$J = 3/2$ charmed hypertriton

H. Garcilazo[*]
Escuela Superior de Física y Matemáticas, Instituto Politécnico Nacional, Edificio 9, 07738 México D.F., Mexico

A. Valcarce,[†] and T. F. Caramés[†]
Departamento de Física Fundamental and IUFFyM, Universidad de Salamanca, E-37008 Salamanca, Spain
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By solving exact three-body equations, we study the three-baryon system with charm +1. We look for possible bound states using baryon-baryon interactions obtained from a chiral constituent quark model. The smaller effect of the $\Lambda_c \leftrightarrow \Sigma_c$ conversion reverses the order of the $(I,J) = (0,1/2)$ and $(I,J) = (0,3/2)$ states, rather close on the strange sector. The diminishing of the kinetic energy due to the large reduced mass gives rise to a bound state in the $(I,J) = (0,3/2)$ channel. After correcting for Coulomb effects, the binding energy would be between 140 and 715 keV.

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I. INTRODUCTION

Soon after the discovery of baryons possessing net charm it was suggested that there should also exist charmed nuclei. The observation of a candidate event that could be interpreted in terms of the decay of a charmed nucleus [1] fostered conjectures about the possible existence of charm analogs of strange hypernuclei [2–4]. Three ambiguous candidates of charmed hypernuclei were reported by an emulsion experiment with 250-GeV protons [5]. This gave rise to several theoretical estimates about the binding energies and the potential-well depth of charmed hypernuclei based on one-boson-exchange potentials for the charmed baryon-nucleon potential [6–10]. There were also theoretical estimations of the production of cross sections as well as experimental condition requirements for producing charmed hypernuclei by means of charmed exchange reactions on nuclei [11]. The experiments that may search for charmed hypernuclei are becoming realistic and may be performed in coming years at Hall C of JPARC [12], at the FAIR experiment [13], or at the Super$B$ collider [14]. All these experimental prospects have reinvigorated twenty years later the study of charmed hypernuclei [15–17] and also more recently theoretical studies of charmed dibaryons [18,19]. We show in Fig. 1 a diagram to generate Λ_c^+ baryons by means of antiproton collisions on the deuteron through the intermediate production of charged D mesons ($m_{D^\pm} = 1869.61$ MeV/c^2), feasible at modern factories, that has been proposed to study the existence of charmed hypernuclei at JPARC [12]. A similar reaction, $D^+ + p \to \Lambda_c^+ + \pi^+$, is proposed at Super$B$ to detect charmed supernuclei [14].

We have developed in the past the exact formalism to study three-baryon systems with a heavy flavor baryon [20]. It is our purpose in this work to extend our previous study to three-baryon systems with a unit of charm by considering the region of the $\Lambda_c NN$ bound states, where a charmed hypertriton might exist, and calculating for the first time the $\Lambda_c d$ and $\Sigma_c d$ scattering lengths. We will simultaneously study all $\Lambda_c NN$ and $\Sigma_c NN$ states with $J = 1/2,3/2$ and $I = 0,1,2$.

II. FORMALISM

Let us start by summarizing the description of the three-baryon system with a charmed baryon. We transform the Faddeev equations from being integral equations in two continuous variables into integral equations in just one continuous variable by expanding the two-body t matrices in terms of Legendre polynomials, P_m [21],

$$t_{i;s_i i_i}(p_i, p_i'; e) = \sum_{nr} P_n(x_i)\tau_{i;s_i i_i}^{nr}(e)P_r(x_i'), \tag{1}$$

where $x_i = \frac{p_i - b}{p_i + b}$, $x_i' = \frac{p_i' - b}{p_i' + b}$, p_i, and p_i' are the initial and final relative momenta of the pair jk, and b is a scale parameter of which the results do not depend on.

If we identify particle 1 with the charmed baryon and particles 2 and 3 with the two nucleons, the integral equations for βd scattering at threshold, with $\beta = \Sigma_c$ or Λ_c, are in the case of pure S-wave configurations

$$\begin{aligned} T_{2;SI;\beta}^{ns_2i_2}(q_2) = {} & B_{2;SI;\beta}^{ns_2i_2}(q_2) \\ & + \sum_{ms_3i_3} \int_0^\infty dq_3 \bigg[(-1)^{1+\sigma_1+\sigma_3-s_2+\tau_1+\tau_3-i_2} \\ & \times A_{23;SI}^{ns_2i_2ms_3i_3}(q_2,q_3;E) \\ & + 2\sum_{rs_1i_1} \int_0^\infty dq_1\, A_{31;SI}^{ns_2i_2rs_1i_1}(q_2,q_1;E) \\ & \times A_{13;SI}^{rs_1i_1ms_3i_3}(q_1,q_3;E) \bigg] T_{2;SI}^{ms_3i_3}(q_3), \end{aligned} \tag{2}$$

