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Triton Binding Energy and Three-Nucleon Potential

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Abstract. The method of continued fractions proposed by Horáček and Sasakawa is applied to calculate the binding energy of the triton. By this method, the Faddeev equation is solved very quickly. We solved it with 34 channels for the super-soft core (TRS), Paris (PARIS) and Argonne (AV) potentials with Tucson-Melbourne (TM) three-nucleon potential. The binding energy obtained from these realistic two-nucleon potentials without a three-nucleon potential is at most 7.7 MeV. If we include the TM three-nucleon potential taking a cutoff mass of $\Lambda = 800$ MeV for the dipole πNN form factor, the triton is overbound ($9.3 \sim 9.7$ MeV), whereas for $\Lambda = 700$ MeV, we get binding energies which almost agree with the experimental value; TRS + TM, 8.47; PARIS + TM, 8.32; AV + TM, 8.42 MeV.

1 Introduction

As a bridge between nuclear and particle physics, the triton should yield insight into fundamental questions such as, to what extent the three-nucleon force is important, to what extent the relativistic effect is important, in what manner the effect of quarks manifests itself, etc. As a first step to answer these questions, we might have to be provided with a good triton wave function, which reproduces the correct binding energy by a "classical" approach using potentials. In this case, in principle, the triton wave function is obtained if we solve the Faddeev equations. In practice, however, we should perform an extensive calculation. To meet this necessity, we demonstrate in the present paper a new way of calculating the triton binding energy by solving the Faddeev equations for various realistic two-nucleon potentials with a three-nucleon potential, and discuss the effect of the three-nucleon force to the binding energy of triton.

To make clear our discussions, first let us explain the meaning of the word "channel". The total spin of the triton is 1/2. Under this condition, the interacting pair and the spectator can take various angular momenta. Table 1 shows the first thirty-four possible angular momenta. In solving the Faddeev equations, if we take the first 5 (18, 26, 34) states, we say that we have performed a 5 (18, 26, 34)-channel calculation.

Up to 1983, all efforts to obtain the experimental value 8.48 MeV of the triton binding energy by solving the Faddeev equations with two-nucleon realistic

Table 1. 34 channels in the triton

Channel number	Interacting pair $^{2S+1}L_J$	Spectator l_j	Channel number	Interacting pair $^{2S+1}L_J$	Spectator l_j
1	$^1S_0$	$s_{1/2}$	18	$^3F_2$	$f_{5/2}$
2	$^3S_1$	$s_{1/2}$	19	$^1F_3$	$f_{5/2}$
3	$^3D_1$	$s_{1/2}$	20	$^1F_3$	$f_{7/2}$
4	$^3S_1$	$d_{3/2}$	21	$^3D_3$	$d_{5/2}$
5	$^3D_1$	$d_{3/2}$	22	$^3G_3$	$d_{5/2}$
6	$^3P_0$	$p_{1/2}$	23	$^3D_3$	$g_{7/2}$
7	$^1P_1$	$p_{1/2}$	24	$^3G_3$	$g_{7/2}$
8	$^1P_1$	$p_{3/2}$	25	$^3F_3$	$f_{5/2}$
9	$^3P_1$	$p_{1/2}$	26	$^3F_3$	$f_{7/2}$
10	$^3P_1$	$p_{3/2}$	27	$^3G_4$	$g_{7/2}$
11	$^3D_2$	$d_{3/2}$	28	$^3G_4$	$g_{9/2}$
12	$^3D_2$	$d_{5/2}$	29	$^1G_4$	$g_{7/2}$
13	$^1D_2$	$d_{3/2}$	30	$^1G_4$	$g_{9/2}$
14	$^1D_2$	$d_{5/2}$	31	$^3F_4$	$f_{7/2}$
15	$^3P_2$	$p_{3/2}$	32	$^3H_4$	$f_{7/2}$
16	$^3F_2$	$p_{3/2}$	33	$^3F_4$	$h_{9/2}$
17	$^3P_2$	$f_{5/2}$	34	$^3H_4$	$h_{9/2}$

potentials such as the Reid soft core potential (RSC) [1] or other potentials for not only the 5-channel but also 18-channel calculations [2] was in vain.

