We study the $DN$ and $D^*N$ interactions to probe the inner structure of $\Sigma_c(2800)$ and $\Lambda_c(2940)$ with the chiral effective field theory to the next-to-leading order. We consider the contact term, one-pion-exchange and two-pion-exchange contributions to characterize the short-, long- and mid-range interactions of the $D^{(*)}N$ systems. The low energy constants of the $D^{(*)}N$ systems are related to those of the $NN$ interaction with quark level Lagrangian that inspired by the resonance saturation model. The $\Delta(1232)$ degree of freedom is also included in the loop diagrams. The attractive potential in the $[DN]_{J^P=1/2}^{J^P=1/2}$ channel is too weak to form bound state, which indicates the explanation of $\Sigma_c(2800)$ as the compact charmed baryon is more reasonable. Meanwhile, the potentials of the isoscalar channels are deep enough to yield the molecular states. We obtain the masses of the $[DN]_{J^P=0}^{J^P=0}$, $[D^*N]_{J^P=0}^{J^P=0}$ and $[D^*N]_{J^P=0}^{J^P=0}$ systems to be 2792.0 MeV, 2943.6 MeV and 2938.4 MeV, respectively. The $\Lambda_c(2940)$ is probably the isoscalar $D^*N$ molecule considering its low mass puzzle. Besides, the $\Lambda_c(2940)$ signal might contain the spin-$\frac{1}{2}$ and spin-$\frac{3}{2}$ two structures, which can qualitatively explain the significant decay ratio to $D^0p$ and $\Sigma_c\pi$. We also study the $B^{(*)}N$ systems and predict the possible molecular states in the isoscalar channels. We hope experimentalists could hunt for the open charmed molecular pentaquarks in the $\Lambda_c^0\pi^+\pi^-$ final state.

I. INTRODUCTION

Hadron spectroscopy plays an important role in understanding the low energy behaviors of QCD. Quark model is very successful in describing the hadron spectra [1]. But it is rather difficult to assign the near-threshold states, such as $X(3872)$ [2] and $D_{s0}(2317)$ [3] to the quark model predictions [4–9]. In the charmed baryon family, a state $\Lambda_c(2940)$ also falls into the same situation as the $X(3872)$ and $D_{s0}(2317)$.

In 2007, the BaBar Collaboration observed a charmed baryon $\Lambda_c(2940)$ in the $D^0p$ invariant mass spectrum [10], which is an isosinglet since no signal is observed in the $D^+p$ final state. It was subsequently confirmed by the Belle experiment in the decay mode $\Lambda_c(2940) \rightarrow \Sigma_c\pi$ [11]. In 2017, the $J^P$ quantum numbers of $\Lambda_c(2940)$ was constrained by the LHCb measurement, and the most likely spin-parity assignment for $\Lambda_c(2940)$ is $J^P = \frac{3}{2}^-$ [12] (The mass and width of $\Lambda_c(2940)$ obtained by the BaBar, Belle and LHCb experiments are shown in Table I).

| TABLE I. The mass and width of $\Lambda_c(2940)$ in experiments (in units of MeV). |
| BaBar   | 2939.8 $\pm$ 1.3 $\pm$ 1.0 | 17.5 $\pm$ 5.2 $\pm$ 5.9 |
| Belle   | 2938.0 $\pm$ 1.3 $^{+2.0}_{-0.4}$ | 13. $^{+1.2}_{-1.3}$ $^{+2.7}_{-0.5}$ |
| LHCb   | 2944.8$^{+3.5}_{-2.5}$ $\pm$ 0.4$^{+0.3}_{-0.4}$ $^{+0.6}_{-0.5}$ | 27.7$^{+0.8}_{-0.6}$ $^{+0.9}_{-0.7}$$^{+5.2}_{-10.4}$ |

Up to now, there are two different interpretations of the internal structure of $\Lambda_c(2940)$. One is the ordinal charmed baryon, and the other one is the $D^*N$ molecular state. However, it is difficult to arrange $\Lambda_c(2940)$ to the $2P$ state in the charmed baryon spectroscopy, since its mass is about 60–100 MeV smaller than the calculations of the quark models [13–16]. Considering the $\Lambda_c(2940)$ lies about 6 MeV below the $D^*p$ threshold, the molecular explanation was firstly proposed in Ref. [17], where the $\Lambda_c(2940)$ as the $\frac{3}{2}^-$ molecular state is preferred by analyzing its decay behaviors. In Ref. [18], He et al. studied the $D^*N$ interaction with the one-boson-exchange model, and their calculation supports the interpretation of the $\Lambda_c(2940)$ as the $D^*N$ bound state with $I(J^P) = 0(\frac{1}{2}^+)$ or $0(\frac{3}{2}^+)$. In Ref. [19], Ortega et al. investigated the $\Lambda_c(2940)$ as a $D^*N$ molecule in the constituent quark model, and they obtain the binding solution in isoscalar $J^P = \frac{3}{2}^-$ channel. In Refs. [20, 21], the strong and radiative decays of $\Lambda_c(2940)$ are calculated in the molecular picture. A QCD sum rule study in Ref. [22] indicates the $\Lambda_c(2940)$ is not a compact state. Some recent calculations based on the chiral quark model also support the molecular explanation for $\Lambda_c(2940)$ [23, 24], (see Refs. [25–29] for review and Refs. [30–39] for other related works).