where σ_1 (τ_1) and σ_3 (τ_3) stand for the spin (isospin) of the charmed baryon and the nucleon respectively, while s_i and i_i are the spin and isospin of the pair jk. $T_{2;SI;\beta}^{ns_2i_2}(q_2)$ is a two-component vector,

$$T_{2;SI;\beta}^{ns_2i_2}(q_2) = \begin{pmatrix} T_{2;SI;\Sigma_c\beta}^{ns_2i_2}(q_2) \\ T_{2;SI;\Lambda_c\beta}^{ns_2i_2}(q_2) \end{pmatrix}, \tag{3}$$

[*]humberto@esfm.ipn.mx
[†]valcarce@usal.es

while the kernel of Eq. (2) is a 2×2 matrix defined by

$$A^{ns_2i_2ms_3i_3}_{23;SI}(q_2,q_3;E) = \begin{pmatrix} A^{ns_2i_2ms_3i_3}_{23;SI;\Sigma_c\Sigma_c}(q_2,q_3;E) & A^{ns_2i_2ms_3i_3}_{23;SI;\Sigma_c\Lambda_c}(q_2,q_3;E) \\ A^{ns_2i_2ms_3i_3}_{23;SI;\Lambda_c\Sigma_c}(q_2,q_3;E) & A^{ns_2i_2ms_3i_3}_{23;SI;\Lambda_c\Lambda_c}(q_2,q_3;E) \end{pmatrix},$$

$$A^{ns_2i_2rs_1i_1}_{31;SI}(q_2,q_1;E) = \begin{pmatrix} A^{ns_2i_2rs_1i_1}_{31;SI;\Sigma_cN(\Sigma_c)}(q_2,q_1;E) & A^{ns_2i_2rs_1i_1}_{31;SI;\Sigma_cN(\Lambda_c)}(q_2,q_1;E) \\ A^{ns_2i_2rs_1i_1}_{31;SI;\Lambda_cN(\Sigma_c)}(q_2,q_1;E) & A^{ns_2i_2rs_1i_1}_{31;SI;\Lambda_cN(\Lambda_c)}(q_2,q_1;E) \end{pmatrix},$$

$$A^{rs_1i_1ms_3i_3}_{13;SI}(q_1,q_3;E) = \begin{pmatrix} A^{rs_1i_1ms_3i_3}_{13;SI;N\Sigma_c}(q_1,q_3;E) & 0 \\ 0 & A^{rs_1i_1ms_3i_3}_{13;SI;N\Lambda_c}(q_1,q_3;E) \end{pmatrix}, \tag{4}$$

where

$$A^{ns_2i_2ms_3i_3}_{23;SI;\alpha\beta}(q_2,q_3;E) = h^{s_2i_2s_3i_3}_{23;SI}\sum_r \tau^{nr}_{2;s_2i_2;\alpha\beta}\big(E - q_2^2/2\nu_2\big)\frac{q_3^2}{2}\int_{-1}^{1} d\cos\theta\, \frac{P_r(x_2')P_m(x_3)}{E + \Delta E\,\delta_{\beta\Lambda_c} - p_3^2/2\mu_3 - q_3^2/2\nu_3 + i\epsilon}; \quad \alpha,\beta = \Sigma_c,\Lambda_c,$$

$$A^{ns_2i_2ms_1i_1}_{31;SI;\alpha N(\beta)}(q_2,q_1;E) = h^{s_2i_2s_1i_1}_{31;SI}\sum_r \tau^{nr}_{3;s_2i_2;\alpha\beta}\big(E - q_2^2/2\nu_2\big)\frac{q_1^2}{2}\int_{-1}^{1} d\cos\theta\, \frac{P_r(x_3')P_m(x_1)}{E + \Delta E\,\delta_{\beta\Lambda_c} - p_1^2/2\mu_1 - q_1^2/2\nu_1 + i\epsilon}; \quad \alpha,\beta = \Sigma_c,\Lambda_c,$$

