

Scattering of $\Lambda_c\Lambda_c$ and $\Lambda_c\bar{\Lambda}_c$ in chiral effective field theory

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We investigate the S -wave scatterings of $\Lambda_c\Lambda_c$ and $\Lambda_c\bar{\Lambda}_c$ systems within a unified chiral effective field theory framework up to next-to-leading order. The contact low-energy coupling constants are determined by fitting to the lattice QCD results for the $\Lambda_c\Lambda_c$ scattering phase shift at an unphysical pion mass. After extrapolating to the physical pion mass, we find a repulsive interaction in the $I(J^P) = 0(0^+)$ $\Lambda_c\Lambda_c$ channel, consistent with the lattice QCD simulation. On the $\Lambda_c\bar{\Lambda}_c$ side, using the fitted contact low-energy constants, we predict the phase shifts and potentials for $\Lambda_c\bar{\Lambda}_c$ scattering in the $I(J^{PC}) = 0(0^{++})$ and $0(1^{--})$ channels. Attractive interactions are found in both channels, each allowing for the formation of bound states. In particular, the attraction in the $0(1^{--})$ $\Lambda_c\bar{\Lambda}_c$ channel is stronger. In addition, our analysis reveals that the spin-spin term caused by the two-pion exchange contributes significantly to the interactions, leading to a distinct mass splitting between the $0(0^{++})$ and $0(1^{--})$ $\Lambda_c\bar{\Lambda}_c$ channels.

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I. INTRODUCTION

In hadron physics, the study of dibaryon (or hexaquark) states has a long and intricate history [1]. Inspired by studies of the nuclear force, the interactions of dibaryon systems have been extensively analyzed using various theoretical methods [2–15]. Meanwhile, in recent years, the HAL QCD method provides an approach for deriving baryon-baryon interaction potentials from first principles [16–22].

Motivated by the aforementioned studies of light dibaryons, it is natural to extend the investigation to dibaryons containing heavy quarks, particularly two-heavy-baryon systems. The combination of chiral symmetry and heavy quark symmetry makes heavy hadron systems ideal platforms for studying hadron interactions. Among these, the interactions of the $\Lambda_c\Lambda_c$ and $\Lambda_c\bar{\Lambda}_c$ systems have been

investigated using various theoretical frameworks. Lee *et al.* systematically studied heavy dibaryon and heavy baryonium systems within the one-boson-exchange model [23]. They concluded that there are no H -dibaryon-like states in the $\Lambda_c\Lambda_c$ sector but that $\Lambda_c\bar{\Lambda}_c$ bound states can exist. In contrast, Refs. [24,25] suggested the possible existence of a double-charm $\Lambda_c\Lambda_c$ bound state. Chen *et al.* investigated the $\Lambda_c\Lambda_c$ interaction in the framework of chiral effective theory (ChEFT) [26]. In addition, various other methods have been applied to study double-charm dibaryon systems [27–40]. For instance, in Ref. [41], Huang, Ping, and Wang studied H -like dibaryon systems with heavy quarks within the quark delocalization color screening model, and their results indicated that the $\Lambda_c\Lambda_c$ interaction is repulsive. On the other hand, several theoretical studies have focused on the $\Lambda_c\bar{\Lambda}_c$ interaction or hidden-charm hexaquark states [42–45]. In Ref. [46], Song and collaborators studied possible molecular states arising from the $\Lambda_c\bar{\Lambda}_c$ interaction mediated by ω and σ exchanges. They found two bound states with quantum numbers $J^{PC} = 0^{-+}$ and 1^{--} and discussed their production in nucleon-antinucleon collisions.

On the experimental side, the $\Lambda_c\bar{\Lambda}_c$ system has attracted considerable interest. The Belle Collaboration first reported the $Y(4630)$ resonance, with mass

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$M = 4634_{-7}^{+8}(\text{stat})_{-8}^{+5}(\text{sys})$ MeV and total width $\Gamma_{\text{tot}} = 92_{-24}^{+40}(\text{stat})_{-21}^{+10}(\text{sys})$ MeV, in the cross section of $e^+e^- \rightarrow \Lambda_c \bar{\Lambda}_c$ [47]. More recently, the BESIII Collaboration measured the cross sections for this process [48,49]. They found no indication of the $Y(4630)$ resonant structure reported by Belle. Additionally, they searched for $H_c^\pm$, a possible $\Lambda_c \bar{\Sigma}_c$ bound state, via the $e^+e^- \rightarrow \pi^+\pi^-\Lambda_c^+\bar{\Lambda}_c^-$ process but observed no statistically significant signal [50].

This leads to a natural question: What is the relation between the $\Lambda_c \bar{\Lambda}_c$ interaction and the observed charmoniumlike states? In Refs. [51,52], final state interactions were introduced to explain the pronounced enhancements around the $\Lambda_c \bar{\Lambda}_c$ threshold in the BESIII data [48,49]. In Ref. [53], the mixing of S -wave and D -wave components of the $\Lambda_c \bar{\Lambda}_c$ wave function due to the tensor part of the final state interactions was considered to explain the near-threshold energy dependence of the $e^+e^- \rightarrow \Lambda_c \bar{\Lambda}_c$ cross section observed by the Belle and BESIII Collaborations [47,48]. Furthermore, the $\Lambda_c \bar{\Lambda}_c$ structure has also been considered in studies of the production and decay of the $Y(4260)$ [54–56]. A systematic investigation of the $\Lambda_c \bar{\Lambda}_c$, $\Sigma_c \bar{\Sigma}_c$, and $\Lambda_b \bar{\Lambda}_b$ interactions was carried out in Ref. [57], where the $Y(4260)$ and $Y(4360)$ were interpreted as possible $\Lambda_c \bar{\Lambda}_c$ baryonia. Several studies based on direct analyses of experimental data have investigated the $\Lambda_c \bar{\Lambda}_c$ molecular state [58–60].

In summary, it is evident that studies of exotic and charmoniumlike states can benefit significantly from investigations of the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ interactions. Moreover, given the intrinsic relation between the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ systems, it is desirable to describe both within a unified framework, ideally with a common set of parameters.

In this work, we therefore present a simultaneous study of the S -wave $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ interactions using ChEFT in the heavy hadron formalism. Studies of low-energy baryon-baryon interactions have achieved considerable success within the ChEFT framework, e.g., for nucleon-nucleon systems [61–66] and hyperon-nucleon systems [67–70]. By incorporating heavy quark symmetry, ChEFT has also demonstrated its effectiveness in describing the low-energy dynamics of two heavy hadrons [26,71–77].

