

Identifying the two-pole structure of the $\Lambda(1405)$ using an SU(3) flavor filter

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We propose a novel method to identify the two-pole structure of the $\Lambda(1405)$. The two poles owe their origin to different quark flavor irreducible representations in the meson-baryon coupled-channel interactions, thus they should be individually manifested in reactions that provide good flavor eigenstate sources. Hadronic decays of charmonia into $\bar{\Lambda}\Sigma\pi$ and $\bar{\Lambda}(1520)\Sigma\pi$ are such reactions, and the flavor octet and singlet poles can be approximately singled out in these two decay modes. This SU(3) flavor filter works even considering the flavor symmetry breaking. With the huge charmonium data sets collected, it is therefore promising to solve the long-standing $\Lambda(1405)$ puzzle employing the proposed flavor filter.

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Introduction. The $\Lambda(1405)$ is a hyperon resonance with quantum numbers $J^P = 1/2^-$ and $I = 0$. It was first predicted as a $\bar{K}N$ bound state by Dalitz and Tuan [1] before it was observed in 1961 [2]. Although it is currently quoted by the Particle Data Group (PDG) [3] as a four-star resonance, its underlying structure is still under debate. In the constituent quark model, it is difficult to assign the $\Lambda(1405)$ as the first orbital excited state of the Λ hyperon, because its mass is much lower than its nucleon counterpart $N(1535)$, and the mass splitting between the $\Lambda(1405)$ and its spin partner $\Lambda(1520)$ with $J^P = 3/2^-$ is much larger than that in the nucleon sector. Thus, it has been a long-standing puzzle what the internal structure of the $\Lambda(1405)$ is. The modern theory concerning the $\Lambda(1405)$ is the

coupled-channel theory constrained by chiral and SU(3) flavor symmetries as well as unitarity [4–8]. In this theory, the $\Lambda(1405)$ is interpreted as a dynamically generated state, or a meson-baryon molecular state. Another fascinating finding in the coupled-channel framework is the two-pole structure, i.e., the single $\Lambda(1405)$ four-star resonance listed in the review of particle physics (RPP) in versions up to 2018 [9], actually corresponds to two distinct poles on the unphysical Riemann sheet of the complex energy plane [7,10–15], which are now listed as two items $\Lambda(1380)$ and $\Lambda(1405)$, respectively, in the latest version of RPP [3]. The two pole positions were updated in Ref. [16] using the state-of-the-art next-to-next-to-leading order chiral potential. We refer to Refs. [3,17–19] for recent reviews about the $\Lambda(1405)$ state.

The two-pole structure of the $\Lambda(1405)$ can be well understood by means of the SU(3) flavor symmetry and group theory [20]. For the scattering of Nambu-Goldstone (NG) bosons off the ground-state light baryons, the Weinberg-Tomozawa (WT) term [5,8] gives the most important contribution at low energies. The WT potential projected to the S wave is given by

$$V_{ij}^{\text{WT}}(\sqrt{s}) = -\frac{C_{ij}}{4f^2} (2\sqrt{s} - M_i - M_j) \mathcal{N}_i \mathcal{N}_j, \quad (1)$$

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with the normalization factor $\mathcal{N}_i = \sqrt{(M_i + E_i)/2M_i}$, where M_i and E_i are the baryon mass and energy in channel i , and f is the NG boson decay constant. The C_{ij} coefficients reflect the interaction strength of pertinent channels, which are calculated from the leading order (LO) meson-baryon chiral Lagrangian [5]. For the $I = 0$ and $S = -1$ sector of interest here, there are four coupled channels: $\pi\Sigma$, $\bar{K}N$, $\eta\Lambda$, and $K\Xi$. In the SU(3) limit, the product of the NG boson octet and the light baryon octet can be decomposed into irreducible representations. The decomposition reads

$$8 \otimes 8 = 1 \oplus 8_s \oplus 8_a \oplus 10 \oplus \bar{10} \oplus 27, \quad (2)$$

where the subscripts “s” and “a” refer to the symmetric and antisymmetric representations, respectively [20]. The WT potential can also be given in a basis of the SU(3) states, and the corresponding coefficient matrix is then diagonalized as

$$C_{\alpha\beta}^{\text{SU}(3)} = \sum_{i,j} \mathcal{D}_{ai} C_{ij} \mathcal{D}_{\beta j} = \text{diag}(6, 3, 3, 0, 0, -2), \quad (3)$$

