The continuum structure of the Borromean halo nucleus \( ^{11}\text{Li} \)

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Abstract

We solve the Faddeev equations for \( ^{11}\text{Li} \) (\( n+n+^{9}\text{Li} \)) using hyperspherical coordinates and analytical expressions for distances much larger than the effective ranges of the interactions. The lowest resonances are found at 0.65 MeV \( (\frac{1}{2}^{+}, \frac{3}{2}^{+}, \frac{5}{2}^{+}) \) and 0.89 MeV \( (\frac{3}{2}^{+}) \) with widths of about 0.35 MeV. A number of higher-lying broader resonances are also obtained and related to the Efimov effect. The dipole strength function and the Coulomb dissociation cross section are also calculated.

PACS numbers: 21.45.+v, 11.80.Jy, 21.60.Gx

Introduction. The new class of nuclear states, called halos, are presently intensively investigated \[1\]. Their structure and behavior in reactions differ substantially from those of ordinary nuclei. They are unusually large and loosely bound. Most of them are presumably not identified yet. Large electromagnetic dissociation cross section, or a concentration of \( 1^{-} \) strength function at low energies exemplify the unusual character of these nuclei. Descriptions as two- and three-body systems have been successful \[2, 3\]. The continuum structure of halo nuclei is still almost unknown even for the most studied halo nucleus \( ^{11}\text{Li} \). However almost inevitably low-lying resonances must be present, since the Borromean nature implies that the two-body subsystems are close to threshold and perhaps nearly fulfill the Efimov condition \[4, 5\].

A recent experiment, where protons are scattered on \( ^{11}\text{Li} \), was interpreted as evidence for a \( 1^{-} \)-resonance in \( ^{11}\text{Li} \) at 1.0 ± 0.1 MeV above the two-neutron threshold with a width of about 0.75 ±0.6 MeV \[6\]. An other interpretation of the same experiment is given in terms of the shake-off process \[7\]. Reaction experiments also indicate an excited \( ^{11}\text{Li} \)-state at about the same energy \[8, 9\]. The measured \[6, 8, 10\] and computed \[11, 12, 13\] \( 1^{-} \)-strength functions are available in the literature, where the absence of calculated \( 1^{-} \) resonances is quoted without details.

Clearly reliable three-body continuum calculations are desirable for \( ^{11}\text{Li} \), but this is not easy due to both technical difficulties and lack of information of \( ^{10}\text{Li} \).
However, recently a method dealing with the necessary large distances was formulated for s-waves \cite{5, 14} and developed and applied for three α-particles and two neutrons surrounding a spin zero core \cite{15}. The purpose of the present letter is to extend the method to finite core-spin and provide continuum structure calculations for $^{11}$Li.

**Method.** The method is described for bound states in \cite{2, 5, 14} and for continuum states in \cite{15}. We shall here briefly sketch the necessary generalizations.

The $k$'th particle has mass $m_k$, charge $eZ_k$, spin $s_k$ and coordinate $r_k$. The two-body potentials are $V_{ij}$. The spins $s_j$ and $s_k$ are coupled to $s$ which in turn coupled to $s_i$ gives the total spin. The corresponding wave function is $\chi_s(i)$. The total wavefunction is expanded in a complete set of hyperangular functions

$$\Psi(\rho, \Omega) = \frac{1}{\rho^{5/2}} \sum_{n=1}^{\infty} f_n(\rho) \phi_n(i) \sin(2\alpha_i)$$

where each of the three components $\phi_n(i)$ is expressed in the corresponding set of hyperspherical coordinates $(\rho, \Omega_i) = (\rho_i, \alpha_i, \Omega_{xi}, \Omega_{yi})$. They satisfy the angular part of the Faddeev equations with the eigenvalues $\lambda_n(\rho)$.