Because of the lack of experimental data or lattice QCD (LQCD) inputs, the unknown low-energy coupling constants (LECs) in ChEFT often have to be fixed using other models, as in Refs. [26,71,73,74,78–82]. In this work, however, we utilize the recently available LQCD data. Recently, Xing *et al.* performed the first LQCD study of $\Lambda_c \Lambda_c$ scattering using Lüscher's finite-volume method, extracting the scattering length and effective range at a pion mass $m_\pi \sim 303$ MeV and lattice spacing $a = 0.07746$ fm [83]. Their results indicate a repulsive interaction for the $I(J^P) = 0(0^+) \Lambda_c \Lambda_c$ system. Coupled-channel effects involving $N\Xi_{cc}$ and $\Sigma_c \Sigma_c$ were also discussed. This LQCD calculation enables us to determine the corresponding LECs in the chiral

Lagrangians by fitting to the lattice data and then to extrapolate the results to the physical point.

This paper is organized as follows: In Sec. II, we introduce the chiral Lagrangians under flavor SU(2) and derive the effective baryon-baryon potentials. In Sec. III, we determine the contact-term LECs by fitting to the S -wave $\Lambda_c \Lambda_c$ phase shifts from LQCD. Using these fitted LECs, we then perform calculations at the physical pion mass, obtaining the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ scattering phase shifts and potentials, and we search for possible bound states by solving the Schrödinger equation. We also investigate the SU(3) effects induced by the two-kaon-exchange contribution and examine their impact on the binding properties of the $\Lambda_c \bar{\Lambda}_c$ systems. Finally, in Sec. IV, we summarize our findings and discuss their implications for future studies of doubly heavy hexaquark systems.

II. CHIRAL EFFECTIVE POTENTIALS

A. Effective Lagrangians

In this work, we adopt the ChEFT Lagrangian in the heavy hadron formalism under the SU(2) case. We use Weinberg's power counting to study the $\Lambda_c \Lambda_c (\bar{\Lambda}_c)$ interactions up to the second chiral order $\mathcal{O}(\epsilon^2)$ at the one-loop level, and the order parameters are $\epsilon = q/\Lambda_\chi$, where q represents the momentum of a Goldstone boson and parameter Λ_χ denotes the chiral symmetry broken scale or the mass of the heavy hadrons.

In heavy hadron formalism, the lowest charmed baryons can be categorized into the following multiplets. The spin- $\frac{1}{2}$ isosinglet is defined as [26]

$$\psi_1 = \begin{bmatrix} 0 & \Lambda_c^+ \\ -\Lambda_c^+ & 0 \end{bmatrix} \quad \text{and} \quad \psi_1^c = \begin{bmatrix} 0 & \bar{\Lambda}_c^- \\ -\bar{\Lambda}_c^- & 0 \end{bmatrix} \quad (1)$$

for the antiparticle. The superscript c represents charge conjugation.

The spin- $\frac{1}{2}$ and spin- $\frac{3}{2}$ isospin triplets are defined as [26,73]

$$\psi_3 = \begin{bmatrix} \Sigma_c^{++} & \frac{\Sigma_c^+}{\sqrt{2}} \\ \frac{\Sigma_c^+}{\sqrt{2}} & \Sigma_c^0 \end{bmatrix}, \quad \psi_3^\mu = \begin{bmatrix} \Sigma_c^{*++} & \frac{\Sigma_c^{*+}}{\sqrt{2}} \\ \frac{\Sigma_c^{*+}}{\sqrt{2}} & \Sigma_c^{*0} \end{bmatrix}^\mu, \quad (2)$$

and

$$\psi_3^c = \begin{bmatrix} \bar{\Sigma}_c^{--} & \frac{\bar{\Sigma}_c^-}{\sqrt{2}} \\ \frac{\bar{\Sigma}_c^-}{\sqrt{2}} & \bar{\Sigma}_c^0 \end{bmatrix}, \quad \psi_3^{c\mu} = \begin{bmatrix} \bar{\Sigma}_c^{*-} & \frac{\bar{\Sigma}_c^{*0}}{\sqrt{2}} \\ \frac{\bar{\Sigma}_c^{*0}}{\sqrt{2}} & \bar{\Sigma}_c^{*+} \end{bmatrix}^\mu \quad (3)$$

for the antiparticle.

The heavy baryon (anti)fields can be projected into the heavy components via the operators $\frac{1 \pm \not{v}}{2}$ [84]:

$$\mathcal{B}_i = e^{iM_i v \cdot x} \frac{1 + \not{v}}{2} \psi_i, \quad \mathcal{B}_i^c = e^{-iM_i v \cdot x} \psi_i^c \frac{1 - \not{v}}{2}, \quad (4)$$

where $\psi_i^{(c)}$ represents the heavy baryon fields $\psi_1^{(c)}$, $\psi_3^{(c)}$, or $\psi_{3^*}^{(c)\mu}$ and M_i is the mass of the heavy baryon. v is the four-velocity. We introduce the leading-order (LO) chiral Lagrangian [26,52,73]

$$\begin{aligned} \mathcal{L}_{\mathcal{B}\phi} = & \frac{1}{2} \text{Tr}[\bar{\mathcal{B}}_1 (i v \cdot D) \mathcal{B}_1] + 2g_2 \text{Tr}[\bar{\mathcal{B}}_3 \mathcal{S} \cdot u \mathcal{B}_1 + \text{H.c.}] \\ & + g_3 \text{Tr}[\bar{\mathcal{B}}_{3^*}^\mu u_\mu \mathcal{B}_3 + \text{H.c.}] + g_4 \text{Tr}[\bar{\mathcal{B}}_{3^*}^\mu u_\mu \mathcal{B}_1 + \text{H.c.}] \\ & + \text{Tr}[\bar{\mathcal{B}}_3 (i v \cdot D - \delta_b) \mathcal{B}_3] + 2g_5 \text{Tr}[\bar{\mathcal{B}}_{3^*}^\mu \mathcal{S} \cdot u \mathcal{B}_{3^*}^\mu] \\ & - \text{Tr}[\bar{\mathcal{B}}_{3^*}^\mu (i v \cdot D - \delta_c) \mathcal{B}_{3^*}^\mu] + 2g_1 \text{Tr}[\bar{\mathcal{B}}_3 \mathcal{S} \cdot u \mathcal{B}_3] \end{aligned} \quad (5)$$