where $\mathcal{D}_{ai}$ and $\mathcal{D}_{\beta j}$ represent the SU(3) Clebsch-Gordan coefficients and the index $\alpha(\beta)$ runs over the **1**, **8_s**, **8_a**, **10**, **$\bar{10}$** , and **27** representations [20,21]. The crucial observation from the above equation is that the meson-baryon interaction is attractive in the singlet and octet representations, with the corresponding coefficients being positive, and correspondingly one pole is dynamically generated in each of these three representations. This coupled-channel chiral dynamics [7,20,22,23] is the origin of the two-pole structure. The $\Lambda(1380)$ and $\Lambda(1405)$ are found to be connected to the dynamically generated singlet and octet (**8_s**) poles, respectively; the assignment could be interchanged at the next-to-leading order (NLO) [15] (see Ref. [24] for an alternative analysis but yielding conclusions similar to those obtained at LO). The degeneracy of the two octets in the SU(3) limit in Eq. (3) is broken at NLO [15]. Using physical masses of the baryons and mesons, the **8_a** octet pole evolves into the $\Lambda(1680)$ resonance.

Besides the $\Lambda(1405)$, there are some other candidates that may also have two-pole structures, such as $K_1(1270)$ [25,26], $D_0^*(2300)$, $D_1(2430)$ [27–30], and $\Xi(1820)$ [31]. For instance, the $D_0^*(2300)$, similar to the $\Lambda(1405)$, is supposed to be dynamically generated from the NG boson octet scattering off the ground state charmed meson antitriplet. As a result, not one single pole, but one sextet and one antitriplet D_0^* pole exist in the complex energy plane [30]. As stressed in, e.g., Refs. [19,22,23], the emergence of two-pole structures is very common in the coupled-channel chiral dynamics. Since the positions of the two poles are usually close to each other, one may only see a single effective resonant structure in the pertinent invariant mass spectrum. Subsequently, one important question is how the two poles living in the same channels

can be unambiguously identified in experiments. For the two-pole structure of $\Lambda(1405)$, considering that one pole couples more strongly to the $\pi\Sigma$ channel, while the other couples more strongly to the $\bar{K}N$ channel, the authors in Ref. [20] suggested that it may be possible to find out the existence of the two resonances by performing different experiments. However, this two-pole structure has still not been satisfactorily experimentally resolved up to now, though there are some hints [32]. We refer to Ref. [19] for a concise review concerning the two-pole structure phenomena. In this Letter, we propose a novel method to identify the two-pole structures using an SU(3) flavor filter.

Heavy quarkonium decays. The predicted masses of both $\Lambda(1380)$ and $\Lambda(1405)$ are lower than the $\bar{K}N$ threshold, therefore one can only observe both resonances simultaneously in the $\pi\Sigma$ spectrum. We consider the three-body decay process $Y \rightarrow \bar{\Lambda}\Sigma\pi$, where Y represents a charmonium or bottomonium state with suitable quantum numbers. If the final state $\pi\Sigma$ is produced from an intermediate hyperon decay, since Y is an SU(3) singlet while the $\bar{\Lambda}$ is an SU(3) flavor octet member, the intermediate state producing the final $\pi\Sigma$ has to belong to an SU(3) octet in the SU(3) symmetric limit. Likewise, for the process $Y \rightarrow \bar{\Lambda}(1520)\Sigma\pi$, since $\bar{\Lambda}(1520)$ is generally supposed to be an SU(3) singlet with $J^P = 3/2^-$ [33–37], the intermediate state producing the final $\pi\Sigma$ has to belong to an SU(3) singlet in this case. Consequently, it is possible to single out the flavor octet and singlet poles in the $Y \rightarrow \bar{\Lambda}\Sigma\pi$ and $Y \rightarrow \bar{\Lambda}(1520)\Sigma\pi$ process, respectively. Considering that the SU(3) flavor symmetry is an approximate symmetry in quantum chromodynamics (QCD) with corrections of about $m_s/\Lambda_{\text{QCD}} \approx 30\%$, where m_s is the strange quark mass and Λ_{QCD} is the nonperturbative QCD scale, the above proposed filter should work well.

Such a flavor filter is feasible at currently running experiments. For instance, the BESIII experiment has collected around 10×10^9 J/ψ events and 3×10^9 $\psi(2S)$ events [38]. In a future Super Tau-Charm Facility, 3.4×10^{12} J/ψ or 6.4×10^{11} $\psi(2S)$ events per year are expected [39]. Furthermore, the branching fractions of relevant decay channels to be discussed here are quite sizable [3], and thus a clear identification of the two-pole structure of the $\Lambda(1405)$ can be expected.