$$A^{ns_1i_1ms_3i_3}_{13;SI;N\beta}(q_1,q_3;E) = h^{s_1i_1s_3i_3}_{13;SI}\sum_r \tau^{nr}_{1;s_1i_1;NN}\big(E + \Delta E\,\delta_{\beta\Lambda_c} - q_1^2/2\nu_1\big)\frac{q_3^2}{2}\int_{-1}^{1} d\cos\theta\, \frac{P_r(x_1')P_m(x_3)}{E + \Delta E\,\delta_{\beta\Lambda_c} - p_3^2/2\mu_3 - q_3^2/2\nu_3 + i\epsilon};$$
$$\beta = \Sigma_c,\Lambda_c, \tag{5}$$

with the isospin and mass of particle 1 (the charmed baryon) being determined by the subindex β. μ_i and ν_i are the usual reduced masses and the subindex $\alpha N(\beta)$ indicates a transition $\alpha N \to \beta N$ with a nucleon as spectator followed by a $NN \to NN$ transition with β as spectator. $\tau^{nr}_{2;s_2i_2;\alpha\beta}(e)$ are the coefficients of the expansion in terms of Legendre polynomials of the charmed baryon-nucleon t-matrix $t_{2;s_2i_2;\alpha\beta}(p_2,p_2';e)$ for the transition $\alpha N \to \beta N$, i.e.,

$$\tau^{nr}_{i;s_ii_i;\alpha\beta}(e) = \frac{2n+1}{2}\,\frac{2r+1}{2}\int_{-1}^{1} dx_i \int_{-1}^{1} dx_i'\, P_n(x_i)t_{i;s_ii_i;\alpha\beta} \times (p_i,p_i';e)P_r(x_i'). \tag{6}$$

The energy shift ΔE, which is usually taken as $M_\alpha - M_\beta$, will be chosen instead such that at the βd threshold the momentum of the αd system has the correct value, i.e.,

$$\Delta E = \frac{[(m_\beta+m_d)^2-(m_\alpha+m_d)^2][(m_\beta+m_d)^2-(m_\alpha-m_d)^2]}{8\mu_{\alpha d}(m_\beta+m_d)^2}, \tag{7}$$

where $\mu_{\alpha d}$ is the αd reduced mass.

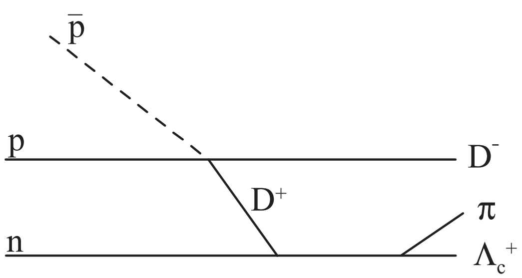

FIG. 1. Diagram to generate Λ_c^+ by means of antiproton collisions on the deuteron.

The inhomogeneous term of Eq. (2), $B^{ns_2i_2}_{2;SI;\beta}(q_2)$, is a two-component vector

$$B^{ns_2i_2}_{2;SI;\beta}(q_2) = \begin{pmatrix} B^{ns_2i_2}_{2;SI;\Sigma_c\beta}(q_2) \\ B^{ns_2i_2}_{2;SI;\Lambda_c\beta}(q_2) \end{pmatrix}, \tag{8}$$

where

$$B^{ns_2i_2}_{2;SI;\alpha\beta}(q_2) = h^{s_2i_210}_{31;SI}\phi_d(q_2) \times \sum_r \tau^{nr}_{2;s_2i_2;\alpha\beta}\big(E^{th}_\beta - q_2^2/2\nu_2\big)P_r(x_2'). \tag{9}$$

$h^{s_2i_2s_1i_1}_{31;SI}$ with $s_1 = 1$ and $i_1 = 0$ are the spin-isospin transition coefficients corresponding to a charmed baryon-deuteron initial state (see Eq. (30) of Ref. [21]), $\phi_d(q_2)$ is the deuteron wave function, E^{th}_β is the energy of the βd threshold, $P_r(x_2')$ is a Legendre polynomial of order r, and

$$x_2' = \frac{\frac{\eta_2}{m_3}q_2 - b}{\frac{\eta_2}{m_3}q_2 + b}. \tag{10}$$

Finally, after solving the inhomogeneous set of Eqs. (2), the βd scattering length is given by

$$A_{\beta d} = -\pi\mu_{\beta d}T_{\beta\beta}, \tag{11}$$

with

$$T_{\beta\beta} = 2\sum_{ns_2i_2} h^{10s_2i_2}_{13;SI}\int_0^\infty q_2^2dq_2\phi_d(q_2)P_n(x_2')T^{ns_2i_2}_{2;SI;\beta\beta}(q_2). \tag{12}$$