and

$$\begin{aligned} \mathcal{L}_{\mathcal{B}^c\phi} = & \frac{1}{2} \text{Tr}[\bar{\mathcal{B}}_1^c (i v \cdot D) \mathcal{B}_1^c] - 2g_2 \text{Tr}[\bar{\mathcal{B}}_3^c \mathcal{S} \cdot u \mathcal{B}_1^c + \text{H.c.}] \\ & - g_3 \text{Tr}[\bar{\mathcal{B}}_{3^*}^{c\mu} u_\mu \mathcal{B}_3^c + \text{H.c.}] - g_4 \text{Tr}[\bar{\mathcal{B}}_{3^*}^{c\mu} u_\mu \mathcal{B}_1^c + \text{H.c.}] \\ & + \text{Tr}[\bar{\mathcal{B}}_3^c (i v \cdot D - \delta_b) \mathcal{B}_3^c] - 2g_5 \text{Tr}[\bar{\mathcal{B}}_{3^*}^{c\mu} \mathcal{S} \cdot u \mathcal{B}_{3^*}^{c\mu}] \\ & - \text{Tr}[\bar{\mathcal{B}}_{3^*}^{c\mu} (i v \cdot D - \delta_c) \mathcal{B}_{3^*}^{c\mu}] - 2g_1 \text{Tr}[\bar{\mathcal{B}}_3^c \mathcal{S} \cdot u \mathcal{B}_3^c], \end{aligned} \quad (6)$$

for the pion-anticharmed baryon vertices, where $\text{Tr}[\dots]$ denotes the trace in flavor space and we have the spin operator $\mathcal{S}^\mu = (i/2)\gamma_5 \sigma^{\mu\nu} v_\nu$, the mass splittings $\delta_b = M_{\Sigma_c} - M_{\Lambda_c}$ and $\delta_c = M_{\Sigma_c^*} - M_{\Lambda_c}$, and the covariant derivative

$$D_\mu \psi = \partial_\mu \psi + \Gamma_\mu \psi + \psi \Gamma_\mu^T, \quad (7)$$

where Γ_μ^T is the transpose of Γ_μ . The chiral connection Γ_μ and axial current u_μ are defined, respectively, as

$$\Gamma_\mu = \frac{1}{2} [\xi^\dagger, \partial_\mu \xi], \quad u_\mu = \frac{i}{2} \{\xi^\dagger, \partial_\mu \xi\}, \quad (8)$$

with $\xi^2 = U = \exp(i\phi/f_\pi)$, and

$$\phi = \begin{bmatrix} \pi^0 & \sqrt{2}\pi^+ \\ \sqrt{2}\pi^- & -\pi^0 \end{bmatrix}. \quad (9)$$

ψ_3 and $\psi_3^{*\mu}$ are degenerate in the heavy quark limit, which allows us to combine them into the superfields ψ^μ and $\bar{\psi}^\mu$ [26,73]:

$$\psi^\mu = \mathcal{B}_{3^*} - \frac{1}{\sqrt{3}} (\gamma^\mu + v^\mu) \gamma^5 \mathcal{B}_3, \quad (10)$$

$$\bar{\psi}^\mu = \bar{\mathcal{B}}_{3^*} + \frac{1}{\sqrt{3}} \bar{\mathcal{B}}_3 \gamma^5 (\gamma^\mu + v^\mu). \quad (11)$$

In terms of the superfields, the LO contact interaction Lagrangian for the charmed baryons takes a compact form [26]:

$$\begin{aligned} \mathcal{L}_{\mathcal{B}\mathcal{B}} = & C_a \text{Tr}[\bar{\mathcal{B}}_1 \mathcal{B}_1] \text{Tr}[\bar{\mathcal{B}}_1 \mathcal{B}_1] + C_b \text{Tr}[\bar{\mathcal{B}}_1 \mathcal{B}_1] \text{Tr}[\bar{\psi}_\mu \psi^\mu] \\ & + C_c \text{Tr}[\bar{\psi}_\mu \psi^\mu] \text{Tr}[\bar{\psi}_\nu \psi^\nu] + D_a \text{Tr}[\bar{\psi}_\mu \tau^i \psi^\mu] \text{Tr}[\bar{\psi}_\nu \tau^i \psi^\nu] \\ & + i C_d \epsilon_{\mu\nu\alpha\beta} v^\alpha \text{Tr}[\bar{\psi}_\rho \gamma^\beta \gamma_5 \tau^i \psi^\rho] \text{Tr}[\bar{\psi}^\mu \psi^\nu] \\ & + i D_b \epsilon_{\mu\nu\alpha\beta} v^\alpha \text{Tr}[\bar{\psi}_\rho \gamma^\beta \gamma_5 \tau^i \psi^\rho] \text{Tr}[\bar{\psi}^\mu \tau^i \psi^\nu], \end{aligned} \quad (12)$$

where the C_a , C_b , C_c , D_a , C_d , and D_b are six independent LECs and τ^i represents the Pauli matrix in isospin space. For the $\Lambda_c \bar{\Lambda}_c$ systems, the LO contact Lagrangian reads

$$\mathcal{L}_{\mathcal{B}^c\mathcal{B}_1} = C_a \text{Tr}[\bar{\mathcal{B}}_1^c \mathcal{B}_1^c] \text{Tr}[\bar{\mathcal{B}}_1 \mathcal{B}_1]. \quad (13)$$

B. Effective potentials

Under the Breit approximation, the calculated scattering amplitudes $\mathcal{M}(\mathbf{q})$ constructed from the above Lagrangians can be related to the effective potential

$$\mathcal{V}(\mathbf{q}) = - \frac{\mathcal{M}(\mathbf{q})}{\sqrt{2M_1 2M_2 2M_3 2M_4}}. \quad (14)$$

By expanding the Lagrangian Eq. (12), we obtain the LO contact $\Lambda_c \bar{\Lambda}_c$ potential

$$\mathcal{V}_{\text{Cont}}^{\Lambda_c \bar{\Lambda}_c} = -8C_a. \quad (15)$$

Now we turn to the next-to-leading order (NLO). At this order, there are two-pion-exchange (TPE) contributions; the diagrams are shown in Fig. 1. Note that the football and triangle diagrams vanish, since the $\Lambda_c \Lambda_c \pi\pi$ vertex does not exist. Because of the large mass splitting $\delta_c = M_{\Sigma_c^*} - M_{\Lambda_c} = 233.5$ MeV, we do not consider the Σ_c^* contributions explicitly; i.e., their effects are integrated