Formalism and results. A diagrammatic description for the $Y \rightarrow \bar{\Lambda}\Sigma\pi$ and $Y \rightarrow \bar{\Lambda}(1520)\Sigma\pi$ decays is presented in Fig. 1. The final states can be produced either through a direct short-distance process, as shown in Fig. 1(a), or via a rescattering mechanism involving the meson-baryon final-state interaction (FSI), as illustrated in Fig. 1(b). The meson-baryon FSI will be treated using the chiral unitary approach, and the two $\Lambda(1405)$ poles will be dynamically generated in the coupled-channel chiral dynamics. The χ_{c0} decays into $\bar{\Sigma}\Sigma\pi$ and $\bar{\Lambda}\Sigma\pi$ have been studied in Refs. [40,41]. However, the flavor filter to identify the two-pole structure of $\Lambda(1405)$ has not been discussed.

The short-distance decay vertex of J/ψ or $\psi(2S)$ into a baryon-antibaryon pair accompanied by a pseudoscalar meson can be parametrized by the effective Lagrangian

$$\mathcal{L}_\psi^I = \tilde{D} \langle \bar{B} \gamma_\mu \gamma_5 \{ \Phi, B \} \rangle \psi^\mu + \tilde{F} \langle \bar{B} \gamma_\mu \gamma_5 [\Phi, B] \rangle \psi^\mu, \quad (4)$$

where Φ and B represent the ordinary SU(3) matrices for the pseudoscalar meson octet and light baryon octet, respectively, the symbol $\langle \dots \rangle$ stands for the trace in the flavor space, and ψ^μ represents the vector charmonium field. The coupling constants $\tilde{D}$ and $\tilde{F}$ can be determined from experimental measurements. For the vector charmonia decays in this work, we use the PDG branching ratios [3] of four channels $\bar{\Lambda}\Sigma\pi$, $\bar{\Lambda}N\bar{K}$, $\bar{\Lambda}\Lambda\eta$, and $\bar{\Sigma}N\bar{K}$ to fit the coupling constants in Eq. (4). The $\pi\Sigma$ line shape depends on the ratio $\mathcal{R}_{\tilde{F}/\tilde{D}}$ of the two coupling constants. For the J/ψ and $\psi(2S)$ decays, the fit gives $\mathcal{R}_{\tilde{F}/\tilde{D}} = 0.18 \pm 0.03$ and $\mathcal{R}_{\tilde{F}/\tilde{D}} = 0.50 \pm 0.06$, respectively.

For the vector charmonia decaying into $\bar{\Lambda}(1520)B\Phi$, the effective Lagrangian can be constructed as

$$\mathcal{L}_\psi^{\text{II}} = g_0 \bar{\Lambda}_\mu \gamma_5 \langle \Phi B \rangle \psi^\mu, \quad (5)$$

where the Rarita-Schwinger field Λ_μ represents the SU(3) singlet spin-3/2 hyperon $\Lambda(1520)$, and g_0 is the coupling constant. There are currently no data concerning J/ψ and $\psi(2S)$ decaying into $\bar{\Lambda}(1520)B\Phi$. However, since only one single coupling g_0 appears in $\mathcal{L}_\psi^{\text{II}}$, the $\pi\Sigma$ line shape will not be affected by the g_0 value.

Taking into account the meson-baryon FSI using the chiral unitary approach, the amplitude for the decay of a charmonium into channel i can be written as

$$t_i = \tilde{V}_i + \sum_j \tilde{V}_j G_j T_{ji}, \quad (6)$$

where $\tilde{V}_i$ corresponds to the short-distance decay amplitude shown in Fig. 1(a) and can be derived from Eqs. (4) and (5). The index j runs over four coupled channels: $\pi\Sigma$, $\bar{K}N$, $\eta\Lambda$, and $K\Xi$. According to the Lagrangian $\mathcal{L}_\psi^{\text{I}}$, for $\psi \rightarrow \bar{\Lambda}B\Phi$, we have $\tilde{V}_i = h_i \bar{u}_B \gamma_\mu \gamma_5 v_\Lambda \epsilon_\mu^\nu$, and the effective coupling constant h_i for each channel reads

$$\begin{aligned} h_{\pi\Sigma} &= -\sqrt{2}\tilde{D}, & h_{\bar{K}N} &= -\sqrt{\frac{1}{3}}\tilde{D} - \sqrt{3}\tilde{F}, \\ h_{\eta\Lambda} &= -\sqrt{\frac{2}{3}}\tilde{D}, & h_{K\Xi} &= \sqrt{\frac{1}{3}}\tilde{D} - \sqrt{3}\tilde{F}. \end{aligned} \quad (7)$$

