Strangeness -2 Hypertriton

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We solve for the first time, the Faddeev equations for the bound state problem of the coupled $\Lambda\Lambda N - \Xi NN$ system to study whether or not a hypertriton with strangeness -2 may exist. We make use of the interactions obtained from a chiral quark model describing the low-energy observables of the two-baryon systems with strangeness 0, -1, and -2 and three-baryon systems with strangeness 0 and -1. The $\Lambda\Lambda N$ system alone is unbound. However, when the full coupling to ΞNN is considered, the strangeness -2 three-baryon system with quantum numbers $(I, J^P) = (\frac{1}{2}, \frac{1}{2}^+)$ becomes bound, with a binding energy of about 0.5 MeV. This result is compatible with the nonexistence of a stable ${}^3_\Lambda$H with isospin one.

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The strangeness $\hat{S} = -2$ sector has become an important issue for theoretical and experimental studies of the strangeness nuclear physics. The $\Xi N - \Lambda\Lambda$ interaction accounts for the existence of doubly strange hypernuclei, which is a gateway to strange hadronic matter. The (K^-, K^+) reaction is one of the most promising ways of studying doubly strange systems. $\Lambda\Lambda$ hypernuclei can be produced through the reaction $K^- p \to K^+ \Xi^-$ followed by $\Xi^- p \to \Lambda\Lambda$. Strangeness -2 baryon-baryon interactions also account for a possible six-quark H dibaryon, which has yet to be experimentally observed. The future E07 experiment from J–PARC [1,2] is expected to improve our knowledge of the $\hat{S} = -2$ sector, giving ten times more emulsion events for double-Λ hypernuclei.

On the experimental side, there are very few data in the $\hat{S} = -2$ sector coming from the inelastic $\Xi^- p \to \Lambda\Lambda$ cross section at a lab momentum of around 500 MeV$/c$, and from the elastic $\Xi^- p \to \Xi^- p$ and inelastic $\Xi^- p \to \Xi^0 n$ cross sections for lab momenta in the range of 500–600 MeV$/c$ [3–5]. The relevant information we have is indirect and comes from double-Λ hypernuclei. Their binding energies, $B_{\Lambda\Lambda}$, provide upper limits for that of the H dibaryon, i.e., $B_H < B_{\Lambda\Lambda}$. The first hypernuclear events are quite old and admit several interpretations [6–8]. In 2001 the so–called Nagara event was reported [9], interpreted uniquely as the sequential decay of ${}^{6}_{\Lambda\Lambda}$He emitted from a Ξ^--hyperon nuclear capture at rest. The mass and the values of $B_{\Lambda\Lambda}$ and of the $\Lambda\Lambda$ interaction energy, $\Delta B_{\Lambda\Lambda}$, were determined without ambiguities. It gave the most stringent constraint to the mass of the H dibaryon to date, i.e., $M_H > 2223.7$ MeV at a 90% confidence level. It took almost one decade, but four more double-Λ hypernuclear events were reported, from KEK E176 and E373 experiments [1], still with preliminary results. All the details are summarized in Table I.

Besides the double-Λ hypernuclei quoted in Table I, there is a general consensus that the mirror $\Lambda\Lambda$ hypernuclei ${}^{5}_{\Lambda\Lambda}$H-${}^{5}_{\Lambda\Lambda}$He are particle stable [10]. The existence of a ${}^{4}_{\Lambda\Lambda}$H bound state has been claimed by the AGS-E906 experiment [11], from correlated weak-decay pions emitted sequentially by $\Lambda\Lambda$ hypernuclei produced in a (K^-, K^+) reaction on ^{9}Be. The stability of the $\Lambda\Lambda N$ system was discarded long ago [12] by using symmetry considerations with respect to the ${}^{3}_{\Lambda}$H, and therefore without considering the important coupled channel effect due to the existence of the ΞNN system.

Theoretically, the $\hat{S} = -2$ sector was recently put back on the agenda by lattice QCD calculations of different collaborations, NPLQCD [13] and HAL QCD [14], providing evidence for a $\Lambda\Lambda$ bound state for nonphysical values of the pion mass ($m_\pi = 389$ MeV and $m_\pi = 673 \to 1010$ MeV, respectively). When performing quadratic and linear extrapolations to the physical point [15], a bound dibaryon (around 7 MeV) and a H at threshold, respectively, are predicted. Reference [15] presents preliminary results for $m_\pi = 230$ MeV, much closer to the physical pion mass, pointing to a H dibaryon at threshold, as also experimentally suggested by the enhancement of the $\Lambda\Lambda$ production near threshold found in Ref. [16].