In the case of the Σ_cNN system, even for energies below the Σ_cd threshold, one encounters the three-body singularities of the Λ_cNN system so that to solve the integral equations (2) one has to use the contour rotation method where the momenta are rotated into the complex plane $q_i \to q_ie^{-i\phi}$ since as pointed

TABLE I. Two-body $\Sigma_c N$ channels $(i_{\Sigma_c}, s_{\Sigma_c})$, $\Lambda_c N$ channels $(i_{\Lambda_c}, s_{\Lambda_c})$, NN channels with Σ_c spectator $(i_{N(\Sigma_c)}, s_{N(\Sigma_c)})$, and NN channels with Λ_c spectator $(i_{N(\Lambda_c)}, s_{N(\Lambda_c)})$ that contribute to a given $\Sigma_c NN - \Lambda_c NN$ state with total isospin I and spin J.

I	J	$(i_{\Sigma_c}, s_{\Sigma_c})$	$(i_{\Lambda_c}, s_{\Lambda_c})$	$(i_{N(\Sigma_c)}, s_{N(\Sigma_c)})$	$(i_{N(\Lambda_c)}, s_{N(\Lambda_c)})$
0	1/2	(1/2,0),(1/2,1)	(1/2,0),(1/2,1)	(1,0)	(0,1)
1	1/2	(1/2,0),(3/2,0),(1/2,1),(3/2,1)	(1/2,0),(1/2,1)	(0,1),(1,0)	(1,0)
2	1/2	(3/2,0),(3/2,1)		(1,0)	
0	3/2	(1/2,1)	(1/2,1)		(0,1)
1	3/2	(1/2,1),(3/2,1)	(1/2,1)	(0,1)	
2	3/2	(3/2,1)			

out in Ref. [21] the results do not depend on the contour rotation angle ϕ.

III. RESULTS AND DISCUSSION

In order to solve the integral equations (2) for the coupled $\Sigma_c NN - \Lambda_c NN$ system we consider all configurations where the baryon-baryon subsystems are in an S wave and the third particle is also in an S wave with respect to the pair. However, to construct the two-body t matrices that serve as input of the Faddeev equations we considered the full interaction including the contribution of the D waves and of course the coupling between the $\Sigma_c N$ and $\Lambda_c N$ subsystems (this is known as the truncated t-matrix approximation [22]). This approximation in the case of the NNN system with the NN interaction taken as the Reid soft-core potential leads to a triton binding energy which differs less than 1 MeV from the exact value [23]. We give in Table I the two-body channels that are included in our calculation. For a given three-body state (I,J) the number of two-body channels that enter is determined by the triangle selection rules $|J - \frac{1}{2}| \leqslant s_i \leqslant J + \frac{1}{2}$ and $|I - \frac{1}{2}| \leqslant i_i \leqslant I + \frac{1}{2}$. For the parameter b we found that $b = 3$ fm^{-1} leads to very stable results while for the expansion (1) we took twelve Legendre polynomials, i.e., $0 \leqslant n \leqslant 11$.

The two-body interactions are obtained from the chiral constituent quark model of Ref. [24]. The NN potentials perfectly describe the S-wave phase shifts [25] and had also been used in the study of the $\Lambda NN - \Sigma NN$ coupled channel problem [20,21]. The charmed baryon-nucleon potential is derived as explained in Ref. [26]. For the case of heavy quarks (c or b) chiral symmetry is explicitly broken and therefore boson exchanges do not contribute. The absence of strange quarks also eliminates the contribution of K or κ exchanges as compared to the strange baryon case. This simplifies a lot the interaction, see Eqs. (8)–(13) of Ref. [26], and gives rise to a potential whose only free parameter would be the harmonic oscillator width of the charm quark. We will present our results for different values of b_c to get parameter-free predictions. No bound states are found for the charmed two-body subsystems.

TABLE II. $\Lambda_c d$, $A_{0,3/2}$, and $A_{0,1/2}$ and $\Sigma_c d$, $A'_{1,3/2}$, and $A'_{1,1/2}$ scattering lengths, in fm.