For $\psi \rightarrow \bar{\Lambda}(1520)B\Phi$, the effective couplings read

$$\begin{aligned} h_{\pi\Sigma} &= -\sqrt{3}g_0, & h_{\bar{K}N} &= \sqrt{2}g_0, \\ h_{\eta\Lambda} &= g_0, & h_{K\Xi} &= -\sqrt{2}g_0. \end{aligned} \quad (8)$$

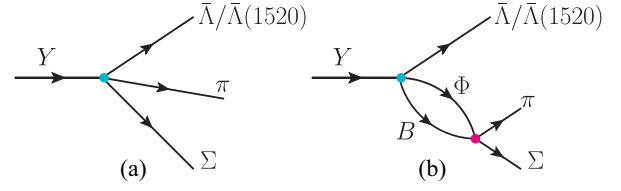


FIG. 1. Diagrams for the heavy quarkonium state Y decaying into $\bar{\Lambda}\Sigma\pi$ or $\bar{\Lambda}(1520)\Sigma\pi$: (a) short-distance process, and (b) rescattering process. Here, Φ and B represent the intermediate NG boson and light baryon, respectively.

The unitarized scattering amplitude T_{ij} takes the form

$$T_{ij} = V_{ij} + V_{ik} G_k T_{kj}. \quad (9)$$

The loop function G_k can be regularized using the dimensional regularization. For the LO WT interaction, we adopt the same subtraction constants as those in Refs. [8,20], i.e., $a_{\pi\Sigma} = -2.00$, $a_{\bar{K}N} = -1.84$, $a_{\eta\Lambda} = -2.25$, and $a_{K\Xi} = -2.67$, and use them for the loop function in Eq. (6) as well. With these parameters, one finds two poles located in the vicinity of 1.4 GeV: Pole 1 at $(1390 - i66)$ MeV and pole 2 at $(1426 - i16)$ MeV [20]. The physical masses of pertinent mesons and baryons are used to obtain the above predictions, which implies that the SU(3) symmetry is already broken. The analysis in Ref. [20] shows that pole 1 and pole 2 develop from the singlet and octet ($\mathbf{8}_s$) poles, respectively. For clarity, we name pole 1 and pole 2 as $\Lambda(1380)$ and $\Lambda(1405)$ in the following discussion, although their positions vary a bit in different approaches. Since SU(3) is a good approximate symmetry and only broken at the level of 30% or so, we expect that in the physical situation the main components of these poles are still SU(3) singlet and octet, respectively.

The $\pi\Sigma$ invariant mass spectra for the two J/ψ decay modes are displayed in Fig. 2. First we focus on the LO results (solid lines). From Fig. 2(a), it can be seen that the narrower octet pole $\Lambda(1405)$ is much more prominent compared with the broader singlet pole $\Lambda(1380)$. This implies that the $\Lambda(1405)$ state can be effectively singled out in the $J/\psi \rightarrow \bar{\Lambda}\Sigma\pi$ decay. On the other hand, for the $J/\psi \rightarrow \bar{\Lambda}(1520)\Sigma\pi$ decay, as expected, the SU(3) singlet pole $\Lambda(1380)$ is much more prominent compared with the octet pole, as can be seen from Fig. 2(b). As a result, we can effectively single out the singlet $\Lambda(1380)$ state in such a process.