Another charmed baryon related with the $DN$ threshold is $\Sigma_c(2800)$, which is an isospin triplet and firstly observed by the Belle Collaboration in the $\Lambda_c\pi$ mass spectrum [40]. The neutral state $\Sigma_c(2800)^0$ was possibly confirmed by the BaBar experiment [41], but the measured mass from BaBar is about 50 MeV larger. The $J^P$ of $\Sigma_c(2800)$ is still undetermined yet [1]. Like the $\Lambda_c(2940)$, $\Sigma_c(2800)$ is interpreted as the $P$-wave excitation of the charmed baryon in the $\lambda$ mode [14, 36, 42–44], and $DN$ molecule [22–24, 45], respectively.

Investigating the $DN$ and $D^*N$ interaction is essential to disentangle the puzzles of $\Sigma_c(2800)$ and $\Lambda_c(2940)$. Besides, understanding the $D^{(*)}N$ interaction is also crucial to probe the $D$-mesic nuclei [46, 47] and the properties of the charmed mesons in the nuclear matter [48, 49]. An alternative approach based on the meson-exchange model [50] has been employed to construct the $DN$ and $D^*N$ interaction by the Jülich group [51–53].

Instead of the boson-exchange model, the modern theory
of nuclear force is built upon the pioneer work of Weinberg [54, 55] and largely developed in the framework of effective field theory. The chiral effective field theory was extensively exploited to study the $NN$ interaction with great success [56–61]. The chiral effective field theory was also utilized to study the systems with heavy flavors in Refs. [62–68], which is a powerful tool in predicting the $BB^*$ and $B^*B^*$ bound states [64], reproducing the newly observed pentaquarks [66], extrapolating the $\Sigma_cN$ potential from lattice QCD to the physical pion mass [67], and so on. As a natural extension of the $NN$ interaction, in this work, we use the chiral effective field theory to study the $D(1^{*})N$ interaction up to the next-to-leading order. We simultaneously consider the long-, mid- and short-range interactions, and include the contribution of $\Delta(1232)$ in the loops as an intermediate state. With the chiral effective field theory, we calculate the $D(1^{*})N$ effective potentials and search for the possible bound states. The numerical results can be compared with the experimental data of $\Lambda_c(2940)$ and $\Sigma_c(2800)$ to see whether they are the genuine charmed baryons or the molecular nature.

This paper is organized as follows. In Sec. II, we give the Lagrangians and effective potentials of the $D(1^{*})N$ systems. In Sec. III, we illustrate our numerical results and discussions. In Sec. IV, we conclude with a short summary. In Appendix A, we relate the low energy constants to those of the $NN$ system with quark model.

II. LAGRANGIANS AND EFFECTIVE POTENTIALS

A. Effective chiral Lagrangians

We first show the leading order Lagrangian of the nucleon and pion interaction under the heavy baryon reduction [69], which reads

$$\mathcal{L}_{N\varphi} = \mathcal{N}(iv \cdot \mathcal{D} + 2g_\alpha S \cdot \alpha\varphi)\mathcal{N},$$

(1)

where $\mathcal{N} = (p, n)^T$ denotes the the large component of the nucleon field under the nonrelativistic reduction. $v = (1, 0)$ is the 4-velocity of the nucleon and $\mathcal{D}_\mu = \partial_\mu + \Gamma_\mu$. $g_\alpha \approx 1.29$ is the axial-vector coupling constant. $S^\mu = \frac{i}{2}\gamma_5\sigma^{\mu\nu}v_\nu$ stands for the spin operator of the nucleon. $\Gamma_\mu$ and $u_\mu$ are the chiral connection and axial-vector current, respectively. Their expressions read

$$\Gamma_\mu = \frac{1}{2}[\xi^\dagger, \partial_\mu \xi] \equiv \tau^i\Gamma^i, \quad u_\mu = \frac{i}{2}\{\xi^\dagger, \partial_\mu \xi\} \equiv \tau^i\omega^i, \quad \tau^+ = \frac{1}{2}[\xi^\dagger, \partial_\mu \xi] \equiv \tau^i\omega^i, \quad \tau^+ = \frac{1}{2}[\xi^\dagger, \partial_\mu \xi] \equiv \tau^i\omega^i,$$

(2)

where $\tau^i$ is the Pauli matrix,

$$\xi^2 = U = \exp \left( \frac{i\varphi}{f_\pi} \right), \quad \varphi = \left[ \begin{array}{c} \pi^0 \\ \sqrt{2}\pi^- \\ -\pi^0 \end{array} \right],$$

(3)

and $f_\pi = 92.4$ MeV is the pion decay constant.