The purpose of this Letter is twofold. On the one hand we present the solution of the Faddeev equations for the

TABLE I. Double Λ hypernuclear events.

Event	Nuclide	$B_{\Lambda\Lambda}$ (MeV)	$\Delta B_{\Lambda\Lambda}$ (MeV)
1963	${}^{10}_{\Lambda\Lambda}$Be	17.7 ± 0.4	4.3 ± 0.4
1966	${}^{6}_{\Lambda\Lambda}$He	10.9 ± 0.5	4.7 ± 1.0
1991	${}^{13}_{\Lambda\Lambda}$B	27.5 ± 0.7	4.8 ± 0.7
NAGARA	${}^{6}_{\Lambda\Lambda}$He	6.91 ± 0.16	0.67 ± 0.17
MIKAGE	${}^{6}_{\Lambda\Lambda}$He	10.06 ± 1.72	3.82 ± 1.72
DEMACHIYANAGI	${}^{10}_{\Lambda\Lambda}$Be	11.90 ± 0.13	-1.52 ± 0.15
HIDA	${}^{11}_{\Lambda\Lambda}$Be	20.49 ± 1.15	2.27 ± 1.23
	${}^{12}_{\Lambda\Lambda}$Be	22.23 ± 1.15	...
E176	${}^{13}_{\Lambda\Lambda}$Be	23.3 ± 0.7	0.6 ± 0.8

bound state problem of the coupled $\Lambda\Lambda N$-ΞNN system. The system has been formally studied and its Faddeev equations written down [17,18], although they have never been applied in a numerical calculation with realistic two-body interactions. This is basically due to the fact that one requires a model of the baryon-baryon interaction which should be able to simultaneously describe two-baryon states with strangeness 0, -1, and -2 within a single consistent theoretical framework. Afterwards, we will apply the formalism by means of the interactions obtained from a chiral quark model describing the low-energy observables of the two-baryon systems with strangeness 0, -1, and -2 and also three-baryon systems with strangeness 0 and -1, trying to elucidate the nature of the three-baryon system with strangeness -2.

The coupled $\Lambda\Lambda N$-ΞNN system is peculiar because it has two identical particles in each of its two components although they are of different type, which complicates considerably its analysis. The Faddeev equations for the bound-state problem of the coupled $\Lambda\Lambda N$-ΞNN system have been derived in Ref. [18]. We have obtained these same equations by an independent method [19], they read:

$$
\begin{aligned}
T^{\Xi NN} &= t^{NN,NN}(1-P_{23})P_{13}P_{23}G_0^{N\Xi N}T^{N\Xi N},\\
T^{N\Xi N} &= t^{N\Xi,N\Xi}P_{12}P_{23}G_0^{\Xi NN}T^{\Xi NN} - t^{N\Xi,N\Xi}P_{13}G_0^{N\Xi N}T^{N\Xi N} + t^{N\Xi,\Lambda\Lambda}(1-P_{23})P_{13}P_{23}G_0^{\Lambda N\Lambda}T^{\Lambda N\Lambda},\\
T^{N\Lambda\Lambda} &= t^{\Lambda\Lambda,N\Xi}P_{12}P_{23}G_0^{\Xi NN}T^{\Xi NN} - t^{\Lambda\Lambda,N\Xi}P_{13}G_0^{N\Xi N}T^{N\Xi N} + t^{\Lambda\Lambda,\Lambda\Lambda}(1-P_{23})P_{13}P_{23}G_0^{\Lambda N\Lambda}T^{\Lambda N\Lambda},\\
T^{\Lambda N\Lambda} &= t^{N\Lambda,N\Lambda}P_{12}P_{23}G_0^{N\Lambda\Lambda}T^{N\Lambda\Lambda} - t^{N\Lambda,N\Lambda}P_{13}G_0^{\Lambda N\Lambda}T^{\Lambda N\Lambda},
\end{aligned}
\tag{1}
$$

where G_0^{ijk} is the propagator for three free particles ijk, $t^{ij,kl}$ are the two-body t matrices for the different transitions $ij \to kl$, and P_{ij} is the exchange operator for particles i and j. The first superscript in the T functions is the spectator and the other two are the interacting pair. We will solve these equations including all the S-wave configurations $\ell_i = \lambda_i = 0$, where ℓ_i is the orbital angular momentum between particles j and k, and λ_i is the orbital angular momentum between particle i and the pair jk. Therefore, the total angular momentum $J = 1/2$ is equal to the total spin.