$A_{0,3/2}$	$A_{0,1/2}$	$A'_{1,3/2}$	$A'_{1,1/2}$
−10.27	9.02	0.74+ i 0.18	2.08+ i 0.47

There are several facts that should be noted before presenting our results and must be considered for our discussion. The kinetic energy associated to the Λ_c^+ is reduced compared with that of the Λ, which would imply that the Λ_c^+ would be more strongly bound than the Λ in the case of having identical interactions with the nucleons. However, the charmed baryon-nucleon interaction is weaker than that of the strange sector due to the absence of the strange boson exchanges, as also noted in Ref. [8]. Finally, note that the Λ_c^+ has a positive charge, whereas the Λ is neutral. Therefore, Coulomb effects may play a non-negligible role in charmed nuclei. In fact, the hypertriton would be unbound if the Λ were to have a positive charge.

Bearing these considerations in mind, we have first proceeded to solve the Faddeev equations for the $\Lambda_c NN$ and $\Sigma_c NN$ systems using the charmed baryon-nucleon and nucleon-nucleon interactions derived from the chiral constituent quark model with full inclusion of the $\Lambda_c \leftrightarrow \Sigma_c$ conversion. Let us first present the results for the $\Lambda_c d$ and $\Sigma_c d$ scattering lengths, shown in Table II. Although it might be difficult to measure them in the near future, they neatly informed us about the possible existence of bound states in the different channels. The $\Sigma_c d$ scattering lengths are complex since the inelastic $\Lambda_c NN$ channels are always open. Both scattering lengths $A'_{1,1/2}$ and $A'_{1,3/2}$, have a positive real part, indicating that the interaction is repulsive. The spin 1/2 $\Lambda_c d$ scattering length, $A_{0,1/2}$, is also positive. However, the spin 1/2 $\Lambda_c d$ scattering length, $A_{0,3/2}$, is negative, giving rise to a bound state with an energy of 271 keV.

We show in Fig. 2 the Fredholm determinant of the different $\Lambda_c NN$ (I,J) states.[1] The $I = 1$ channels are repulsive; only the $I = 0$ channels present attraction. Curiously, as already noted in the scattering lengths, the order of $I = 0$ channels is reversed with respect to the strange sector, with $J = 3/2$ being the most attractive one. This difference can be easily understood due to the importance of the $\Lambda \leftrightarrow \Sigma$ conversion in the strange sector [27]. The contribution of the $\Lambda \leftrightarrow \Sigma$ conversion should be even stronger than $\Delta \leftrightarrow N$ conversion in ordinary nuclei, because it is not suppressed in S waves and the $\Lambda - \Sigma$ mass difference is much smaller. When the

[1]In all $I = 0$ and $I = 1$ cases the zero energy corresponds to the binding energy of the deuteron below the corresponding threshold, $\Lambda_c NN$ or $\Sigma_c NN$. The deuteron is perfectly reproduced by our model for the NN interaction [25]. For $I = 2$ the zero energy corresponds to the $\Sigma_c NN$ mass because the deuteron channel does not contribute (see Table I).

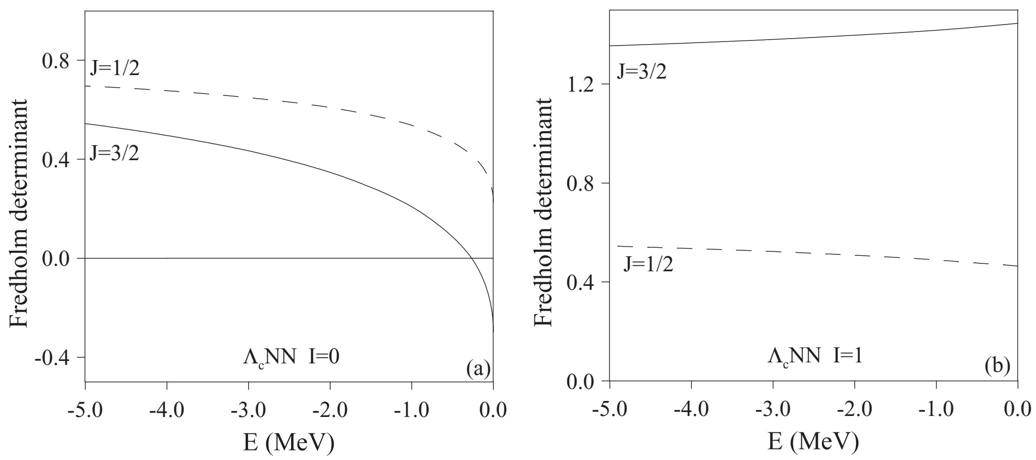

FIG. 2. (a) Fredholm determinant for the $J = 1/2$ and $J = 3/2$ $I = 0$ $\Lambda_c NN$ channels. (b) Fredholm determinant for the $J = 1/2$ and $J = 3/2$ $I = 1$ $\Lambda_c NN$ channels.