Considering the importance of $\Delta(1232)$ in the $NN$ interaction [50, 70–72], we adopt the small scale expansion method [73] to explicitly include the $\Delta(1232)$ in the Lagrangians. The Lagrangian that delineates the $\Delta-N$ coupling is given as

$$\mathcal{L}_{\Delta N\varphi} = -T^i \left( i\gamma^5 \partial^i \delta_\alpha + 2g_1 S \cdot u^i \right) g_{\mu \nu} T^\nu, \quad \mathcal{L}_{\Delta N\varphi} = 2g_2 \left( T^i \omega^i_\mu N \right. \left. + \bar{N} \omega^i_\mu g_{\mu \nu} T^\nu \right),$$

(4)

where $\delta_\alpha = m_\Delta - m_N$, $g_1 = \frac{g_0}{2}$, $g_0$ is estimated with the quark model [73]. $g_1 \approx 1.05$ is the coupling constant for $\Delta N\pi$ vertex. $T^i$ denotes the spin-$\frac{3}{2}$ and isospin-$\frac{1}{2}$ field $\Delta(1232)$ after performing the nonrelativistic reduction. Its matrix form reads

$$T^1 = \frac{1}{\sqrt{2}} \left( \frac{\Delta^{++}}{\sqrt{3}} - \frac{1}{3} \Delta^{-} \right), \quad T^2 = \frac{i}{\sqrt{2}} \left( \frac{\Delta^{++} + \sqrt{3} \Delta^{0}}{\sqrt{3}} \right), \quad T^3 = \sqrt{\frac{2}{3}} \left[ \frac{\Delta^{-}}{\Delta^{0}} \right]$$

(6)

The leading order Lagrangian that depicts the interaction between the charmed mesons and light Goldstones reads [74, 75]

$$\mathcal{L}_{\varphi} = i(\mathcal{H} v \cdot \mathcal{D})\mathcal{H} - \frac{1}{8} \delta_0 \mathcal{H} \sigma^{\mu \nu} \mathcal{R} \sigma_{\mu \nu} + g(\mathcal{H}^5 \gamma_5 \mathcal{H}),$$

(7)

where $\mathcal{H}^5 \gamma_5 \mathcal{H}$ represents the trace in spinor space. $\delta_0$ is defined as $\delta_0 = \langle m_D - m_D \rangle$, $g \approx -0.59$ stands for the axial coupling, whose sign is determined with the help of quark model. The $\mathcal{H}$ is the super-field for the charmed mesons, which reads

$$\mathcal{H} = \frac{1}{2} \left( P^{\ast} \gamma^5 + i P \gamma_5 \right),$$

$$\mathcal{H} = \gamma_5 \gamma_0 \gamma_0 = \left( P^{\ast} \gamma^5 + i P \gamma_5 \right) \frac{1 + \gamma_5}{2},$$

(8)

with $P = (D^0, D^{-})^T$ and $P^{\ast} = (D^{0}, D^{-})^T$, respectively. We construct the leading order contact Lagrangian to describe the short distance interaction between the nucleon and charmed meson

$$\mathcal{L}_{\mathcal{N}\mathcal{H}} = D_{\alpha} \bar{N} \mathcal{H} \mathcal{H} + D_{\alpha} \bar{N} \gamma_5 \gamma_7 \mathcal{H} + E_{\alpha} \bar{N} \gamma_5 \gamma_7 \mathcal{H} + E_{\alpha} \bar{N} \gamma_5 \gamma_7 \mathcal{H},$$

(9)

where $D_{\alpha}$ and $E_{\alpha}$ are four low energy constants (LECs). $D_{\alpha}$ and $E_{\alpha}$ contribute to the central potential and spin-spin interaction, respectively. $E_{\alpha}$ and $E_{\beta}$ are related with the isospin-isospin interaction and contribute to the central and spin-spin interaction in spin space, respectively. With the quark model, we fix their values with the $NN$ interaction as inputs, which is given in the Appendix A.

B. Expressions of the effective potentials

In the framework of heavy hadron chiral perturbation theory, the scattering amplitudes of the $D(1^{*})N$ systems can be expanded order by order in powers of a small parameter $\varepsilon = \frac{\rho}{\Lambda_{\chi}}$, where $\rho$ is either the momentum of Goldstone bosons or the residual momentum of heavy hadrons, and $\Lambda_{\chi}$ represents either the chiral breaking scale or the mass of a heavy hadron.
The expansion is organized by the power counting rule in Refs. [54, 55]. The \( \mathcal{O}(\varepsilon^0) \) Feynman diagrams for the \( DN \) and \( D^*N \) systems are shown in Fig. 1, which contain the contact and one-pion-exchange diagrams. The one-pion-exchange diagram for the \( DN \) system vanishes since the \( DD\pi \) vertex is forbidden. The corresponding momentum-space potentials of graphs in Fig. 1 read

\[
\mathcal{V}_{D^N_{1,1}} = D_a - 4E_a(I_1 \cdot I_2),
\]

\[
\mathcal{V}_{D^N_{1,2}} = D_a + D_b \sigma \cdot T - 4(E_a + E_b \sigma \cdot T)(I_1 \cdot I_2),
\]

\[
\mathcal{V}_{D^N_{1,3}} = (I_1 \cdot I_2) \frac{g_s}{f_\pi} \frac{(q \cdot \sigma)(q \cdot T)}{q^2 + m_\pi^2},
\]

where \( I_1 \) and \( I_2 \) are the isospin operators of \( D \) and \( N \), respectively. The operators \( \sigma \) and \( T \) are related to the spin operators of the spin-\( \frac{1}{2} \) baryon, spin-1 meson as \( \frac{1}{2} \sigma \) and \( -T \), respectively (see Ref. [66] for details). The Breit approximation

\[
\mathcal{V}(q) = -\frac{\mathcal{M}(q)}{\sqrt{2\Pi}m_tr_m_f}
\]

is used to relate the scattering amplitude \( \mathcal{M}(q) \) to the effective potential \( \mathcal{V}(q) \) in momentum space (\( m_t \) and \( m_f \) are the masses of the initial and final states, respectively).