The set of Eqs. (1) are integral equations in two continuous variables which couple the nine two-body channels obtained from Table II. In order to solve these equations we use the method applied in our previous works [20,21], where the two-body t matrices are expanded in terms of Legendre polynomials leading to integral equations in only one continuous variable coupling the various Legendre components required for convergence.

In each of the two components of the coupled $\Lambda\Lambda N$-ΞNN system we take particles 2 and 3 to be the two identical ones and particle 1 to be the different one. We will take the basis states 1 and 3 using a cyclic coupling scheme, i.e., $1 = (2+3)+1$, and $3 = (1+2)+3$, while for the basis state 2 we use the anticyclic scheme $2 = (1+3)+2$. With these conventions, Eqs. (1) take the explicit form [19],

$$
\begin{aligned}
T^{\Xi NN}_{\alpha_1 m}(q_1) &= 2\sum_{\alpha_3 n}\int_0^\infty q_3^2 dq_3 K^{NN,NN;N\Xi N}_{mn;\alpha_1\alpha_3;13}(q_1 q_3)T^{N\Xi N}_{\alpha_3 n}(q_3),\\
T^{N\Xi N}_{\alpha_3 m}(q_3) &= \sum_{\alpha_1 n}\int_0^\infty q_1^2 dq_1 K^{N\Xi,N\Xi;\Xi NN}_{mn;\alpha_3\alpha_1;31}(q_3 q_1)T^{\Xi NN}_{\alpha_1 n}(q_1) - \sum_{\alpha_3' n}\int_0^\infty q_3'^2 dq_3' K^{N\Xi,N\Xi;N\Xi N}_{mn;\alpha_3\alpha_3';23}(q_3 q_3')T^{N\Xi N}_{\alpha_3' n}(q_3')\\
&\quad + 2\sum_{\alpha_3' n}\int_0^\infty q_3'^2 dq_3' K^{N\Xi,\Lambda\Lambda;\Lambda N\Lambda}_{mn;\alpha_3\alpha_3';13}(q_3 q_3')T^{\Lambda N\Lambda}_{\alpha_3' n}(q_3'),\\
T^{N\Lambda\Lambda}_{\alpha_1 m}(q_1) &= \sum_{\alpha_1' n}\int_0^\infty q_1'^2 dq_1' K^{\Lambda\Lambda,N\Xi;\Xi NN}_{mn;\alpha_1\alpha_1';31}(q_1 q_1')T^{\Xi NN}_{\alpha_1' n}(q_1') - \sum_{\alpha_3 n}\int_0^\infty q_3^2 dq_3 K^{\Lambda\Lambda,N\Xi;N\Xi N}_{mn;\alpha_1\alpha_3;23}(q_1 q_3)T^{N\Xi N}_{\alpha_3 n}(q_3)\\
&\quad + 2\sum_{\alpha_3 n}\int_0^\infty q_3^2 dq_3 K^{\Lambda\Lambda,\Lambda\Lambda;\Lambda N\Lambda}_{mn;\alpha_1\alpha_3;13}(q_1 q_3)T^{\Lambda N\Lambda}_{\alpha_3 n}(q_3),\\
T^{\Lambda N\Lambda}_{\alpha_3 m}(q_3) &= \sum_{\alpha_1 n}\int_0^\infty q_1^2 dq_1 K^{N\Lambda,N\Lambda;N\Lambda\Lambda}_{mn;\alpha_3\alpha_1;31}(q_3 q_1)T^{N\Lambda\Lambda}_{\alpha_1 n}(q_1) - \sum_{\alpha_3' n}\int_0^\infty q_3'^2 dq_3' K^{N\Lambda,N\Lambda;\Lambda N\Lambda}_{mn;\alpha_3\alpha_3';23}(q_3 q_3')T^{\Lambda N\Lambda}_{\alpha_3' n}(q_3'),
\end{aligned}
\tag{2}
$$

where

$$K^{\beta\gamma,\delta\zeta;\kappa\lambda\rho}_{mn;\alpha_i\alpha_j;kl}(q_iq_j) = \frac{2m+1}{4}A^{\alpha_i\alpha_j}_{kl}\int_{-1}^{1}d\cos\theta\int_{-1}^{1}dx_iP_m(x_i)P_n(x_j)t^{\beta\gamma,\delta\zeta}_{\alpha_i}(p_i,p'_i;E+\Delta E-q_i^2/2\nu_i) \times \frac{1}{E+\Delta E-p'^2_j/2\eta_j-q_j^2/2\nu_j}. \quad (3)$$