For Borromean systems we select those solutions $\Psi_n$ to eq.(1) where the large-distance ($\rho \to \infty$) boundary conditions for $f_n$ are given by \cite{16}

$$f_n'(-\rho) \to \delta_{n,n'} F_n(-\kappa \rho) - S_{n,n'} F_n(\kappa \rho)$$

where $\kappa^2 = 2mE/\hbar^2$ and $F_n(\pm)$ are related to the Hankel functions of integer order

$$F_n(\pm)(\kappa \rho) = \sqrt{\frac{m\rho}{4\hbar^2}} H_n^{(\pm)}(\kappa \rho)$$

where $K_n(K_n + 4) = \lambda_n(\rho = \infty)$. The continuum wave functions are orthogonal and normalized to delta functions in energy.

By diagonalization of the $S$-matrix we obtain eigenfunctions and eigenphases. Resonance energies and widths can be obtained by use of the complex energy method, where the Faddeev equations are solved for $E = E_r - i\Gamma/2$ with the boundary condition $f_n' \propto \delta_{n,n'} \sqrt{\frac{m\rho}{4\hbar^2}} H_n^{(\pm)}(\kappa \rho)$ These solutions correspond to poles of the $S$-matrix \cite{16}.

**Large-distance behavior.** We shall now specialize to the system of interest, i.e. two neutrons ($k = 2, 3$) and a core ($k = 3$) with spin $s_c$. The three independent orbital s-wave components, $\phi_{L,s_i}^{(i)}$, $i = 1, 2, 3$, are characterized by $l_x = 0$, $l_y = L$ and the spin wave functions $\chi_s^{(1)}$, $\chi_s^{(2)}$, $\chi_s^{(3)}$.
All other partial waves decouple for large $\rho$ over a distance scale defined by the range of the interactions. The s-waves, however, are coupled and feel the interactions over a distance defined by the scattering lengths. The remaining two s-wave components are determined by antisymmetry, i.e. $\phi^{(2)}_{L,s} = -\phi^{(3)}_{L,s}$, $\phi^{(3)}_{L,s_2} = \phi^{(2)}_{L,s_2}$. The last combination, $\phi^{(1)}_{L,s=1}$, is not allowed due to the Pauli principle.

The three components, $\phi^{(i)}_{L,s}$, obey for large $\rho$ the coupled angular Faddeev equations which, to leading order in $\alpha$, explicitly are given by

$$\left( \frac{\partial^2}{\partial \alpha^2} + \kappa_i^2(\alpha) \right) \phi^{(i)}_{L,s}(\rho, \alpha) = 2\alpha(-1)^L C^{(i)} \rho^2 v_i(\rho \sin \alpha) ,$$

$$\kappa^2(\alpha) = -\left[ \frac{L(L+1)}{\cos^2 \alpha} - \rho^2 v_i(\rho \sin \alpha) - v_i^2 \right] ,$$

where we defined $v_i^2 = \lambda_i + 4$, the s-wave interactions $v_i(x_i) = V_{jk}(x_i/\mu_{jk})^{2\mu_{jk}}$, $m\mu_{jk}^2 = m_j m_k/(m_j + m_k)$ and the coupling terms

$$C^{(i)} = \sum_j \sum_{s} C^{ij}_{ss'} \phi^{(j)}(\rho, \varphi_k) \sin(2\varphi_k) , \quad C^{ij}_{ss'} = \langle \chi^{(i)}_s | \chi^{(j)}_s \rangle ,$$

where $\varphi_k = \arctan(m_k(m_i+m_j+m_k)/m_i m_j)$. The s-summation runs over the two possible intermediate spin couplings for the component $j$. The spin overlap matrix elements are diagonal for $i = j$, i.e. $C^{ii}_{ss'} = \delta_{ss'}$ and symmetric, i.e. $C^{ij}_{ss'} = C^{ji}_{ss'}$.