The next-to-leading order two-pion-exchange diagrams for the \( DN \) system are illustrated in Fig. 2. The effective potentials from these graphs read

\[
\mathcal{V}_{D^N_{1,1}} = (I_1 \cdot I_2) \frac{1}{f_\pi^2} J_{22}^F(m_\pi, q),
\]

\[
\mathcal{V}_{D^N_{1,2}} = (I_1 \cdot I_2) \frac{g_s^2}{f_\pi^2} \left[ \frac{2}{d-1} J_{24}^T - q^2 \frac{d-2}{d-1} \right] \left( J_{24}^T + J_{33}^T \right)(m_\pi, \mathcal{E}, q),
\]

\[
\mathcal{V}_{D^N_{1,3}} = (I_1 \cdot I_2) \frac{g_s^2}{f_\pi^2} \left[ \frac{2}{d-1} J_{24}^T - q^2 \right] \left( J_{24}^T + J_{33}^T \right)(m_\pi, \mathcal{E}, q),
\]

\[
\mathcal{V}_{D^N_{1,4}} = (I_1 \cdot I_2) \frac{g_s^2}{f_\pi^2} \left[ \frac{2}{d-1} J_{24}^T - q^2 \right] \left( J_{24}^T + J_{33}^T \right)(m_\pi, \mathcal{E}, q),
\]

The effective potentials of the isospin operators of \( D \) and \( N \) systems are illustrated in Fig. 2. The effective potentials of the system are shown in Fig. 3. Their analytical expressions are written as

\[
\mathcal{V}_{D^N_{1,1}} = (I_1 \cdot I_2) \frac{1}{f_\pi^2} J_{22}^F(m_\pi, q),
\]

\[
\mathcal{V}_{D^N_{1,2}} = (I_1 \cdot I_2) \frac{g_s^2}{f_\pi^2} \left[ \frac{2}{d-1} J_{24}^T - q^2 \frac{d-2}{d-1} \right] \left( J_{24}^T + J_{33}^T \right)(m_\pi, \mathcal{E}, \delta_a, q),
\]

\[
\mathcal{V}_{D^N_{1,3}} = (I_1 \cdot I_2) \frac{g_s^2}{f_\pi^2} \left[ \frac{2}{d-1} J_{24}^T - q^2 \right] \left( J_{24}^T + J_{33}^T \right)(m_\pi, \mathcal{E}, \delta_b, q),
\]

\[
\mathcal{V}_{D^N_{1,4}} = (I_1 \cdot I_2) \frac{g_s^2}{f_\pi^2} \left[ \frac{2}{d-1} J_{24}^T - q^2 \right] \left( J_{24}^T + J_{33}^T \right)(m_\pi, \mathcal{E}, \delta_b, q),
\]
FIG. 2. The two-pion-exchange diagrams of the $DN$ system at $O(\varepsilon^2)$. These diagrams are classified as the football diagram ($F_{1,1}$), triangle diagrams ($T_{1,1}$), box diagrams ($B_{1,1}$) and crossed box diagrams ($R_{1,1}$). We use the heavy-thick line to denote the $\Delta(1232)$ in the loops. Other notations are the same as those in Fig. 1.

$V^{F_{1,1}}_{DN} = (I_1 \cdot I_2) \frac{4g_s^2}{3f_{\pi}^2} \left[ (2 - d)J_{34}^T - \frac{q^2}{d - 1} (J_{24}^T + J_{53}^T) \right] (m_\pi, E - \delta_a, q), \quad (25)$

$V^{B_{1,1}}_{DN} = \left[ \frac{1}{8} - \frac{1}{3} I_1 \cdot I_2 \right] \frac{3g_s^2g_\sigma^2}{2f_{\pi}^2} \left[ \frac{4d^2 - 10d + 6}{d - 1} J_{24}^B - q^2 \frac{2d^2 + 3d - 8}{d - 1} (J_{31}^B + J_{42}^B) \right.$

$- q^2 \frac{d - 2 + \sigma \cdot T}{d - 1} J_{21}^B + q^2 \frac{d - 2}{d - 1} (J_{22}^B + 2J_{32}^B + 4J_{43}^B) \left. \right] (m_\pi, E, E, q), \quad (26)$

$V^{B_{2,1}}_{DN} = \left[ \frac{1}{8} - \frac{1}{3} I_1 \cdot I_2 \right] \frac{3g_s^2g_\sigma^2}{2f_{\pi}^2} \left[ -2q^2 \frac{d + 1}{d - 1} (J_{31}^B + J_{42}^B) - q^2 \frac{1}{d - 1} (1 + \sigma \cdot T) J_{21}^B \right.$

$+ (d + 1) J_{41}^B + q^4 \frac{1}{d - 1} (J_{22}^B + 2J_{32}^B + J_{43}^B) \left. \right] (m_\pi, E, E + \delta_b, q), \quad (27)$

$V^{B_{2,2}}_{DN} = \left[ \frac{1}{2} + \frac{2}{3} I_1 \cdot I_2 \right] \frac{g_\sigma^2}{f_{\pi}^2} \left[ -q^2 (d - 2) - \sigma \cdot T J_{21}^B - q^2 \frac{(d - 2)(d^2 + 3d - 8)}{(d - 1)^2} (J_{31}^B + J_{42}^B) \right.$