η_j and ν_j are the usual reduced masses and $P_n(x)$ is a Legendre polynomial. $p_i = b(1+x_i)/(1-x_i)$, $x_j = (p_j - b)/(p_j + b)$, and b is a scale parameter on which the solution does not depend. $p'_i = [q_j^2 + (\eta_i q_i/m_k)^2 + 2(\eta_i q_i q_j/m_k)\cos\theta]^{1/2}$, $p_j = [q_i^2 + (\eta_j q_j/m_k)^2 + 2(\eta_j q_i q_j/m_k)\cos\theta]^{1/2}$, and $\Delta E = 0$ if the corresponding state (either i or j) belongs to the $N\Lambda\Lambda$ component, while $\Delta E = 2m_\Lambda - m_N - m_\Xi$ if the corresponding state belongs to the ΞNN component. Finally, $A^{\alpha_i\alpha_j}_{kl}$ are the usual spin-isospin transition coefficients [20], where $\delta\zeta$ is the interacting pair in the state i and $\lambda\rho$ is the interacting pair in the state j.

For practical purposes, we took into account all the S-wave two-body amplitudes that contribute in Eqs. (2) as shown in Table II. Even though our calculation will include only two-body S waves, the corresponding two-body amplitudes will be obtained from a full model, including D waves in spin-triplet channels and the coupling to higher mass states in those cases where the quantum numbers allow for it.

Once the method to solve the bound state problem of the $\Lambda\Lambda N$-ΞNN system has been designed, we apply it to the chiral quark model of the baryon-baryon interaction developed in Ref. [22]. The model is capable of describing the low-energy parameters of the two-nucleon system, the S-wave phase shifts, and the triton binding energy [23]. It reproduces the elastic and inelastic scattering cross sections of the $\hat{S} = -1$ two-baryon systems and the hypertriton binding energy [20,21]. As can be seen in Fig. 2 of Ref. [21], the isospin one ΛNN system is unbound. Finally, the model provides parameter-free predictions for the elastic and inelastic scattering cross sections of the $\hat{S} = -2$ two-baryon systems [24] that are consistent with the scarce available data. In particular, the relevant $\Xi^- p \to \Lambda\Lambda$ is correctly described (see Fig. 2 of Ref. [24]). Thus, we are confident that the interactions are realistic enough to allow for the study of the existence (or nonexistence) of the strangeness -2 hypertriton.

TABLE II. S-wave two-body channels $(i,\ j)$ of the various subsystems that contribute to the strangeness -2 $(I, J^P) = (\frac{1}{2}, \frac{1}{2}^+)$ three-body state.

Subsystem	$(i,\ j)$ channels
NN	(0,1),(1,0)
$N\Lambda$	$(\frac{1}{2}, 0)$, $(\frac{1}{2}, 1)$
$\Lambda\Lambda$	(0,0)
$N\Xi$	(0,0),(0,1),(1,0),(1,1)

The H dibaryon has strangeness -2, positive parity, and isospin and spin $(i, j) = (0, 0)$. It appears in our model as a bound state of the coupled $\Lambda\Lambda$-$N\Xi$-$\Sigma\Sigma$ system with a binding energy of 6.928 MeV [24]. Therefore, the main configuration of the strangeness -2 hypertriton will be a H dibaryon as the interacting pair and a S-wave nucleon as spectator, which leads to total isospin and spin $(I, J) = (\frac{1}{2}, \frac{1}{2})$ and positive parity. This configuration is also favored by having a deuteron as interacting pair and a S-wave Ξ hyperon as spectator. We give in Table II all the S-wave two-body channels that contribute to the $(I, J^P) = (\frac{1}{2}, \frac{1}{2}^+)$ three-body state. The NN channels have, of course, strangeness 0, the $N\Lambda$ channels have strangeness -1, and the $\Lambda\Lambda$ and $N\Xi$ channels both have strangeness -2. As can be seen from this table, the $\Lambda\Lambda N$ and ΞNN systems are coupled together through the $(i, j) = (0, 0)$ two-body channel.

In Ref. [21] we calculated the hyperon-deuteron (Yd) scattering lengths as well as the hypertriton binding energy taking into account all the S- and D-wave components that contribute in the various three-body channels. From a combined analysis of the nucleon-hyperon (NY) two-body data, the Yd scattering lengths, and the hypertriton binding energy, we were able to constrain the allowed values of the ΛN spin-triplet and spin-singlet scattering lengths as $1.41 \le a^{\Lambda N}_{1/2,1} \le 1.58$ fm, and $2.33 \le a^{\Lambda N}_{1/2,0} \le 2.48$ fm. Therefore, we now make use of the NY interacting models satisfying these constraints to calculate the binding energy of the strangeness -2 hypertriton. The results obtained in the full coupled channel problem $\Lambda\Lambda N$-ΞNN are shown in Table III. As one can see from this table, the strangeness -2 three-baryon system with quantum numbers $(I, J^P) = (\frac{1}{2}, \frac{1}{2}^+)$ is bound, the binding energy varying between 0.4 and 0.6 MeV. However, as predicted in Ref. [12] due to the nonexistence of an isospin one ${}^{3}_{\Lambda}$H bound state, the $\Lambda\Lambda N$ system alone is not bound. The bound state only appears when the coupling between