The rescaled potentials $\rho^2 v_i(\rho \sin \alpha_i)$ approach for sufficiently large $\rho$ the zero-range potentials, where the sensitivity to the shape disappears. Any potential with the same scattering length and effective range would then lead to results accurate to the order $\rho^{-2}$. We shall therefore for convenience use square well potentials [14], where $v_i(\rho \sin \alpha_i) = v_i^{(0)}(\text{region I})$ when $\alpha_i < \alpha_i^{(0)} = \arcsin(R_i \mu_{jk}/\rho) \ll 1$ and zero otherwise (region II).

The solutions to eq.(4) are then for $\alpha_i > \alpha_i^{(0)}$ given by

$$\phi^{(i,11)}_{L,s_i}(\rho, \alpha_i) = A^{(i)}_{L} P_L(\nu, \alpha_i) ,$$

$$P_L(\nu, \alpha) \equiv \cos^L \alpha \left( \frac{\partial}{\partial \alpha \cos \alpha} \right)^L \sin \left[ \nu \left( \alpha - \frac{\pi}{2} \right) \right]$$

and for $\alpha_i < \alpha_i^{(0)}$ to leading order in $\alpha_i$ given by

$$\phi^{(i,1)}_{L,s_i}(\rho, \alpha) = B^{(i)}_{L} \sin(\kappa_i(0)\alpha) - 2\alpha(-1)^L \rho^{2\mu_{ij}^{(i)}} \kappa_i^2(0) C^{(i)}$$

where the wave functions in $C^{(i)}_L$ in eq.(6) must be $\phi^{(i,11)}_{L,s_i}$.
Matching the solutions, eqs. (7) and (9), and their derivatives at \( \alpha_i = \alpha_i^{(i)} \) gives a linear set of equations for \( A_L^{(i)} \) and \( B_L^{(i)} \). Large-distance physical solutions are obtained when the corresponding determinant is zero. This is the quantization condition for \( \lambda \) and the eigenvalue equation determining the asymptotic behavior of \( \lambda(\rho) \).

**Input parameters.** We use the neutron-neutron interaction which reproduces the low-energy properties of the nucleon-nucleon system \(^3\). The neutron-core, n–\(^9\)Li, effective interaction assumes that the spin of both \(^9\)Li and \(^{11}\)Li is \( \frac{3}{2} \). We use the mean-field spin-orbit term \( \mathbf{l} \cdot \mathbf{s} \) and we include a spin-spin term to differentiate between the two spin couplings for a given orbital angular momentum \( l \). Such a spin splitting term is hard to avoid due to the strong spin dependence of the underlying basic interaction. Several phase equivalent parametrizations are possible. They differ in the number of two-body bound states occupied by the core neutrons which therefore subsequently must be excluded in the computation \(^3\). The results are very close and we shall therefore here only use the shallow potentials without bound states.

We use the radial form, \( \exp(- (r/2.55\text{fm})^2) \), for all neutron-core potentials. The strengths \( V_s(lj_n) \) depend on orbital \( l \) and total neutron angular momentum \( j_n \) for zero core spin, and for finite core spin also on the total spin \( s = 1, 2 \). All possible s- and p-waves are included whereas other waves can be ignored to the accuracy we need. We use \( V_1(s_{1/2}) = -6.89 \text{ MeV}, V_2(s_{1/2}) = -7.51 \text{ MeV}, V_1(p_{1/2}) = -38.59 \text{ MeV}, V_2(p_{1/2}) = -35.65 \text{ MeV}, V_1(p_{3/2}) = 43.91 \text{ MeV} \) and \( V_2(p_{3/2}) = 46.85 \text{ MeV} \) corresponding to virtual s-states at 0.247 MeV and 0.140 MeV and \( p_{1/2} \)-resonances at 0.75 MeV and 1.60 MeV. The \( p_{3/2} \)-potential is purely repulsive thereby simulating the Pauli blocking by the core neutrons.