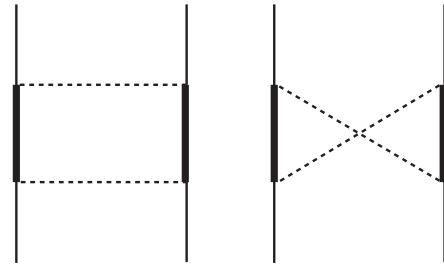


FIG. 1. Two-pion-exchange diagrams for the $\Lambda_c \Lambda_c$ system: the planar box diagram and the crossed box diagram. The thin, thick, and dashed lines denote the Λ_c , Σ_c , and π , respectively.

out and absorbed into the contact LECs. Therefore, the expression for the TPE potential reads

$$\begin{aligned} \mathcal{V}_{2\pi}^{\Lambda_c \Lambda_c} = & \frac{3g_2^4}{4f_\pi^4} \left[-\mathbf{q}^2 J_{21}^B + \frac{2}{3} \mathbf{q}^2 (\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2) J_{21}^B + \mathbf{q}^4 (J_{22}^B + 2J_{32}^B \right. \\ & + J_{43}^B) + 15J_{41}^B - 10\mathbf{q}^2 (J_{31}^B + J_{42}^B) - \mathbf{q}^2 J_{21}^R \\ & - \frac{2}{3} \mathbf{q}^2 (\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2) J_{21}^R + \mathbf{q}^4 (J_{22}^R + 2J_{32}^R + J_{43}^R) \\ & \left. + 15J_{41}^R - 10\mathbf{q}^2 (J_{31}^R + J_{42}^R) \right] (-\delta_b, -\delta_b), \end{aligned} \quad (16)$$

where the definitions and the calculation procedures of the loop functions $J_{ij}^{B/R}$ are given in Refs. [71,74,79,81,82]. The value of the $\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2$ operator is calculated with the relation

$$\langle \boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2 \rangle = 2[S(S+1) - S_1(S_1+1) - S_2(S_2+1)].$$

Expanding the Lagrangian in Eq. (13), we obtain the LO $\Lambda_c \bar{\Lambda}_c$ contact potential

$$\mathcal{V}_{\text{Cont}}^{\Lambda_c \bar{\Lambda}_c} = 4C_a, \quad (17)$$

where the C_a is taken as in Eq. (15). The contribution from the annihilation term is ignored due to the lack of experimental data or LQCD calculations for $\Lambda_c \bar{\Lambda}_c$ scattering. The TPE potential for $\Lambda_c \bar{\Lambda}_c$ system is calculated as

$$\mathcal{V}_{2\pi}^{\Lambda_c \bar{\Lambda}_c} = \mathcal{V}_{2\pi}^{\Lambda_c \Lambda_c}. \quad (18)$$

In numerical calculations, a Gaussian regulator of the form $\mathcal{F}(\mathbf{q}) = \exp(-\mathbf{q}^4/\Lambda^4)$ is introduced to regularize the ultraviolet divergence in the chiral potentials, where Λ is the cutoff parameter [26,71,74,79,85].

For the single-channel elastic scattering process, we calculate the two-body scattering amplitudes T by solving the Lippmann-Schwinger equation

$$\begin{aligned} T(p', p; \sqrt{s}) = & \mathcal{V}(p', p) + \int \frac{p''^2 dp''}{(2\pi)^3} \mathcal{V}(p', p'') \\ & \times \frac{2\mu}{p^2 - p''^2 + i\epsilon} T(p'', p; \sqrt{s}), \end{aligned} \quad (19)$$

with $\sqrt{s}$ the total energy of the two-body system in the center-of-mass frame and μ the reduced mass. The transferred momentum $\mathbf{q}$ is related to the initial momentum $p = |\mathbf{p}|$ and the final momentum $p' = |\mathbf{p}'|$ in the center-of-mass frame via

$$\mathbf{q}^2 = p^2 + p'^2 - 2pp' \cos \theta, \quad (20)$$

where θ is the scattering angle. Based on the effective potentials, we can solve the Lippmann-Schwinger equation

in Eq. (19) and then evaluate the S -wave $\Lambda_c \Lambda_c (\bar{\Lambda}_c)$ scattering phase shift using

$$\tan \delta_0 = \frac{\text{Im}(T_0)}{\text{Re}(T_0)}. \quad (21)$$

To understand the potential behaviors of the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ interactions, we need to make Fourier transformation on $\mathcal{V}(\mathbf{q})$ to get $\mathcal{V}(r)$, the potentials in coordinate space:

$$\mathcal{V}(r) = \int \frac{d^3 \mathbf{q}}{(2\pi)^3} \mathcal{V}(\mathbf{q}) e^{-i\mathbf{q} \cdot \mathbf{r}} \mathcal{F}(\mathbf{q}). \quad (22)$$

Furthermore, based on $\mathcal{V}(r)$, we can solve the Schrödinger equations to investigate whether the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ systems can form bound states.

III. NUMERICAL RESULTS AND DISCUSSION

A. $\Lambda_c \Lambda_c$ interaction

The constant C_a in the contact potential (15) shall be determined by fitting to the LQCD data [83]. For the first time in Ref. [83], a LQCD calculation was performed for the $I(J^P) = 0(0^+) \Lambda_c \Lambda_c$ scattering on two $N_f = 2 + 1$ Wilson-Clover ensembles with the pion mass $m_\pi \sim 303$ MeV and lattice spacing $a = 0.07746$ fm. They obtained the scattering phase shift and determined the scattering length.

In Table I, we list the corresponding pion and charmed baryon masses given in Ref. [83], which will be used in the fitting procedure later. For g_2 in Eq. (16), we adopt $g_2 = -0.6$ determined from the decay widths of the process $\Sigma_c \rightarrow \Lambda_c \pi$ [26,86]. Note that C_a and g_2 are assumed to be pion mass independent [10,87]. The dependence of the pion decay constant f_π on the pion mass has been studied in Ref. [88], and we adopt $f_\pi \sim 93$ MeV for the physical pion mass $m_\pi \sim 138$ MeV and $f_\pi \sim 101$ MeV for $m_\pi \sim 303$ MeV. We have checked that neglecting the pion mass dependence of f_π does not affect our conclusions.