$N\Lambda \leftrightarrow N\Sigma$ potential is disconnected, the $J = 3/2$ channel is almost not modified while the $J = 1/2$ losses great part of its attraction. Thus, the ordering between the $J = 1/2$ and $J = 3/2$ channels is reversed in such a way that the hypertriton would not be bound (see Fig. 6(a) of Ref. [20]). The $\Lambda_c \leftrightarrow \Sigma_c$ conversion is less important than in the strange sector first

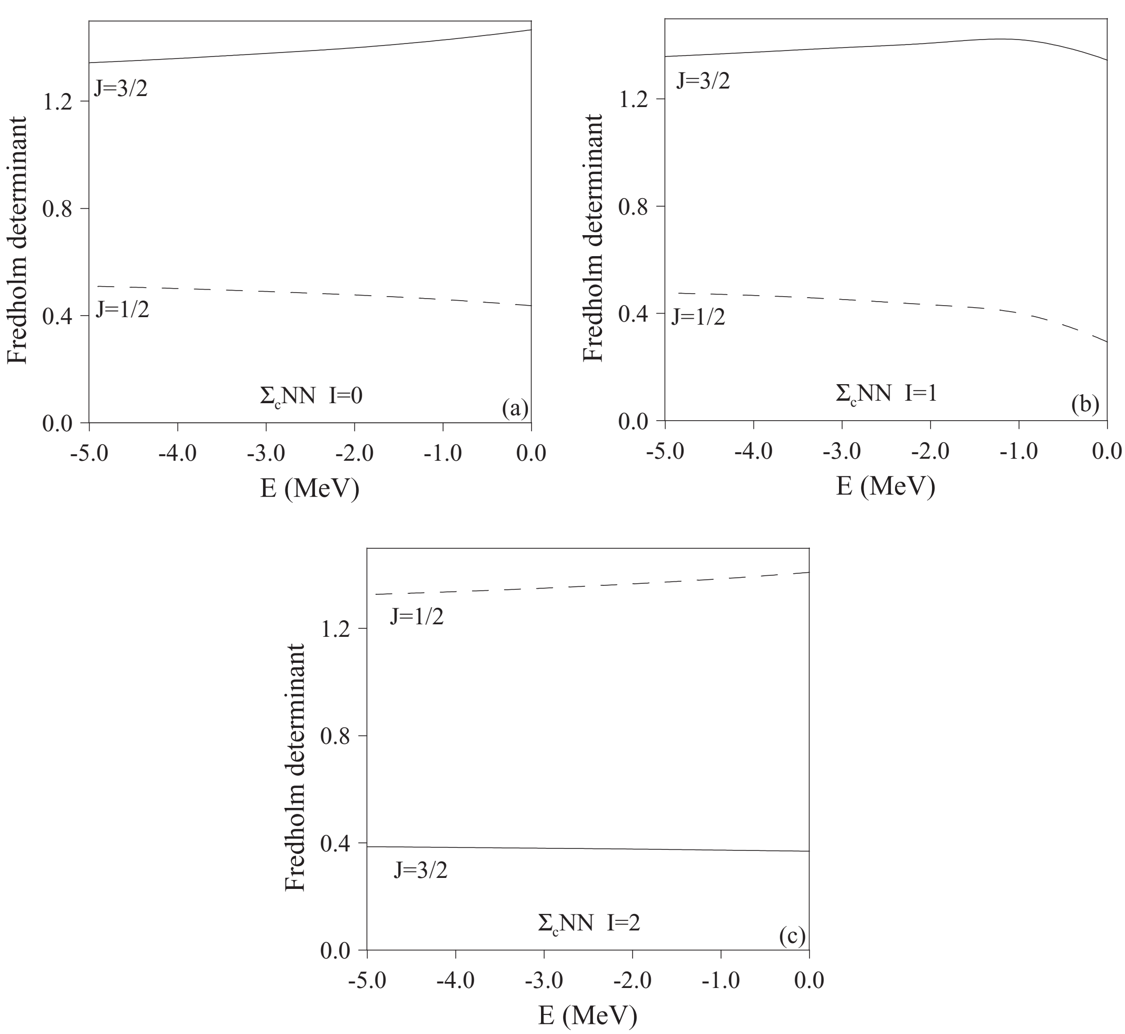

FIG. 3. (a) Fredholm determinant for the $J = 1/2$ and $J = 3/2$ $I = 0$ $\Sigma_c NN$ channels. (b) Fredholm determinant for the $J = 1/2$ and $J = 3/2$ $I = 1$ $\Sigma_c NN$ channels. (c) Fredholm determinant for the $J = 1/2$ and $J = 3/2$ $I = 2$ $\Sigma_c NN$ channels.