Moreover, the signal for the octet pole $\Lambda(1405)$ is shielded beneath a dip around the $N\bar{K}$ threshold. This is a consequence of the universal mechanism for the emergence of a dip proposed in Ref. [42]: For the $\Lambda(1520)$ being an SU(3) singlet, the $B\Phi$ pair produced from the $J/\psi \rightarrow \bar{\Lambda}(1520)B\Phi$ must be largely an SU(3) singlet, and the octet enters only as an intermediate state via SU(3) breaking; in that case, the strong S -wave attraction between $N\bar{K}$ in the octet naturally leads to a dip (for a recent application of

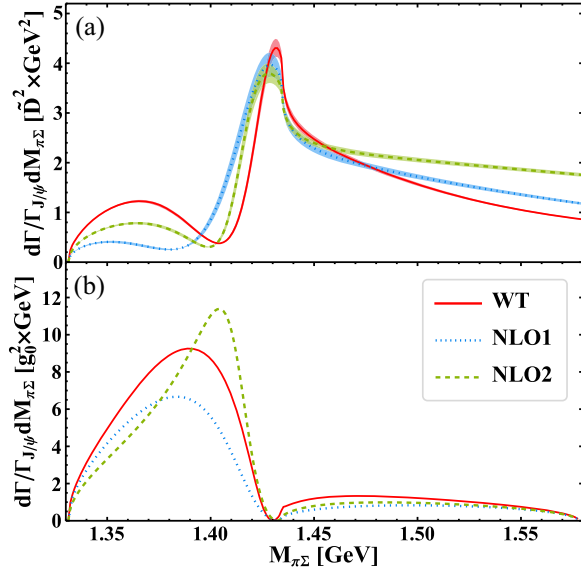


FIG. 2. $\pi\Sigma$ invariant mass distributions in (a) $J/\psi \rightarrow \bar{\Lambda}\Sigma\pi$ and (b) $J/\psi \rightarrow \bar{\Lambda}(1520)\Sigma\pi$. The narrow bands in (a) are obtained by considering the uncertainty in the ratio $\mathcal{R}_{\bar{F}/\bar{D}}$. The curves correspond to different FSI models: WT (solid, LO WT interaction only); NLO1 (dotted, includes LO, NLO, and Born terms) [14]; NLO2 (dashed, includes LO and NLO interactions) [15].

the same mechanism in explaining the $X(3872)$ as a dip in the cross section of $e^+e^- \rightarrow J/\psi\pi^+\pi^-$, we refer to Ref. [43]. It is a consequence of the specific interference pattern dictated by coupled-channel unitarity and the large S -wave $N\bar{K}$ scattering length (and hence a near-threshold pole, corresponding to the narrow octet one).

It should be mentioned that the other octet pole $\Lambda(1680)$ is far from the $\Lambda(1380)$ and $\Lambda(1405)$ and will not pollute the two-pole structure we are interested in here.

The above results were obtained with the LO WT interaction. Next, we further consider the NLO contributions for the NG bosons scattering off the light baryons. There have been various NLO calculations in the literature [13–15, 44–46]. Two sets of parameters from Ref. [14] and Ref. [15] are employed in our calculation. In Ref. [14], the pole positions of the $\Lambda(1380)$ and $\Lambda(1405)$ in the full NLO scheme (NLO1) were found as $(1381 - i81)$ MeV and $(1424 - i26)$ MeV, respectively. In Ref. [15], the pole positions at NLO without the Born terms (NLO2) were found as $(1415 - i165.7)$ MeV and $(1417.9 - i15.6)$ MeV. The NLO chiral Lagrangian can be found explicitly in Ref. [13]. Both the NLO contact and Born terms break the $SU(3)$ symmetry. The numerical NLO results in the two schemes NLO1 and NLO2 are displayed in Fig. 2. One may notice that although the pole positions vary in different schemes, the line-shape behaviors in the NLO schemes remain similar to the LO ones. These numerical results show the reliability of the flavor filter method against NLO corrections. However, the higher-order corrections in the chiral expansion may introduce non-negligible uncertainties in the pole positions

and line shapes, as indicated by the differences between the NLO1 and NLO2 results. Its discriminatory power needs to be further assessed by future theoretical studies including higher-order effects and by improved experimental data.