The data points of the LQCD simulations, as well as our fitted results of the $0(0^+) \Lambda_c \Lambda_c$ scattering phase shifts, are presented in Fig. 2, where the F32P30 and F48P30 ensembles are both considered. Note that we exclude the data points at high energy from each ensemble as in Ref. [83]. The parameters C_a with uncertainties are determined by minimizing the $\chi^2/\text{d.o.f.}$ It is observed that the fitted C_a gradually decreases as the cutoff Λ increases.

TABLE I. Baryon masses for different pion masses adopted in the LQCD simulation in Ref. [83].

Ensemble	m_π/MeV	M_{Λ_c}/GeV	M_{Σ_c}/GeV
F32P30	303.2	2.413	2.572
F48P30	303.4	2.410	2.566

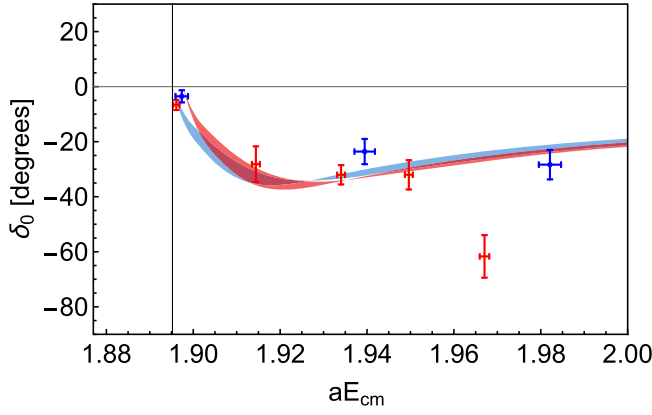


FIG. 2. $0(0^+)\Lambda_c\Lambda_c$ scattering phase shifts as a function of aE_{cm} . The blue and red dots denote the LQCD data from Ref. [83] for the F32P30 and F48P30 ensembles, respectively. The gray-blue and light-red bands correspond to our ChEFT fits of the F32P30 and F48P30 ensembles, respectively, obtained by varying the cutoff Λ from 0.56 to 0.66 GeV.

Here, we give a comment about the finite-volume effects of the lattice data we used. Although the Lüscher quantization condition converts the leading power-law volume dependence of the two-particle energy levels into the infinite-volume phase shifts, residual corrections of the order of $e^{-m_\pi L}$ remain. The two ensembles used in Ref. [83], $m_\pi L = 3.81$ and 5.72 , correspond to exponential suppression factors of about 2.2% and 0.33%, respectively. Therefore, such residual finite-volume effects are expected to be small for the present analysis.

Then, the fitted C_a can be used to estimate the $\Lambda_c\Lambda_c$ scattering phase shifts at the physical pion mass. We therefore calculate the $0(0^+)\Lambda_c\Lambda_c$ scattering phase shifts for $m_\pi = 138$ MeV using two sets of C_a , which are presented in Fig. 3.

The phase shifts obtained at the physical pion mass show only a weak dependence on C_a and cutoff. Our calculation confirms that the $0(0^+)\Lambda_c\Lambda_c$ interaction remains repulsive, consistent with the LQCD studies [83].

To illustrate the characteristics of this interaction, we present in Fig. 4 the $0(0^+)\Lambda_c\Lambda_c$ effective potentials in coordinate space, computed with $\Lambda = 0.60$ GeV at the physical pion mass. The calculation shows that the attractive TPE contribution is considerably weaker than the repulsive contact term, resulting in an overall repulsive potential. It contrasts the results obtained with those in Refs. [23,25] but is consistent with the results in Refs. [26,41].

Finally, we discuss the $\chi^2/\text{d.o.f.}$ obtained in Table II. It should be noted that only a limited number of phase shift data points are available in the fitting region, resulting in a small number of degrees of freedom. Consequently, the obtained $\chi^2/\text{d.o.f.}$ values are relatively large. The limited number of the data points can indeed affect the precision of the fitted parameters and the obtained interaction potentials.

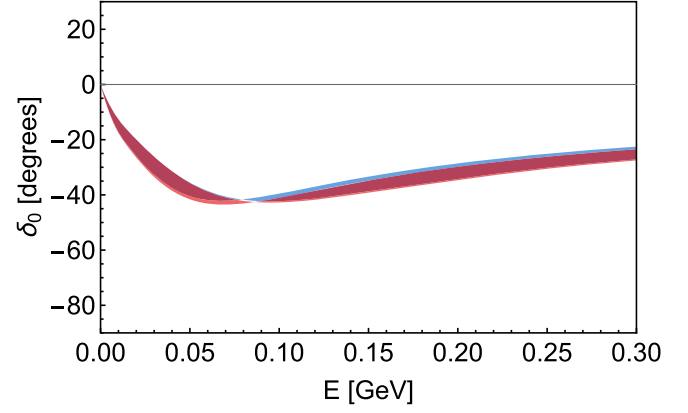


FIG. 3. $0(0^+)\Lambda_c\Lambda_c$ scattering phase shifts as a function of the kinetic energy in the center-of-mass frame at the physical pion mass. The gray-blue and light-red bands correspond to our ChEFT fits of the F32P30 and F48P30 ensembles, respectively, obtained by varying the cutoff Λ from 0.56 to 0.66 GeV.

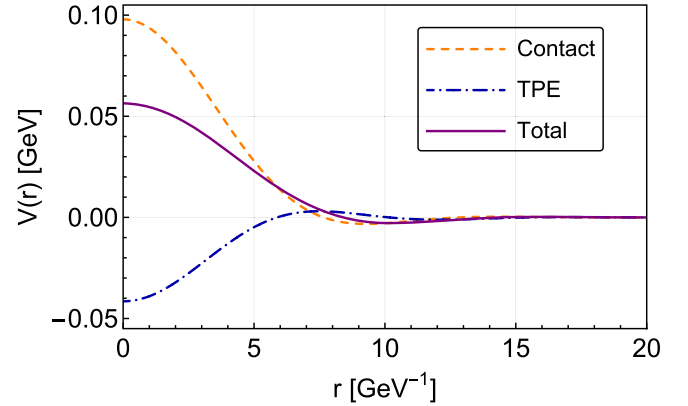


FIG. 4. $0(0^+)\Lambda_c\Lambda_c$ effective potentials in coordinate space at the physical pion mass, with cutoff $\Lambda = 0.60$ GeV. Here, we set $C_a = -3.65 \text{ GeV}^{-2}$ from Table II.

However, we believe the general behaviors of our results are not affected.