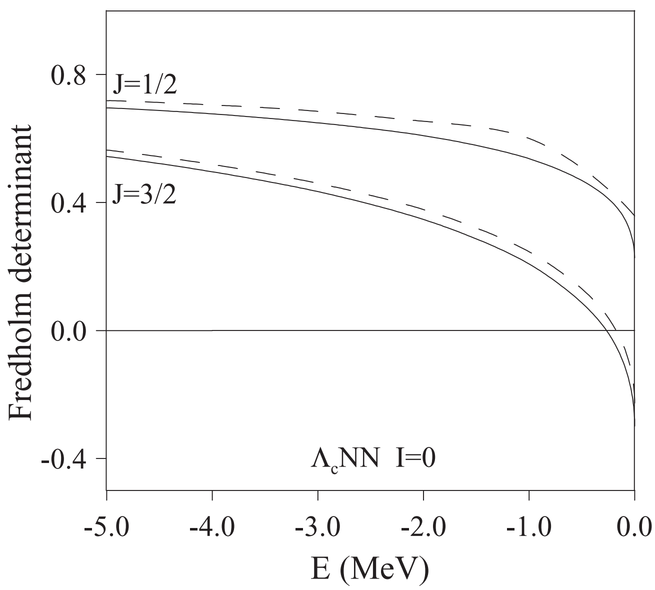

FIG. 4. Fredholm determinant for the $J = 1/2$ and $J = 3/2$ $I = 0$ $\Lambda_c NN$ channels. The solid line stands for the full calculation and the dashed line for when the $\Lambda_c \leftrightarrow \Sigma_c$ transition is taken to be zero.

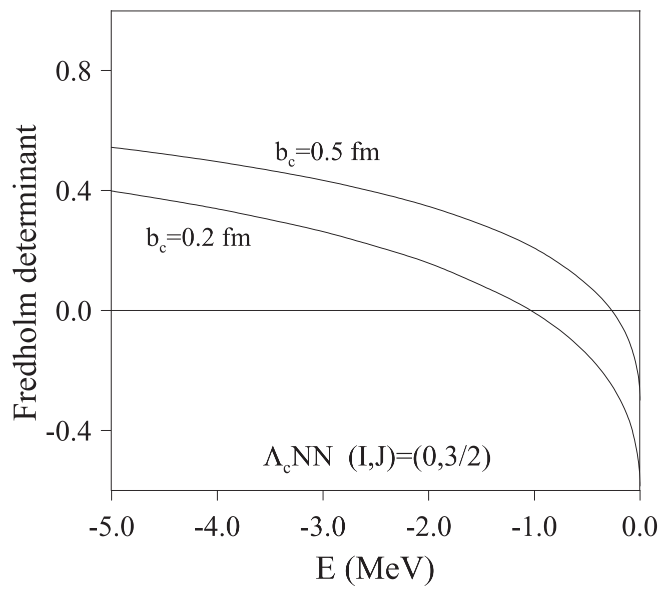

FIG. 5. Fredholm determinant for the $(I,J) = (0,3/2)$ $\Lambda_c NN$ channel for the different values of the width parameter for the charmed quark wave function.

due to their mass difference, 168 MeV as compared to the 73 MeV of the strange sector. Besides, it comes reduced with respect to the strange sector due to the absence of the strange meson exchanges [8], giving rise to a smaller $N\Lambda_c \leftrightarrow N\Sigma_c$ transition potential. One should also have in mind that in the $(I,J) = (0,3/2)$ channel the charmed baryon-nucleon interaction with spin singlet does not contribute, being much more repulsive than the spin-triplet one. Note that the ratio for the relative contribution of the spin-singlet to the spin-triplet partial waves comes determined from the strange sector through the Λp scattering cross section and the hypertriton binding energy [20].