$+ \frac{2(d^2 - 2d + 2)}{d - 1} J_{41}^B + q^4 \frac{(d - 2)^2}{(d - 1)^2} (J_{22}^B + 2J_{32}^B + J_{43}^B) \left. \right] (m_\pi, E - \delta_a, E, q), \quad (28)$

$V^{B_{2,3}}_{DN} = \left[ \frac{1}{2} + \frac{2}{3} I_1 \cdot I_2 \right] \frac{g_\sigma^2}{f_{\pi}^2} \left[ -q^2 (d + 1)(d - 2) - q^2 \frac{(d - 2)(d^2 + 3d - 8)}{(d - 1)^2} (J_{31}^B + J_{42}^B) \right.$

$+ q^4 \frac{d - 2}{d - 1} (J_{22}^B + 2J_{32}^B + J_{43}^B) + (d^2 - 2d + 2) J_{41}^B \left. \right] (m_\pi, E - \delta_a, E + \delta_b, q). \quad (29)$

FIG. 3. The two-pion-exchange diagrams of the $D^* N$ system at $O(\varepsilon^2)$. The notations are the same as those in Fig. 2.
The expressions of the diagrams ($R_{2,i}$) can be obtained with $\mathcal{V}_{D^{*}N}^{R_{2,i}} = \mathcal{V}_{D^{*}N}^{B_{2,i}} |_{j^p_{1}\rightarrow j^p_{2}, \, 1_{11}\rightarrow 1_{12}, \, \sigma \cdot T \rightarrow -\sigma \cdot T}$.  

III. NUMERICAL RESULTS AND DISCUSSIONS

With the momentum-space potentials $\mathcal{V}(q)$ obtained in Sec. II B, we make the following Fourier transformation to get the effective potential $V(r)$ in the coordinate space,

$$V(r) = \int \frac{d^3q}{(2\pi)^3} e^{iq\cdot r} \mathcal{V}(q) \mathcal{F}(q). \tag{30}$$

We need to introduce a regulator $\mathcal{F}(q)$ to suppress the high momentum contribution. We choose the Gauss form $\mathcal{F}(q) = \exp(-q^{2n}/\Lambda^{2n})$ as used in the $NN$ and $NN$ systems [76, 77]. The power $n = 3$ and cutoff $\Lambda \simeq 0.5 \pm 0.1$ GeV are always adopted to fit the experimental data and make predictions [78–81].

A. Numerical results

In order to get the numerical results, we also need to know the values of the four LECs in Eq. (9). Generally, these LECs should be determined by fitting the $D^{(*)}N$ scattering data in experiments or in lattice QCD simulations. However, the data in this area are scarce, thus we have to resort to other alternative ways. As proposed in Refs. [67, 68], we estimate the LECs by constructing the contact Lagrangian at the quark level, and then extract the couplings from the $NN$ interaction, which is demonstrated in Appendix A.

We show the effective potentials of each possible $I(J^P)$ configurations in Fig. 4. In the following, we analyze the behaviors of effective potentials for each system.

$DN$ system: The result in Fig. 4(a) shows that the $O(e^0)$ contact and $O(e^2)$ two-pion-exchange potentials of the $[DN]_{J=1/2}^{I=0}$ system are both attractive. But the attraction of two-pion-exchange potential is rather weak. The attractive potential is dominantly provided by the contact interaction. We find a bound state in this channel. The binding energy and mass of this state are predicted, respectively,

$$\Delta E_{[DN]_{J=1/2}^{I=0}} \simeq -11.1 \text{ MeV},$$
$$M_{[DN]_{J=1/2}^{I=0}} \simeq 2792.0 \text{ MeV}. \tag{31}$$

For the $[DN]_{J=1/2}^{I=1}$ system in Fig. 4(b), the $O(e^0)$ contact interaction vanishes in our calculation, and the total potential arises from the two-pion-exchange contribution. We notice the potential in this channel is much shallower than that of the $[DN]_{J=1/2}^{I=0}$ channel, i.e., the attraction is too feeble to form the bound state. Thus the binding solution does not exist in this channel.

$D^{*}N$ system: The contact potential of the $[D^{*}N]_{J=1/2}^{I=0}$ system in Fig. 4(c) is attractive, while the one-pion-exchange and two-pion-exchange interaction are both repulsive. Therefore, the total potential is shallower than that of the $[DN]_{J=1/2}^{I=0}$ channel. However, we still obtain a binding solution in the $[D^{*}N]_{J=1/2}^{I=0}$ system. The binding energy and mass of this state are

$$\Delta E_{[D^{*}N]_{J=1/2}^{I=0}} \simeq -1.5 \text{ MeV},$$
$$M_{[D^{*}N]_{J=1/2}^{I=0}} \simeq 2943.6 \text{ MeV}. \tag{32}$$

For the $[D^{*}N]_{J=3/2}^{I=1}$ system in Fig. 4(d), the one-pion-exchange potential is weakly attractive, but the contact and two-pion-exchange potentials are all repulsive. Thus, the total attractive potential is not strong enough to form molecular states in this channel.

For the channel $[D^{*}N]_{J=3/2}^{I=0}$ in Fig. 4(e), the behavior of its potentials is very interesting. We notice the one-pion- and two-pion-exchange contributions almost cancel each other. Thus the total potential is mainly provided by the contact term, which can reach up to $-80$ MeV at the deepest position. By solving the Schrödinger equation, we find the binding solution in the $[D^{*}N]_{J=3/2}^{I=0}$ system, and the binding energy is

$$\Delta E_{[D^{*}N]_{J=3/2}^{I=0}} \simeq -6.7 \text{ MeV}. \tag{33}$$

The corresponding mass of this bound state is

$$M_{[D^{*}N]_{J=3/2}^{I=0}} \simeq 2938.4 \text{ MeV}, \tag{34}$$
which is in good agreement with the mass of $\Lambda_c(2940)$ measured by BaBar, Belle and LHCb (e.g., see Table I).