B. $\Lambda_c \bar{\Lambda}_c$ interaction

Employing the effective potentials given in Eqs. (17) and (18), we calculate the $0(0^{++})$ and $0(1^{--})\Lambda_c\bar{\Lambda}_c$ scattering phase shifts as functions of the center-of-mass kinetic energy at the physical pion mass, as well as the effective potentials in coordinate space, using a fixed cutoff $\Lambda = 0.60$ GeV. The results are presented in Figs. 5 and 6, respectively.

As illustrated in Fig. 5(a), the phase shift of the $0(0^{++})$ channel begins at 180° and decreases as the center-of-mass kinetic energy increases, indicating an attractive interaction. According to Levinson's theorem [89], this behavior suggests the existence of a bound state. Also, we find that the two curves nearly coincide due to the similar C_a values

TABLE II. Fitted values of constant C_a with uncertainty (in units of GeV^{-2}) in the $\Lambda_c \Lambda_c$ contact potential and the corresponding $\chi^2/\text{d.o.f.}$, obtained using the F32P30 and F48P30 ensembles for various cutoff values Λ (in units of GeV).

Λ	0.56	0.58	0.60	0.62	0.64	0.66
C_a (F32P30)	-4.04 ± 0.19	-3.75 ± 0.18	-3.71 ± 0.17	-3.60 ± 0.16	-3.46 ± 0.16	-3.38 ± 0.15
$\chi^2/\text{d.o.f.}$	6.40	6.33	6.15	6.04	5.94	5.84
C_a (F48P30)	-3.90 ± 0.19	-3.73 ± 0.18	-3.65 ± 0.17	-3.55 ± 0.16	-3.44 ± 0.16	-3.36 ± 0.15
$\chi^2/\text{d.o.f.}$	6.39	5.92	5.47	5.11	4.81	4.56

in the two sets. In the $0(1^{--})$ channel [Fig. 5(b)], the phase shift also decreases with increasing center-of-mass kinetic energy, so this channel possesses an attractive interaction. The consistency of the two channels originates from their common contact potentials [Eq. (17)].

Then, we present the effective potentials in coordinate space obtained with the fitted C_a at $\Lambda = 0.60 \text{ GeV}$ in Fig. 6. For both the $0(0^{++})$ and $0(1^{--})$ channels in Fig. 6, the LO contact interactions and NLO TPE interactions are all attractive. Compared to the $0(0^{++})$, the attraction of

the $0(1^{--})$ TPE potential is stronger, suggesting that the $0(1^{--})$ channel may support a more deeply bound state. Now we proceed to solve the Schrödinger equation to locate the $0(0^{++})$ and $0(1^{--})$ $\Lambda_c \bar{\Lambda}_c$ bound states. After the calculations, the dependences of the binding energies on the cutoff Λ are listed in Table III. At $\Lambda = 0.56 \text{ GeV}$, a very shallow bound state starts to appear in the $0(0^{++})$ channel, while a deeper bound state is found in the $0(1^{--})$ channel. As Λ increases, the binding energies of both states increase accordingly.

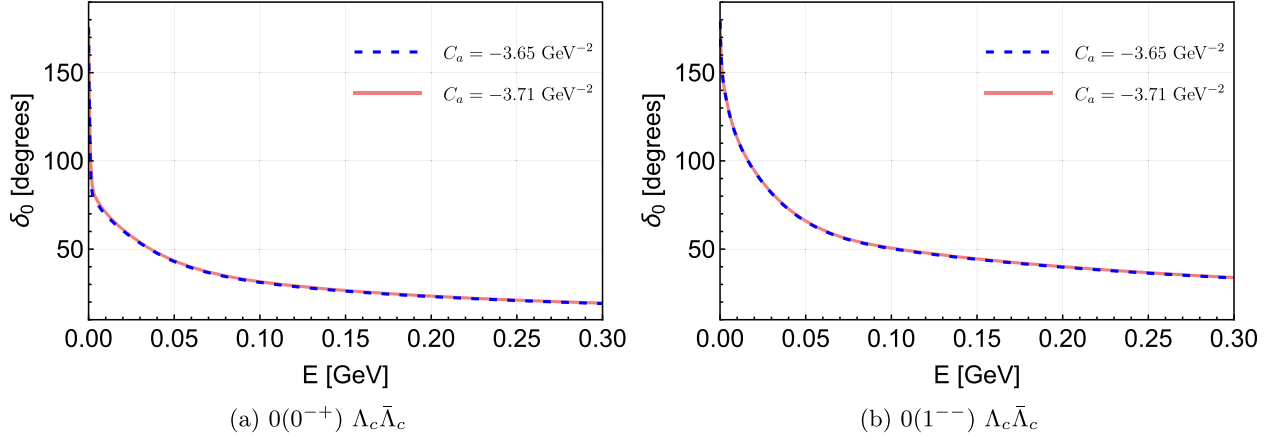


FIG. 5. $0(0^{++})$ and $0(1^{--})$ $\Lambda_c \bar{\Lambda}_c$ scattering phase shifts as a function of the kinetic energy E in the center-of-mass frame at the physical pion mass, with cutoff $\Lambda = 0.60 \text{ GeV}$.

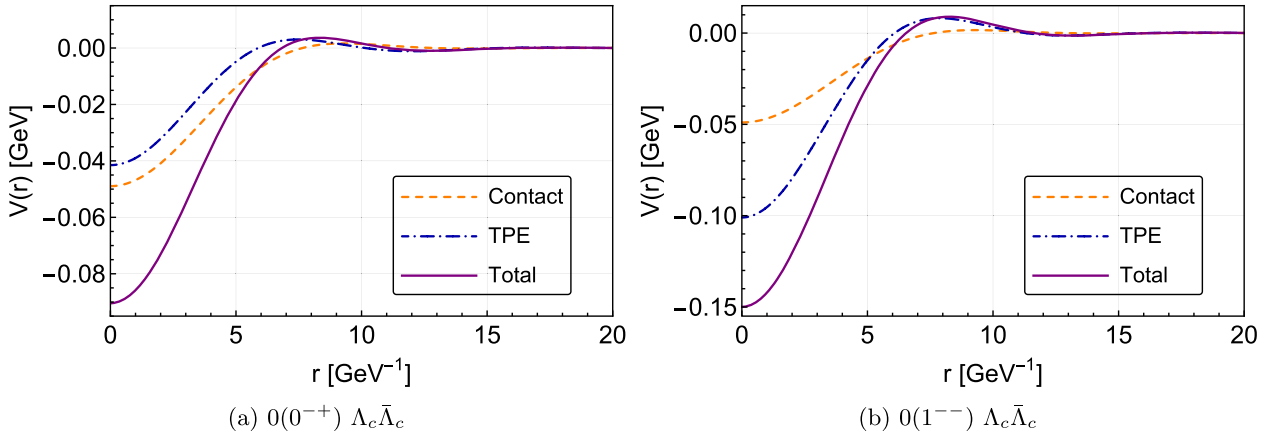


FIG. 6. The $0(0^{++})$ and $0(1^{--})$ $\Lambda_c \bar{\Lambda}_c$ effective potentials in coordinate space at the physical pion mass using the cutoff $\Lambda = 0.60 \text{ GeV}$ and the fitted $C_a = -3.65 \text{ GeV}^{-2}$ in Table II.