We show in Fig. 3 the real part of the Fredholm determinant of the six (I,J) $\Sigma_c NN$ channels that are possible for energies below the $\Sigma_c d$ threshold. The imaginary part of the Fredholm determinant is small. As one can see all channels are repulsive and thus uninteresting from the point of view of possible bound states. The larger attraction is found in the $(I,J) = (1,1/2)$ $\Sigma_c NN$ channel that in the strange sector presented a quasibound state close to the three-body threshold [20]. Such ΣNN quasibound state has been recently suggested in ^{3}He$(K^-,\pi^\mp)$ reactions at 600 MeV/c [28].

To check the relevance of the $\Lambda_c \leftrightarrow \Sigma_c$ conversion we have solved the most interesting $\Lambda_c NN$ channels, $(I,J) = (0,1/2)$ and $(0,3/2)$, switching off the transition between the $\Lambda_c N$ and $\Sigma_c N$ subsystems. We plot in Fig. 4 the Fredholm determinant for both cases. The solid line indicates the result of the full calculation while the dashed one represents the results without $\Lambda_c \leftrightarrow \Sigma_c$ conversion. As can be seen the effect of the $\Sigma_c NN$ channel for $\Lambda_c NN$ is not so important as in the strange case. In particular, the order between the two channels is not reversed. As mentioned above, the reduction of the kinetic energy associated to the Λ_c^+ compared with that of the Λ implies that the $(I,J) = (0,3/2)$ state is more strongly bound than in the strange sector.

Two final remarks are in order. First, there are estimations on the literature about the binding energy of heavier charmed hypernuclei [8] and its dependence on the hard core of the $N\Lambda_c$ interaction. Our quark-model approach presents the advantage of having the hard-core radius fixed by means of quark-antisymmetrization. Second, we have calculated the contribution of the Coulomb potential exactly obtaining a contribution of $E_C = 131$ keV, which would give a final binding energy of 140 keV for the $J = 3/2$ charmed hypertriton. This result could be easily understood by considering the Λnp state as a bound state of a deuteron and a Λ. The Coulomb energy would be given at first order by

$$E_C = \frac{\langle \Psi(\vec{r}\,) | V(r) | \Psi(\vec{r}\,) \rangle}{\langle \Psi(\vec{r}\,) | \Psi(\vec{r}\,) \rangle} , \tag{13}$$

with $V(r) = \frac{\alpha}{r}$. For small binding energies B, the wave function can be represented by $\Psi(\vec{r}\,) = e^{-kr}$, with $k = \sqrt{2\eta B}$, where η is the reduced mass of the $\Lambda - d$ system and B is the binding energy. This would give rise to $E_C = \alpha\, k$, and using the binding energy B, we would obtain $E_C = 172$ keV, comparable to our exact result, and that would not be enough as to destroy the bound state.

To check the dependence of the binding energy of the $J = 3/2$ charmed hypertriton on the free parameter of the quark-model charmed baryon-nucleon interaction, we show in Fig. 5 the results of the $\Lambda_c NN$ $(I,J) = (0,3/2)$ state for several models with slightly different width parameters for the charmed quark wave function. One should have in mind that in Ref. [29] it was argued that the smaller values of b_c are preferred to get consistency with calculations based on infinite expansions, as hyperspherical harmonic expansions [30], where the quark wave function is not postulated. This also agrees with simple harmonic oscillator relations $b_c = b_n\sqrt{\frac{m_n}{m_c}}$. As can be seen the binding energy varies between 1037 and

271 keV, and the repulsive Coulomb contribution would vary between 322 and 131 keV, which makes the competition between electromagnetic and strong contributions crucial for the existence of this state, which would have a binding energy between 715 and 140 keV.

IV. SUMMARY

In summary, we have solved the Faddeev equations for the $\Lambda_c NN$ and $\Sigma_c NN$ systems using the charmed baryon-nucleon and nucleon-nucleon interactions derived from a chiral constituent quark model with full inclusion of the $\Lambda_c \leftrightarrow \Sigma_c$ conversion. We present results for the binding energy and the $\Lambda_c d$ and $\Sigma_c d$ scattering lengths. As compared to the strange sector, the kinetic energy is reduced but the interactions are weaker. The smaller contribution of the $\Lambda_c \leftrightarrow \Sigma_c$ conversion due to the larger mass difference and the smaller transition potential reverses the order of the two only attractive channels, $(I,J) = (0,1/2)$ and $(0,3/2)$, with the spin-3/2 state becoming the most attractive one. After correcting for Coulomb effects the charmed hypertriton would have a binding energy of at least 140 keV. The actual experimental facilities are capable of carrying experiments seeking for these states, which would help us in our progress in the knowledge of the baryon-baryon interaction on the heavy-flavor sector.

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