For the last channel $[D^{*}N]_{J=3/2}^{I=1}$ in Fig. 4(f), the one-pion- and two-pion-exchange potentials almost cancel each other, and the contact contribution is very weakly attractive. Thus no bound state can be found in this channel.

Role of the $\Delta(1232)$: Considering the strong coupling between $\Delta(1232)$ and $NN$, we include the contribution of $\Delta(1232)$ in the loop diagrams (e.g., see Figs. 2 and 3). Here, we discuss the role of $\Delta(1232)$ in the effective potentials of the $DN$ and $D^{*}N$ systems. We try to ignore the effect of $\Delta(1232)$, and notice that the lineshape of the two-pion-exchange potentials changes drastically. Except for the $[DN]_{J=1/2}^{I=0}$, the whole behavior of the other channels is totally reversed. For example, the two-pion-exchange potential of the $[DN]_{J=1/2}^{I=1}$ channel becomes repulsive, which renders the total potential of this channel shallower. But for the $[D^{*}N]_{J=3/2}^{I=2}$ channel, the two-pion-exchange potential becomes attractive. The variation is about $-30$ MeV, which gives rise to a deeper attractive potential, and the binding energy is $-16$ MeV.

In general, the conclusion that there exists the bound state in isoscalar $[D^{(*)}N]_{J}$ channel and no binding solution in isovector channel is robust, no matter we consider the $\Delta(1232)$ or not. However, the $\Delta(1232)$ plays an important role in determining the physical masses of $\Lambda_c(2940)$ and other bound states, since the molecular states are very sensitive to the subtle changes of their internal effective potentials.
B. Discussions

No binding solution in the isovector channel indicates that the \([DN]^{I=1/2}_{J=1/2}\) molecular explanation of \(\Sigma_c(2800)\) is not favored. Although the \(\Sigma_c(2800)\) is near the \(DN\) threshold, its mass is also consistent with the quark model predictions [14, 42, 43, 82, 83]. Thus interpreting the \(\Sigma_c(2800)\) as the \(1P\) charmed baryon seems to be more reasonable.

The situation of \(\Lambda_c(2940)\) is very similar to the \(\Lambda(1405), D_s(2317)\) and \(X(3872)\), i.e., there is large gap between the physical states and quark model predictions \(^1\). Generally, one possible reason is these states per se may be exotic rather than the conventional ones.

A recent analysis from LHCb gives weak constraints on the \(J^P\) quantum numbers of \(\Lambda_c(2940)\), where \(J^P = \frac{3}{2}^-\) is favored [12]. This is consistent with our calculations. Actually, one can notice two peaks in the \(D^0p\) invariant mass spectrum from 2.92 GeV to 2.99 GeV in the results of LHCb (see Fig. 13(a) in Ref. [12]). The one at 2.94 GeV is just the reported \(\Lambda_c(2940)\). The other peak at 2.98 GeV may correspond to the true \(\Lambda_c(2P)\) baryon, since its mass is close to the quark model prediction [13–16]. Our calculation indicates the \(\Lambda_c(2940)\) is probably the \(S\)-wave \(D^*N\) molecular state.

We report three bound states in the \([DN]_{J=1/2}^{I=0}\), \([D^*N]_{J=1/2}^{I=0}\) and \([D^*N]_{J=3/2}^{I=0}\) systems. They are very similar to the newly observed \(P_c(4312), P_c(4440)\) and \(P_c(4457)\) at LHCb [85], which are interpreted as the \([\bar{D}\Sigma_c]_{J=1/2}^{I=1/2}\), \([\bar{D}^*\Sigma_c]_{J=1/2}^{I=1/2}\) and \([\bar{D}^*\Sigma_c]_{J=3/2}^{I=1/2}\) molecular states [65, 66], respectively. If \(\Lambda_c(2940)\) is indeed the \(D^*N\) molecular state, then it should contain two structures, i.e., \([D^*N]_{J=1/2}^{I=0}\) and \([D^*N]_{J=3/2}^{I=0}\). Because the mass splitting between the spin-\(\frac{1}{2}\) and spin-\(\frac{3}{2}\) states is only about 5 MeV, it is very difficult to disassemble these two structures with current accuracy. Similar situation has happened to the \(P_c\) states. The previously reported \(P_c(4450)\) [86] contains two structures, \(P_c(4440)\) and \(P_c(4457)\), after increasing the data sample. More interesting, we find the mass of the spin-\(\frac{3}{2}\) state is larger than that of the spin-\(\frac{1}{2}\) one.

The signal of \(\Lambda_c(2940)\) has been observed in the \(D^0p\) and \(\Sigma_c\pi\) final states [10–12]. However, if the \(J^P\) of \(\Lambda_c(2940)\) is \(\frac{3}{2}^-\) as weakly constrained by the LHCb, then it decays into the \(D^0p\) and \(\Sigma_c\pi\) through the \(D\)-wave, which is strongly suppressed\(^2\). Therefore, as mentioned above, one promising explanation is that the \(\Lambda_c(2940)\) signal actually contains two structures. The spin-\(\frac{1}{2}\) structure can easily decay into \(D^0p\) and \(\Sigma_c\pi\) via the \(S\)-wave.