The choice of these parameters is dictated by the accumulated information about the structure of \(^{10}\)Li, i.e. a \( p \)-resonance at about 0.6 MeV, a low-lying virtual s-state and a small spin splitting of these states \(^3\). We also demand that the \(^{11}\)Li-binding energy is reproduced; we obtain 305 keV with the corresponding root mean square radius of 3.34 fm. The calculated fragment momentum distributions in break-up reactions also compare rather well with measured values suggesting about 20% of the \( p^2 \)-configuration in the \(^{11}\)Li ground state wave function \(^3\). The remaining part is in the \( s^2 \)-configuration. Limited ambiguity, although some, is then left for the parameters of these “realistic” interactions.

The number of Jacobi polynomials in the basis expansion is carefully chosen to give accurate numerical results up to a distance, typically around 40 fm, where the asymptotic behavior is reached and from then on the asymptotic solutions eqs. (7)-(9) are used. The accurate low-energy continuum spectrum calculations require integration of the radial equations up to distances of the order of five times the sum of the scattering lengths. For the n+n+\(^9\)Li system this is about 200 fm.
The key quantities are the angular eigenvalues \( \lambda_n(\rho) \), which, divided by \( \rho^2 \), are the adiabatic effective radial potentials. They depend on total angular momentum and parity as seen in Fig. 1, where we show the spectra for \( \frac{3}{2}^- \) (ground state) and \( \frac{3}{2}^+ \) (1\(^-\)-excitation). The structure is complicated at small distances where avoided level crossings are seen. The lowest level has in both cases an attractive pocket, which is responsible for the ground state (\( \frac{3}{2}^- \)) and several resonances. At large distance the structure is simpler as the hyperspherical spectrum is approached. In the computation we use the asymptotic behavior (shown in Fig. 1) of the three coupled s-waves, one from each Jacobi set.

The \(^{11}\)Li-states can be labeled as arising from couplings of the \( s_{1/2} \) and \( p_{1/2} \) neutron-core states to angular momentum and parity \( 0^\pm \) and \( 1^\pm \), which in turn is coupled with the core spin \( 3/2^- \). The result for each parity \( \pi = \pm 1 \) is three nearly degenerate states of \( 1/2^\pi \), \( 3/2^\pi \), \( 5/2^\pi \) and one non-degenerate \( 3/2^\pi \)-state. This structure is reflected in the eigenvalue spectra where \( 3/2^\pm \) contain additional levels on top of spectra similar to those of \( 1/2^\pm \) and \( 5/2^\pm \). The degeneracy would be complete without spin splitting of the neutron-core potentials.
Phase shifts, poles and the Efimov effect. We show in Fig. 2 the eigenphases for the cases in Fig. 1. For $\frac{3}{2}^-$, we find rapid variations at smaller energy and one crossing of $\pi/2$ at about 1.2 MeV. For $\frac{1}{2}^-$ and $\frac{5}{2}^-$, none of them shown, we find rather smooth phase shifts indicating that the structure in $\frac{3}{2}^-$ is due to the non-degenerate components. For $\frac{3}{2}^+$, four crossings of $\pi/2$ are seen in two pairs. In $\frac{1}{2}^+$ and $\frac{5}{2}^+$, is found a similar structure, but one member of each pair is not present indicating the non-degenerate nature.

For the lowest spins and parities we give in Table 1 the lowest S-matrix poles obtained by the complex energy method. We see the three times nearly degenerate $1^-\text{-resonance}$ at about 0.65 MeV with a width of about 0.35 MeV. We also see degenerate $0^\pm\text{-resonances}$ at 0.89 MeV with widths of 0.33 MeV and 0.43 MeV. In addition degenerate $1^\pm\text{-poles}$ appear at 1.3 - 1.4 MeV, 1.6 - 1.8 MeV and 1.9 - 2.1 MeV with widths of about 0.5 MeV, 0.6 MeV and 0.8 MeV, respectively. Also poles of $0^\pm\text{-type}$ should be present for $3/2^\pm$ at these energies as seen by comparing with the model of zero core spin. They can not be distinguished numerically.

