$$

where g_i and g_j are effective coupling constants between the dynamically generated state and the ith and jth channels, respectively, s_{pole} represents the dynamically generated pole, γ_0, γ_1, $\cdots$ are Laurent expansion coefficients. Then, g_ig_j can be obtained by calculating the residues of corresponding poles [14]

$$g_ig_j=\lim_{s\to s_{\text{pole}}}(s-s_{\text{pole}})T_{ij}. \tag{20}$$

If the pole is not a bound state pole, an analytic continuation of the amplitude should be made up the corresponding Riemann sheet, the residue can then be calculated. The analytically extended amplitude on the a^{th} Riemann sheet can be written as

$$T^{(a)}_{ij}(w)=\frac{N^{(a)}_{ij}(w)}{D^{(a)}_{ij}(w)}, \tag{21}$$

with

$$D^{(a)}_{ij}(w)=|I-V(w)G^{(a)}(w)| \tag{22}$$

and $w=\sqrt{s}$. Finally, the effective coupling constants can be calculated by

$$g_ig_j=\frac{2wN^{(a)}_{ij}(w)}{\dfrac{\mathrm{d}}{\mathrm{d}w}D^{(a)}_{ij}(w)}\Big|_{w\to w_{\text{pole}}}, \tag{23}$$

where w_{pole} specifies the position of the pole. The resultant coupling constants are tabulated in Table 4.

It is shown that $\Lambda^{(1)}_c(2761)$ has relatively stronger couplings to the ND, $\Lambda_c\eta$ and Ξ_cK channels, but whether it would decay in these modes also depends on the phase space. Similarly, $\Lambda^{(2)}_c(2934)$ has sizable couplings with the ND, Ξ_cK and $\Sigma_c\pi$ channels,

$\Lambda_c^{(3)}(3028)$ has visible couplings with the ΛD_s, ND and $\Xi_c' K$ channels, and $\Lambda_c^{(4)}(3130)$ has noticeable couplings with the ND, $\Lambda_c\eta'$ and $\Xi_c' K$ channels.

Table 4. $|g_i|$ in $I=0,\ S=0,\ C=+1$ case (GeV).

channel	$\Sigma_c\pi$	ND	$\Lambda_c\eta$	$\Xi_c K$	$\Xi_c^{'} K$	ΛD_s	$\Lambda_c\eta'$	$\Lambda_c\eta_c$
$\Lambda_c^{(1)}(2761)$	1.432	2.896	2.310	2.045	0	0	0	0
$\Lambda_c^{(2)}(2934)$	1.370	2.499	1.176	1.991	0.7895	0	0	0
$\Lambda_c^{(3)}(3028)$	1.275	2.334	1.420	1.543	2.067	2.802	0	0
$\Lambda_c^{(4)}(3130)$	2.679	4.246	0.9703	2.758	3.565	2.871	3.615	0

4 Summary

Starting from chiral Lagrangian in the lowest order (11) and (12), we study the pseudoscalar-baryon P-wave scattering in the $I=0,\ S=0,\ C=+1$ sector. By searching for the poles on various Riemann sheets, we find four non-strange charmed states labeled as $\Lambda_c^{(i)}$ $(i=1,2,3,4)$. These states are all dynamically generated from the pseudoscalar-baryon P-wave scattering process and have the same quantum number $1/2^+$ or $3/2^+$. Their masses and widths are $M_{\Lambda_c}^{(1)}=2.761$ GeV, $\Gamma_{\Lambda_c}^{(1)}=0.108$ GeV, $M_{\Lambda_c}^{(2)}=2.934$ GeV, $\Gamma_{\Lambda_c}^{(2)}=0.019$ GeV, $M_{\Lambda_c}^{(3)}=3.028$ GeV, $\Gamma_{\Lambda_c}^{(3)}=0.093$ GeV, $M_{\Lambda_c}^{(4)}=3.130$ GeV and $\Gamma_{\Lambda_c}^{(4)}=0.110$ GeV, respectively. The result shows that $\Lambda_c^{(1)}$ is close to $\Lambda_c(2765)^+$, its mass matches the data of $\Lambda_c(2765)^+$ but its width is much larger than the observed one. Since the experimental status for $\Lambda_c(2765)^+$ is poor, we need more accurate measurement on the mass, width and decay branching ratios to specify the nature of observed $\Lambda_c(2765)^+$ and more theoretical studies on its decays. $\Lambda_c^{(2)}$ matches with the mass and width data of $\Lambda_c(2940)^+$ quit well. It is also compatible with many theoretical predictions mentioned in the Introduction, above. However, whether its J^P is $1/2^+$ or $3/2^+$ still requires more studies on its decays, both theoretically and experimentally. In this calculation, we further predict two possible new Λ_c's, $\Lambda_c(3028)^+$ and $\Lambda_c(3130)^+$. Combined with the couplings calculated above, one may guess that it would be possible to find them in the ND decay mode.

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