Borrowing experiences from the discovery of \(P_c\) states, we urge the experimenters to reanalyze the \(\Lambda_c(2940)\) \(\rightarrow D^0p\) signal actually contains two structures. The spin-\(\frac{1}{2}\) structure can easily decay into \(D^0p\) and \(\Sigma_c\pi\) via the \(S\)-wave.

\(^1\)An unquenched study with the channel coupling in Ref. [84] declares the mass of \(\Lambda_c(2P,3/2^-)\) state can be lowered down to match the experimental data of \(\Lambda_c(2940)\).

\(^2\)Based on the \(3P_0\) model calculation, the large decay width of \(\Lambda_c(2940)\) \(\rightarrow D^0p\) is reported in Refs. [87, 88] by treating \(\Lambda_c(2940)\) as a \(2P\) state in the \(\Lambda_c\) family, but the low mass puzzle still exist.
In addition to the mass spectrum, the decay pattern can also give us some important criteria to identify the inner structure of \( \Lambda_c(2940) \). In the molecular scenario, the \( D^*N \) system can easily decay into the \( DN \) channel via the pion exchange, while the \( \Sigma_c\pi \) decay mode requires the exchange of a nucleon or a \( D \) meson. Thus, the decay amplitude of the \( D^0\pi^0 \) mode should be much larger than that of the \( \Sigma_c\pi \), because the heavy hadron exchange is generally suppressed. However, the phase space of the \( \Sigma_c\pi \) mode is larger.

The three body decay mode is also very interesting. We take the decay modes of the \( X(3872) \) and other higher charmonia as an example. The branching fraction of \( X(3872) \rightarrow D^0\pi^0 \) can reach up to 40\% [1]. In contrast, the open charm three body decays of the higher charmonia is only a few percents [89]. Analogously, the branching fraction of \( \Lambda_c(2940) \rightarrow D^0\pi^0(\gamma)p \) should also be conspicuous in the molecular picture.

Besides, our study can be easily extended to the \( \bar{B}^*(s)N \) systems. The axial coupling \( g \) and mass splitting \( \delta_B \) in Eq. (7) should be replaced by the bottomed ones, where we adopt \( g = -0.52 \) [90, 91] and \( \delta_B = 45 \) MeV [1]. The predicted results are listed in Table. II. There also exist bound states in the isoscalar \([\bar{B}^*(s)N]_J^I\) systems. These states might be reconstructed at the \( \Lambda^*_{c0}\pi^+\pi^- \) final states, and the \([\bar{B}^*N]_J^I\) states could also be detected in the \( \bar{B}^-p \) mass spectrum.

### Table II. The predicted binding energies and masses for the isoscalar \([D^*(s)N]_J^I\) and \([\bar{B}^*(s)N]_J^I\) systems (in units of MeV).

| System | \([D^*N]_J^I\) | \([D^*N]_J^I\) | \([D^*N]_J^I\) | \([\bar{B}N]_J^I\) | \([\bar{B}^*N]_J^I\) | \([\bar{B}^*N]_J^I\) |
|--------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| \(\Delta E\) | -11.1 | -1.5 | -6.7 | -8.8 | -3.5 | -8.4 |
| Mass | 2792.0 | 2943.6 | 2938.4 | 6208.8 | 6259.4 | 6254.5 |

### IV. SUMMARY

A sophisticated investigation on the \( D_N \) and \( D^*N \) interactions is crucial to clarify the nature of the charmed baryons \( \Sigma_c(2800) \) and \( \Lambda_c(2940) \). In this work, we systematically study the effective potentials of the \( D_N \) and \( D^*N \) systems with the chiral effective field theory up to the next-to-leading order. We simultaneously consider the contributions of the long-range one-pion-exchange, mid-range two-pion-exchange and short-range contact term. We also include the \( \Delta(1232) \) as an intermediate state in the loop diagrams. The LECs are estimated from the \( N\bar{N} \) interaction with the help of quark model.

For the \( DN \) system, our calculation shows the effective potentials of the \( [D^*N]_{J=1/2}^I \) and \( [DN]_{J=1/2}^I \) channels are both attractive. We find a bound state in the \( [DN]_{J=0}^I \) channel, but the attraction in the \( [DN]_{J=1}^I \) channel is too weak to form a bound state. Thus the explanation of \( \Sigma_c(2800) \) as the \( DN \) molecular state is disfavored in our calculations. The \( \Sigma_c(2800) \) is more likely to be the conventional \( 1P \) charmed baryon, since its mass is well consistent with the quark model prediction.

There are four channels in the \( [D^*N]_J^I \) system. We find only the isoscalar \( [D^*N]^I_J \) potential is deep enough to form the molecular state. We obtain the masses of the bound states in the \( [D^*N]_{J=0}^I \) and \( [D^*N]_{J=1}^I \) channels to be 2943.6 and 2938.4 MeV, respectively, which well accord with the BaBar, Belle and LHCb measurements for \( \Lambda_c(2940) \). Considering the small mass splitting between the spin-\( \frac{1}{2} \) and spin-\( \frac{3}{2} \) states, we conjecture the \( \Lambda_c(2940) \) signal contains two structures.