The height of the effective radial barriers are 1.7 MeV, 0.9 MeV, 0.7 MeV and 0.6 MeV for $\frac{1}{2}^-$, $\frac{1}{2}^+$, $\frac{5}{2}^-$, $\frac{5}{2}^+$, respectively. The pocket and the barrier in this potential is absent for both $\frac{1}{2}^-$ and $\frac{5}{2}^-$. These potentials are all attractive when the centrifugal barrier is removed. The lowest resonances appear around the barrier and their widths are consequently relatively large and rather sensitive to fine tuning of the interactions. Related cross sections would probably
Table 1: The real and imaginary values \((E_r, \Gamma)\) (in MeV) of the lowest \(S\)-matrix poles \(E = E_r - i\Gamma/2\) for \(^{11}\text{Li}\) for various spins and parities \(J^\pi\). The excitation energy \(E^* = E_r + 0.305\) MeV. The interactions for the upper part of the table are the same as in Fig. 1. The lower part of the table are results for a model with \(s_c = 0\) and the same average positions of the neutron-core resonances, i.e. energies of the \(s_{1/2}\) virtual state and the \(p_{1/2}\) resonance at 0.2 MeV and 1.22 MeV, respectively.

| \(J^\pi\)  | \(E_r\)  | \(\Gamma\) | \(E_r\)  | \(\Gamma\) | \(E_r\)  | \(\Gamma\) | \(E_r\)  | \(\Gamma\) |
|-----------|----------|------------|----------|------------|----------|------------|----------|------------|
| \(2^+\)  | -        | -          | 1.37     | 0.51       | 1.56     | 0.56       | 1.98     | 0.65       |
| \(-2^-\) | -0.305   | 0.89       | 0.43     | 1.41       | 0.56     | 1.60       | 0.61     | 2.03       |
| \(1^+\)  | -        | -          | 1.36     | 0.49       | 1.60     | 0.68       | 2.01     | 0.72       |
| \(2^+\)  | 0.65     | 0.35       | -        | 1.28       | 0.48     | 1.74       | 0.64     | 1.95       |
| \(0^+\)  | 0.68     | 0.33       | 0.88     | 0.33       | 1.33     | 0.50       | 1.77     | 0.63       |
| \(3^+\)  | 0.68     | 0.37       | -        | 1.36       | 0.55     | 1.74       | 0.64     | 2.11       |
| \(0^-\)  | -        | -          | 1.35     | 0.45       | 1.62     | 0.61       | 1.96     | 0.92       |
| \(1^-\)  | -        | -          | 1.40     | 0.59       | 1.59     | 0.63       | 2.02     | 0.81       |
| \(0^-\)  | -        | 0.92       | 0.39     | 1.25       | 0.51     | 1.82       | 0.62     | 2.02       |
| \(1^-\)  | 0.64     | 0.31       | -        | 1.46       | 0.53     | 1.76       | 0.59     | 2.08       |

be relatively smooth. We have not attempted to reproduce a \(1^-\)-resonance precisely at 1 MeV, but we still reproduce the properties almost within the experimental uncertainties.

The relatively large number of additional poles could be due to the Efimov effect, which is an anomaly in a three body system when the scattering lengths are much larger that the range of the interactions. With increasing scattering lengths, the infinitely many poles of the three-body S-matrix move towards the point \(E=0\). For very large but finite scattering lengths a number of poles must already appear close to zero. These poles originate from the long distance tail of the effective potential \(\propto \sum_i a_i \mu_{ik}^{-1} r^{-3}\), where \(a_i\) is the scattering length of the \(i\)-th subsystem) and they are not sensitive to the details of the interactions. Since there are no confining barriers for these poles, their corresponding widths must be rather large.

In our case the Efimov condition is almost fulfilled, since the scattering lengths are much larger then the range of the interactions: \(a_{nn}\mu_{nn}^{-1} + 2a_{cn}\mu_{cn}^{-1} \approx 50\) fm. This must necessarily result in a number of broad resonances near the \(E=0\) point.