TABLE III. Dependence of the binding energies (in units of MeV) on the cutoff Λ (in units of MeV) for the $0(0^{++})$ and $0(1^{--}) \Lambda_c \bar{\Lambda}_c$ systems. The results are obtained using the fitted C_a values from F32P30 and F48P30 ensembles, respectively (see Table II).

Λ	$0(0^{++})$		$0(1^{--})$	
	F32P30	F48P30	F32P30	F48P30
560	$0.2^{+0.2}_{-0.1}$	$0.1^{+0.2}_{-0.1}$	$8.6^{+0.8}_{-0.7}$	$8.1^{+0.7}_{-0.8}$
580	$0.3^{+0.2}_{-0.2}$	$0.2^{+0.2}_{-0.1}$	$10.7^{+0.8}_{-0.9}$	$10.5^{+0.9}_{-0.8}$
600	$0.5^{+0.4}_{-0.1}$	$0.5^{+0.3}_{-0.2}$	$14.5^{+0.9}_{-1.0}$	$14.1^{+1.0}_{-0.9}$
620	$1.0^{+0.4}_{-0.3}$	$0.9^{+0.4}_{-0.3}$	$18.6^{+1.1}_{-1.0}$	$18.3^{+1.1}_{-1.0}$
640	$1.5^{+0.5}_{-0.4}$	$1.5^{+0.4}_{-0.5}$	$23.5^{+1.2}_{-1.3}$	$23.3^{+1.2}_{-1.2}$
660	$2.5^{+0.5}_{-0.6}$	$2.3^{+0.6}_{-0.5}$	$29.7^{+1.3}_{-1.3}$	$29.5^{+1.3}_{-1.3}$

It is worth mentioning the mass splitting between the $0(0^{++})$ and $0(1^{--})$ channels. In Refs. [23,46], the $0(0^{++})$ and $0(1^{--}) \Lambda_c \bar{\Lambda}_c$ systems produced the bound states with almost the same binding energy. It is caused by the fact that their spin-dependent terms were negligible in the S -wave $\Lambda_c \bar{\Lambda}_c$ interaction. In our work, however, the spin-spin terms originated from the TPE potentials play a decisive role, leading to a distinct mass splitting between the $0(0^{++})$ and $0(1^{--})$ channels.

So, the splitting caused by the TPE can be treated as a key feature in our ChEFT framework. This obvious discrepancy between ChEFT and the model calculations [23,46] can serve as a benchmark in future lattice and experimental verifications.

Here, we can give some comments about the $\Lambda_c \bar{\Lambda}_c$ annihilation contribution. The annihilation terms are usually treated as short-distance dynamics. For example, the annihilation contribution was parametrized, and the corresponding parameters are determined from antibaryon-baryon scattering phase shifts [90–92]. In our work, due to the lack of direct constraints from $\Lambda_c \bar{\Lambda}_c$ scattering or annihilation observables, the parametrized annihilation terms cannot be determined in the present work. However, Ref. [52] studied the experimental $e^+e^- \rightarrow \Lambda_c \bar{\Lambda}_c$ cross section data, parametrized, and constrained the annihilation contributions. The results showed that the annihilation terms are not the main contribution near threshold.

Regarding experimental candidates, the $Y(4630)$ reported by the Belle Collaboration is unlikely to be a $\Lambda_c \bar{\Lambda}_c$ molecule, as its mass peak lies significantly above the $\Lambda_c \bar{\Lambda}_c$ threshold [46,47]. In Ref. [58], a $\Lambda_c \bar{\Lambda}_c$ molecule was identified, with its pole position determined by fitting to the BESIII data [48]. Subsequently, an analysis of the $e^+e^- \rightarrow \Lambda_c \bar{\Lambda}_c$ data [49] predicted a $\Lambda_c \bar{\Lambda}_c$ bound state ~ 38 MeV below the threshold using a single-channel approach [59], while a coupled-channel study in Ref. [60] found a similar pole but located above the threshold. Our results are consistent with these studies that directly analyze the

experimental data. Based on our calculations, we suggest searching for the $0(0^{++}) \Lambda_c \bar{\Lambda}_c$ molecular state in the $D\bar{D}^*$ and $D^*\bar{D}^*$ final states. Meanwhile, the $0(1^{--})$ state should be sought in the $D\bar{D}$, $D\bar{D}^*$, and $D^*\bar{D}^*$ channels. Such heavy baryonium states may be accessible in, e.g., the B -meson decays, e^+e^- and $p\bar{p}$ collisions at the LHCb, BESIII, Belle II, and the future PANDA experiment.

C. The estimation of the SU(3) effects

To further investigate the flavor effects beyond the SU(2) case, we extend the interaction to the SU(3) symmetry framework by including the strange pseudoscalar mesons. In the SU(3) case, the spin- $\frac{1}{2}$ and spin- $\frac{3}{2}$ charmed baryon fields can be organized as antisymmetric and symmetric 3×3 matrices, respectively. Meanwhile, the η and K mesons are included in the pseudoscalar meson field ϕ . Since the concerned vertices are negligible, we do not consider the η exchange here [52]. Therefore, the dominant SU(3) correction comes from the two-kaon-exchange (TKE) term at NLO, as illustrated by the football diagram in Fig. 7. The corresponding potential is given by

$$\mathcal{V}_{2K}^{\Lambda_c \Lambda_c} = \frac{1}{2f_\pi^4} J_{22}^F. \quad (23)$$

Including the TKE term, we can obtain the total potential

$$\mathcal{V}_{\text{Total}}^{\Lambda_c \Lambda_c} = \mathcal{V}_{\text{Cont}}^{\Lambda_c \Lambda_c} + \mathcal{V}_{2\pi}^{\Lambda_c \Lambda_c} + \mathcal{V}_{2K}^{\Lambda_c \Lambda_c}. \quad (24)$$

Then, we can refit the LQCD results and present the resulting C_a with uncertainties in Table IV. $m_K \sim 524$ MeV is adopted for $m_\pi \sim 303$ MeV [83]. Compared with the fitted values in the SU(2) case, the values of C_a shift after including the SU(3) effects.