It is not so easy to squeeze the \( \Lambda_c(2940) \) into the conventional charmed baryon spectrum, since the 60–100 MeV gap between the physical mass and quark model prediction cannot be readily remedied. However, this problem can be easily reconciled in the molecular picture, i.e., the \( \Lambda_c(2940) \) is probably the isoscalar \( D^*N \) molecule rather than the \( 2P \) charmed baryon.

We also investigate the influence of \( \Delta(1232) \) in the loop diagrams. The binding solutions always exist in the isoscalar \( [D^*N]^I_J \) channels no matter we include the \( \Delta(1232) \) or not. There still do not exist bound states in the isovector channels even we ignore the \( \Delta(1232) \). However, the \( \Delta(1232) \) is important in yielding the shallowly bound isoscalar \( [D^*(s)N]^I_J \) states.

We hope experimentalist could seek for the pentaquark candidates in the open charmed channels, where the \( D^*N \) molecular pentaquarks in the isoscalar systems might be reconstructed at the \( \Lambda^*_{c0}\pi^+\pi^- \) final state.

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### Appendix A: Determining the LECs from \( N\bar{N} \) interaction

One needs to know the values of the LECs in Eq. (9) to study the strength of the short-range interaction. As proposed in Refs. [67, 68] (more details can be found in the appendix of these two references), the LECs of \( D^*(s)N \) systems can be bridged to those of the \( N\bar{N} \) interaction with the help of quark model. The way is analogous to the resonance saturation model [92], but we build the quark level Lagrangian. We assume the contact interaction stems from the heavy meson exchanging. We introduce \( \mathcal{J} \) and \( \mathcal{J}^\mu \) to produce the central potential and spin-spin interaction, respectively. The matrix form of \( \mathcal{J} \) and \( \mathcal{J}^\mu \) can be expressed as

\[
\mathcal{J} = \mathcal{J}_3 \tau^1 + \sqrt{\frac{1}{3}} \mathcal{J}_1, \tag{A1}
\]

\[
\mathcal{J}^\mu = \mathcal{J}_3^{\mu \tau^1} + \sqrt{\frac{1}{3}} \mathcal{J}_1^\mu, \tag{A2}
\]
where $J_3$ and $J_1$ denote the isospin triplet and isospin singlet, respectively. The coefficient $\sqrt{\frac{1}{3}}$ is introduced to satisfy the SU(3) flavor symmetry.

The $q\bar{q}$ contact potential can be written as

$$V_{q\bar{q}} = c_q(1 - 3\tau_1 \cdot \tau_2) + c_t(1 - 3\tau_1 \cdot \tau_2)\sigma_1 \cdot \sigma_2, \quad \text{(A3)}$$

where $c_q$ and $c_t$ are the coupling constants. The minus sign in Eq. (A3) arises since the isospin triplet and the isospin singlet have different $G$-parities.

With the $q\bar{q}$ contact potential $V_{q\bar{q}}$ in Eq. (A3) and the relevant matrix element in Table III, we obtain the $NN$ contact potential as follows,

$$V_{NN} = \langle N\bar{N}|V_{q\bar{q}}|NN\rangle = 9c_q - 3c_t \cdot \tau_2$$

$$+ c_t\sigma_1 \cdot \sigma_2 - \frac{25}{3}c_t(\tau_1 \cdot \tau_2)(\sigma_1 \cdot \sigma_2). \quad \text{(A4)}$$

Similarly, the $D^*N$ contact potential can be easily worked out,

$$V_{D^*N} = \langle D^*N|V_{q\bar{q}}|D^*N\rangle = 3c_q - 12c_t\cdot I_1 \cdot I_2$$

$$- c_t\sigma \cdot T + 20c_t(\sigma \cdot I_1 \cdot I_2)(\sigma \cdot T). \quad \text{(A5)}$$

Matching Eq. (11) and Eq. (A5) one can get the LECs in Eq. (9), which read

$$D_a = 3c_q, \quad D_b = -c_t, \quad E_a = 3c_q, \quad E_b = -5c_t. \quad \text{(A6)}$$

Therefore, once we know the values of $c_q$ and $c_t$, we can capture the short range interaction of the $D^*NN$ systems. The $c_q$ and $c_t$ can be extracted from the $NN$ interaction, and the $NN$ scattering phase shift has been fitted in the framework of chiral effective field theory to the next-to-next-to-leading order in Ref. [78]. Using the values of $C_{3S1}$, in the $I = 0$ and $I = 1$ channels fitted at ($\Lambda, \bar{\Lambda}$) = (450, 500) MeV as inputs, we obtain

$$c_q = -8.1 \text{ GeV}^{-2}, \quad c_t = 0.65 \text{ GeV}^{-2}. \quad \text{(A7)}$$

We notice $|c_q|/|c_t| \approx 12.5$, i.e., the spin-spin interaction only serves as a perturbation to give mass splittings between spin multiplets.

### Table III. The matrix elements of the operator $\sum_i h_i c_{I_i} c_{I_0} \Theta_{IJ}$, where $h_a$ and $h_b$ are two hadrons. $\Theta_{IJ}$ is the two-body interaction operator between quarks.

| $[NN]_{J=0}^{I=1}$ | $[D^*N]_{J=3/2}^{I=1}$ |
|---------------------|---------------------|
| $9$ | $3$ | $1$ | $1$ | $-25/3$ | $\frac{5}{4}$ |
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