**Transition strengths and Coulomb cross sections.** Dipole strength functions for different potentials are shown in Fig. 3 both for computed continuum wave functions and for plane waves. The interaction without spin splitting (ap-
Energy (MeV)

dB(E1)/dE (e² fm²/MeV)

Sackett 1993
Shimoura 1995
Zinser 1997
JJH (Kₘₐₓ = 10)
JJH (Kₘₐₓ = ∞)
sₐ = 0
sc = 3/2

Figure 3: The dipole strength functions dB/dE are shown here for several potentials for comparison. The curves peaked at low energy are for the correct continuum wave functions and the broader curves are for plane waves. The interactions for sₐ = 3/2 are the same as in Fig. 1 and for sₐ = 0 as described in table 1. JJH refers to the neutron-core potential from [18], where the result for the limited basis Kₘₐₓ = 10 is from [12]. The measured curves are from [8, 9, 10]. We normalize the data in [10] to our sum rule value, while the absolute data from [8, 9] is left unchanged.

approximation with zero core spin) gives a distribution shifted about 100 keV towards lower energy compared to the result for the realistic full computation. A lower and broader peak is obtained for the potential from [18] where the p²-content of the three-body wave function is very small. For comparison we also show the result for the same potential with the limited basis employed in [12]. The low-lying 1−-resonances enhance the strength functions at low energies compared to the plane wave computation. The computed strength functions substantially exceed most of the data points [8, 9, 10] in the peak region around 0.55 MeV. (Note that the data in [8, 9] contain much less total strength.) A reduction could be achieved with higher resonance energy and larger width, but this would probably only be provided by a potential with much too small p²-content in the three-body wave function. On the other hand, the proper comparison with the different experimental results is not obvious.

The Coulomb dissociation cross section is obtained by folding the virtual photon spectrum with the strength function. The precise expressions for the dipole approximation adopted here can be found in [13]. The low-energy enhancement places the strength where the virtual photon spectrum is relatively large and therefore necessarily also implies a large Coulomb cross section, since
the dipole is the dominating dissociation mode. The total cross sections for beam energies 42 MeV/A, 180 MeV/A, 280 MeV/A and 800 MeV/A on a lead target are 4.487 b, 2.128 b, 1.429 b and 0.971 b, respectively. This is in agreement with the available measured values [8], but substantially larger than 2.586 b, 1.166 b, 0.933 b and 0.657 b found by using plane waves. For 280 MeV/A we find 1458 b, 1501 b, 1283 b, respectively for the potentials without spin splitting, with small $p^2$-content with complete and limited basis.

Conclusions. A recently formulated method was used to compute low-energy three-body continuum spectra for $^{11}$Li for various angular momentum states. The angular part of the Faddeev equations are treated numerically at short distances, whereas the large distance behavior of eigenvalues and eigenfunctions is computed essentially analytically. Combining the results from these two regions allow accurate computations at large distances.

We employ interactions that reproduce reasonably well the low energy continuum properties of the two-body subsystems and the fragment momentum distributions in break-up reactions. The correct spins of $^9$Li and $^{11}$Li are used implying level splitting and fine structure beyond the approximation of zero core spin. A $1^-$-resonance at about 0.65 MeV with a width of 0.3 MeV is found in fair agreement with a recent measurement. Also several other s-matrix poles were found and related to the Efimov effect. Strength functions, Coulomb dissociation cross sections, phase shifts and S-matrix poles for $J^\pi = \frac{1}{2}^\pm, \frac{3}{2}^\pm, \frac{5}{2}^\pm$ are computed. A proper comparison of measured and computed strength functions would be rewarding either by showing the limitations of, or by selecting, the most correct three-body model.

Acknowledgments. A.C. acknowledges support by the European Union via the Human Capital and Mobility program contract nr. ERBCHBGCT930320.

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