For the process $J/\psi \rightarrow \bar{\Lambda}\Sigma\pi$, the above discussion only considers the situation where $\pi\Sigma$ is produced from intermediate hyperons with $I = 0$. As a three-body decay process, it can also receive contributions from the cascade decay $J/\psi \rightarrow \Sigma\bar{\Sigma}^* \rightarrow \Sigma\bar{\Lambda}\pi$, where $\bar{\Sigma}^*$ is an intermediate isovector antihyperon, such as the $\bar{\Sigma}(1385)$ state. As can be seen from the Dalitz plots of the recent BESIII measurement [47], the $\bar{\Sigma}(1385)$ signal is sizable. Although the $\bar{\Sigma}^*$ peak appears in the $\pi\bar{\Lambda}$ spectrum, the reflection effect may still contaminate the identification of the excited Λ state in the $\pi\Sigma$ spectrum. The influence of reflection effects can be eliminated by applying an appropriate energy cut in the Dalitz plot. Figure 3 displays the $\pi\Sigma$ distribution under the cut $M_{\bar{\Lambda}\pi} < 1.3$ GeV, which largely excludes all $\bar{\Sigma}^*$ resonances. The main feature of the resulting line shape is similar to that in Fig. 2(a)—a narrow peak rising from about 1.40 GeV and reaching its maximum at about 1.42 GeV. Notably, the BESIII Dalitz plot distribution in the $M_{\bar{\Lambda}\pi} < 1.3$ GeV region shows a data accumulation near $M_{\Sigma^+\pi^-}^2 \approx 2$ GeV² [47]. While BESIII has not yet published a complete amplitude analysis of the $J/\psi \rightarrow \bar{\Lambda}\Sigma\pi$ data, we extract the invariant mass spectrum of the $\pi\Sigma$ system from the Dalitz plot presented in Fig. 2 of Ref. [47] and the result is shown as the data points in Fig. 3. The peak position is located at approximately 1.42 GeV. Although this value is larger than the mass given in the latest version of RPP ($1405.1^{+1.3}_{-1.0}$ MeV), it agrees remarkably well with the real part of the $\Lambda(1405)$ pole positions listed in RPP [3], where the central values range from 1417.7 to

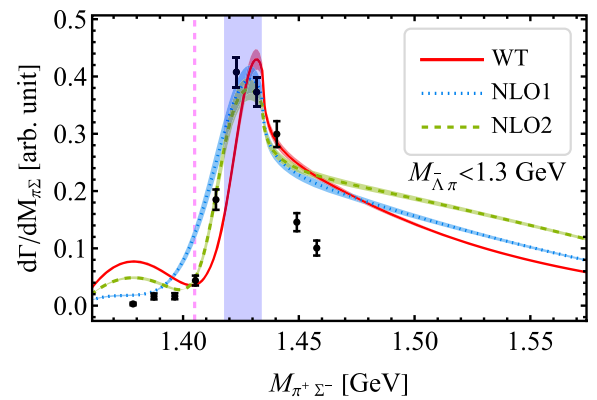


FIG. 3. $\pi\Sigma$ invariant mass distribution in $J/\psi \rightarrow \bar{\Lambda}\Sigma^-\pi^+$ under the cut $M_{\bar{\Lambda}\pi} < 1.3$ GeV. The data are extracted from Fig. 2(a) of Ref. [47] and rescaled to match the distribution curves. The vertical dashed line indicates the $\Lambda(1405)$ mass, while the shaded area represents the range of central values for the real part of the higher $\Lambda(1405)$ pole position, as listed in the current version of RPP [3].

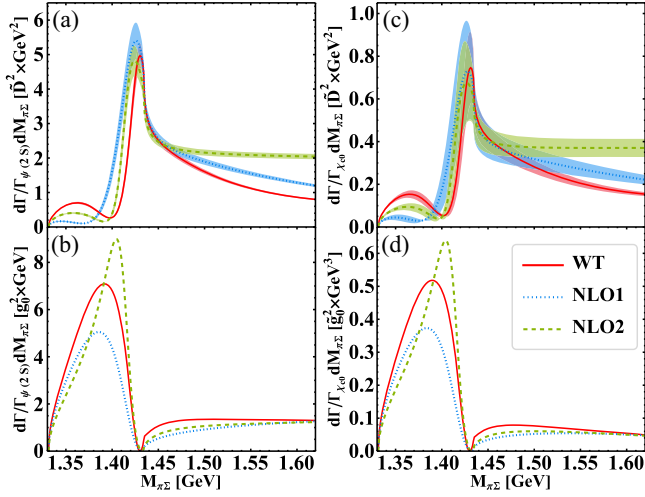


FIG. 4. $\pi\Sigma$ invariant mass distributions in (a) $\psi(2S) \rightarrow \bar{\Lambda}\Sigma\pi$, (b) $\psi(2S) \rightarrow \bar{\Lambda}(1520)\Sigma\pi$, (c) $\chi_{c0} \rightarrow \bar{\Lambda}\Sigma\pi$, and (d) $\chi_{c0} \rightarrow \bar{\Lambda}(1520)\Sigma\pi$. The bands in (a) and (c) are obtained by considering the uncertainty in the ratio $\mathcal{R}_{\bar{F}/\bar{D}}$.