TABLE III. Binding energy of the strangeness -2 hypertriton (in MeV) measured with respect to the NH threshold for several models of the NY interaction satisfying the constraints of Ref. [21] for the $N\Lambda$ scattering lengths, $a^{N\Lambda}_{i,j}$ (in fm).

		$a^{N\Lambda}_{1/2,1}$			
		1.41	1.46	1.52	1.58
	2.33	0.416	0.455	0.495	0.542
$a^{N\Lambda}_{1/2,0}$	2.39	0.424	0.463	0.504	0.551
	2.48	0.447	0.487	0.528	0.577

the $\Lambda\Lambda N$ and ΞNN components is considered, i.e., when the $(i, j) = (0, 0)$ two-body $t^{\Lambda\Lambda,N\Xi}$ amplitude is included in the calculation.

The relevance of the $\Lambda\Lambda$-ΞN coupling for double-Λ hypernuclei has been emphasized for the case of the ${}^{4}_{\Lambda\Lambda}$H hypernucleus [25,26]. If this system is studied with NN, $N\Lambda$, and $\Lambda\Lambda$ interactions improved for the description of the ${}^{6}_{\Lambda\Lambda}$He, it is found to be unbound. In the case of the ${}^{6}_{\Lambda\Lambda}$He the $\Lambda\Lambda$-$N\Xi$ coupling plays a minor role, because the nucleon generated in the transition must occupy an excited p shell, the lowest s shell being forbidden by the Pauli principle. This is not the case for the ${}^{4}_{\Lambda\Lambda}$H, where the nucleon generated by the $\Lambda\Lambda$-$N\Xi$ transition can occupy a hole in the lowest s shell. This effect generates theoretical binding for the ${}^{4}_{\Lambda\Lambda}$H [25] and it is also responsible for generating binding in the strangeness -2 three-baryon system with quantum numbers $(I, J^P) = (\frac{1}{2}, \frac{1}{2}^+)$. It is therefore important to obtain experimental information about the strength of the $\Lambda\Lambda$-ΞN coupling. It could be derived from the measurement of the ${}^{4}_{\Lambda\Lambda}$H binding energy. In the meantime, the only available experimental data are the inelastic cross section $\Xi^- p \rightarrow \Lambda\Lambda$, correctly described by the present model (see Fig. 2 of Ref. [24]).

The possible existence of a strangeness -2 hypertriton will give a strong impact on forthcoming experimental projects as well as ongoing theoretical studies. Experimentally, it could be measured in the J-PARC-E07 experiment, where more than 10^3 $\Lambda\Lambda$ nuclei are expected to be detected by means of Ξ-capture reactions using different target nuclei: C, N, and O [27]. Theoretically, lattice QCD has evolved to the point where the calculation of the binding energy of light nuclei and hypernuclei with $A \leq 4$ and $\hat{S} \leq 2$, at unphysically heavy light-quark masses, is possible [28]. Extrapolations to the physical light-quark masses have not been attempted because the quark mass dependences of the energy levels in the light nuclei are not known. Future calculations at smaller lattice spacings and at lighter quark masses will facilitate such extrapolations and, therefore, comparison with experiment and, thus, the analysis of the strangeness -2 hypertriton.

In summary, we have solved for the first time the Faddeev equations for the bound state problem of the coupled $\Lambda\Lambda N$-ΞNN system to study whether or not a hypertriton with strangeness -2 may exist. We make use of the interactions obtained from a chiral quark model describing the low-energy observables of the two-baryon systems with strangeness 0, -1, and -2 and three-baryon systems with strangeness 0 and -1. The $\Lambda\Lambda N$ system alone is unbound in agreement with the nonexistence of an isospin one ${}^{3}_{\Lambda}$H bound state. However, when the full coupling to ΞNN is considered through the $(i, j) = (0, 0)$ two-body $t^{\Lambda\Lambda,N\Xi}$ amplitude, the strangeness -2 three-baryon system with quantum numbers $(I, J^P) = (\frac{1}{2}, \frac{1}{2}^+)$ becomes bound, with a binding energy of about 0.5 MeV.

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