For the $\Lambda_c \bar{\Lambda}_c$ interaction, the TKE contribution has an opposite sign compared with that in the $\Lambda_c \Lambda_c$ channel:

$$\mathcal{V}_{2K}^{\Lambda_c \bar{\Lambda}_c} = -\mathcal{V}_{2K}^{\Lambda_c \Lambda_c}. \quad (25)$$

Using the C_a in Table IV, we solve the Schrödinger equation to search for the $0(0^{++})$ and $0(1^{--}) \Lambda_c \bar{\Lambda}_c$ bound

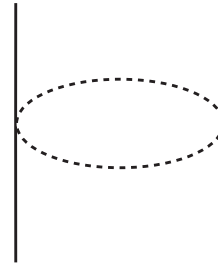


FIG. 7. Two-kaon-exchange football diagram for the $\Lambda_c \Lambda_c$ system. The thin and dashed lines denote the Λ_c and K , respectively.

TABLE IV. Refitted values of C_a with uncertainty (in units of GeV^{-2}) and the corresponding $\chi^2/\text{d.o.f.}$ for various cutoff values Λ (in units of GeV) with $\text{SU}(3)$ effects.

Λ	0.56	0.58	0.60	0.62	0.64	0.66
C_a (F32P30)	-4.98 ± 0.19	-4.83 ± 0.18	-4.69 ± 0.17	-4.57 ± 0.16	-4.46 ± 0.16	-4.37 ± 0.15
$\chi^2/\text{d.o.f.}$	6.41	6.28	6.16	6.06	5.95	5.86
C_a (F48P30)	-4.85 ± 0.19	-4.72 ± 0.18	-4.61 ± 0.17	-4.51 ± 0.16	-4.42 ± 0.16	-4.33 ± 0.15
$\chi^2/\text{d.o.f.}$	6.50	6.00	5.57	5.20	4.90	4.64

TABLE V. Dependence of the binding energies (in units of MeV) on the cutoff Λ (in units of MeV) for the $0(0^{++})$ and $0(1^{--}) \Lambda_c \bar{\Lambda}_c$ systems. The results are obtained using the refitted C_a in Table IV. “...” means that there is no bound state.

Λ	$0(0^{++})$	$0(1^{--})$	
		F32P30	F48P30
560	...	$3.4^{+0.6}_{-0.6}$	$3.0^{+0.6}_{-0.6}$
580	...	$4.9^{+0.7}_{-0.6}$	$4.5^{+0.6}_{-0.7}$
600	...	$7.0^{+0.8}_{-0.8}$	$6.6^{+0.8}_{-0.7}$
620	...	$9.7^{+1.0}_{-0.8}$	$9.4^{+0.9}_{-0.9}$
640	...	$13.2^{+1.0}_{-1.1}$	$13.0^{+1.0}_{-1.1}$
660	...	$17.7^{+1.2}_{-1.2}$	$17.4^{+1.1}_{-1.2}$

states at physical pion and kaon masses. The cutoff dependences of the binding energies are shown in Table V.

Compared with the $\text{SU}(2)$ case, the interactions here in the two channels are still attractive, and the discrepancy among them emphasized in Sec. III B still exists. Note that the inclusion of KK exchange weakens the attraction in both channels, causing the shallow bound state in the $0(0^{++})$ channel to disappear and reducing the binding energy in the $0(1^{--})$ channel. In a word, in the $\text{SU}(3)$ case, although both channels are still attractive, only the $0(1^{--})$ bound state remains.

IV. SUMMARY

Employing a ChEFT framework with a common set of LECs constrained by LQCD simulations, we simultaneously investigate the S -wave interactions of the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ systems. Specifically, we calculate the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ potentials up to NLO, incorporating both contact and TPE contributions. The common LECs are determined by fitting to the LQCD results for the $I(J^P) = 0(0^+) \Lambda_c \Lambda_c$ phase shifts at a pion mass of approximately 303 MeV . Using these LECs, we then explore the interaction properties at the physical pion mass, including phase shifts, coordinate space potentials, and the possible formation of bound states. The robustness of this study lies in its direct connection to LQCD input.

In the $\Lambda_c \Lambda_c$ sector, our results consistently support a repulsive nature for the $0(0^+)$ channel. A detailed analysis of the effective potential in coordinate space reveals that the repulsive contact term dominates over the attractive two-pion-exchange contribution, yielding an overall repulsive potential.

For the $\Lambda_c \bar{\Lambda}_c$ system, the phase shifts clearly indicate attractive interactions. Further analysis of the coordinate space effective potential and a search for bound state solutions show that the attraction is more pronounced in the $0(1^{--})$ channel, where a bound state is more likely to exist. The spin-spin interaction term, which is non-negligible in our effective potential, emerges as the key factor distinguishing the dynamics of the $0(0^+)$ and $0(1^{--}) \Lambda_c \bar{\Lambda}_c$ systems—a refinement not addressed in earlier works [23,46].

In summary, our results reveal that the $\Lambda_c \Lambda_c$ channel with quantum numbers $0(0^+)$ exhibits a repulsive interaction. In contrast, both the $0(0^+)$ and $0(1^{--})$ channels of the $\Lambda_c \bar{\Lambda}_c$ system are found to be attractive, indicating the possible existence of bound states, particularly in the $0(1^{--})$ channel. In addition, there is a distinct mass splitting between the two channels that is attributed to the spin-spin terms in the TPE. Furthermore, we consider the $\text{SU}(3)$ effects induced by the TKE contribution and refit the C_a accordingly. Although both channels are still attractive in the $\text{SU}(3)$ case, only the $0(1^{--})$ bound state remains.

This work provides a consistent and comparative description of the $\Lambda_c \Lambda_c$ and $\Lambda_c \bar{\Lambda}_c$ interactions within a unified framework constrained by LQCD data. The predicted attractive potentials in the $\Lambda_c \bar{\Lambda}_c$ system, along with the possibility of bound states, as well as the predicted splitting brought by the TPE, motivate future experimental searches and more refined LQCD simulations aimed at exploring these exotic states.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [83,93].

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