1434 MeV [14,45,46,48]. Furthermore, the peak exhibits a narrow width of only a few tens of MeV. These findings, presented in Fig. 3, are in excellent agreement with theoretical expectation that the higher-mass, narrower octet pole is selectively enhanced in the $J/\psi \rightarrow \bar{\Lambda}\Sigma\pi$ channel.

In addition to the J/ψ decays, the same idea applies to the $\psi(2S)$ and χ_{c0} decays into $\bar{\Lambda}\Sigma\pi$ and $\bar{\Lambda}(1520)\Sigma\pi$.

The effective Lagrangian for the χ_{c0} decaying into $\bar{B}B\phi$ takes the form

$$\mathcal{L}_{\chi_{c0}}^I = \tilde{D}_\chi \langle \bar{B} \gamma_5 \{ \Phi, B \} \rangle \chi_{c0} + \tilde{F}_\chi \langle \bar{B} \gamma_5 [\Phi, B] \rangle \chi_{c0}. \quad (10)$$

For $\chi_{c0} \rightarrow \bar{\Lambda}(1520)B\Phi$, the Lagrangian is given by

$$\mathcal{L}_{\chi_{c0}}^II = i\tilde{g}_0 \bar{\Lambda}_\mu \gamma_5 \langle \partial^\mu \Phi B + \Phi \partial^\mu B \rangle \chi_{c0}. \quad (11)$$

Likewise, the coupling constants $\tilde{D}_\chi$ and $\tilde{F}_\chi$ in Eq. (10) can be determined from the branching fractions of the χ_{c0} decays into the $\bar{\Lambda}\Sigma\pi$, $\bar{\Lambda}N\bar{K}$, and $\bar{\Sigma}N\bar{K}$ channels, and we obtain the ratio $\mathcal{R}_{\bar{F}_\chi/\bar{D}_\chi} = 0.29 \pm 0.17$.

For the $\psi(2S)$ and χ_{c0} decays, the $\pi\Sigma$ invariant mass distribution curves are displayed in Figs. 4(a), (b) and (c), (d), respectively. As before, the octet $\Lambda(1405)$ pole is much more significant in the $\bar{\Lambda}\Sigma\pi$ final states, while the singlet $\Lambda(1380)$ pole dominates the $\bar{\Lambda}(1520)\Sigma\pi$ final state.

Conclusion and outlook. In summary, we proposed a novel method to unambiguously identify the two-pole structure of

$\Lambda(1405)$. The key point of this method is that the two poles are rooted in different irreducible SU(3) flavor representations (one singlet and one octet). For a heavy quarkonium state Y decaying into $\bar{\Lambda}\Sigma\pi$ and $\bar{\Lambda}(1520)\Sigma\pi$, the octet and singlet poles can be singled out separately. This SU(3) flavor filter works even considering the flavor symmetry breaking. Although the positions of the two poles vary depending on the parameters used in the literature, their production weights in the $\bar{\Lambda}\Sigma\pi$ and $\bar{\Lambda}(1520)\Sigma\pi$ channels are rather different, and the line-shape behaviors remain relatively stable using different parameter sets.

This flavor filter is experimentally accessible, since running experiments, e.g., BESIII, have already accumulated huge data samples. With the already available data in Ref. [47], a solid conclusion about the possible resonant structures in the $\pi\Sigma$ distribution can be made once the cut $M_{\bar{\Lambda}\pi}$ is made. It is therefore promising to solve the long-standing $\Lambda(1405)$ puzzle by employing this flavor filter. The present proposal, if confirmed, will provide the first significant experimental validation of the $\Lambda(1405)$ two-pole structure, complementing recent lattice QCD [49,50] and chiral effective field theory studies [15,22–24,51].

The same flavor filter idea can also be used to identify other multiple-pole structures in coupled-channel dynamics. For instance, in the decay of a bottomonium (such as Υ) into $\bar{D}^*D\pi$, the antitriplet pole of $D_0^*(2300)$ may be singled out, since the bottomonium is an SU(3) singlet and $\bar{D}^*$ belongs to a triplet. Note that in this case, the initial state cannot be in the charmonium mass region since the nature of charmonium(-like) resonances that can decay into $\bar{D}^*D\pi$ is under debate and they might not be ideal SU(3) singlets [52].

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Data availability. The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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