In 1984, we also performed 18-channel calculations [3] for some realistic potentials such as RSC, the velocity dependent potentials (URG) [4] and the Paris potential (PARIS) [5]. Later, we extended the calculation to include the Ueda-Green potential (UG) [6], the super-soft core potential (TRS) [7], and the Argonne potential (AV) [8]. The results show that so far as we take only account of two-nucleon realistic interactions the calculated binding energy of the triton is about 1 MeV short of the experimental value.

By 1983, some authors took account of the Tucson-Melbourne three-nucleon potential (TM) [9] in the first order perturbation energy E_3 for 5-channel calculations. In 1984, we performed, for the first time, the 18-channel calculation taking account of E_3 for TM three-nucleon potential. The results are given in ref. [3]. It shows that the 5-channel calculations cannot give a value beyond 8 MeV, whereas all 18-channel calculations yield values beyond it. This result gave a hope of getting the triton binding energy.

We found the reason why the effect of the three-nucleon potential is almost hidden. The matrix elements which give the largest contribution to the first order perturbational value E_3 are those between the interacting pair in the $^3S_1$ state with the spectator in the $s_{1/2}$ state and (1) the interacting pair in the $^3D_1$ state with the spectator in the $s_{1/2}$ state yielding -1.11 MeV (binding), and (2) the interacting pair in the $^3S_1$ state with the spectator in the $d_{3/2}$ state yielding 1.11 MeV to E_3. Thus, these two largest contributions cancel out. This is for the 18-channel calculation of RSC (RSC18). For RSC5 and for other realistic poten-

tials, the cancellation is not so drastic as in RSC18, but the situation remains almost the same. In RSC5, the sum of these two matrix elements gives even a repulsive effect, as seen in Table 4 of ref. [3]. This is one of the reasons why the effect of the three-nucleon potential was not easily detected. Another reason is the following. Among 171 independent matrix elements for E_3, 16 matrix elements exceed 0.1 MeV in the absolute value. The sum of these matrix elements yields the binding energy $-E_3$ of 0.24 MeV, whereas the sum of all matrix elements yields 0.89 MeV. This shows that the contributions from small components are very important, and the 5-channel calculation does not take enough matrix elements. We also found that E_3 strongly depends on the πNN form factor and its cutoff mass [3].

In the present paper, we shall discuss the result of the 34-channel calculation for realistic two-nucleon potentials (2NP), specifically, TRS, PARIS and AV with the three-nucleon potential TM. To perform this calculation, we extended the method of continued fractions (MCF) [10] to handle the three-nucleon bound state. By this method, we can solve the 34 (26, 18)-channel Faddeev equation with 2NP + TM in 600 (300, 180) seconds by a computer available at Tohoku University (ACOS 1000). Since MCF is so efficient, we think that it is worthwhile to describe this method in some detail. This will be done in Sect. 2. The result of calculations for 2NP as well as 2NP + TM will be discussed in Sect. 3. The summary and conclusion will be given in Sect. 4.

2 Method of Continued Fractions (MCF) Applied to Faddeev Equation

The MCF was originally intended to handle potential scattering. Examples given in ref. [10] show that this method is very efficient both for local and non-local potentials. In the present section, we extend the method to the treatment of bound state of a three-nucleon system.

One of the merits in this method is that we need *not* solve equations in a coupled form for a three-nucleon system, where the masses of the three nucleons are the same. We need only calculate the coupled matrix elements, whatever the number of channels is. Thus, once we store the matrix elements, the rest of the calculations, for which MCF is applied, needs almost no time. Indeed, the computational time that was written in Sect. 1 is for MCF. Another merit is its flexibility within a framework of continued fractions. In fact, each article in ref. [10] presents a different version of the method. The method that will be described in the present section is a further version. By making use of this flexibility, we can choose a version that minimizes the computational time, even when we write out the convergence in terms of order-by-order (and not in terms of the computational time) an adopted version is a little bit slower than other versions. In any case, the big merits are its simplicity of the algorithm and, of course, a very quick convergence.

2.1 Faddeev Equation

The solution Ψ of the three-body Schrödinger equation

$$(E - H_0 - V_{12} - V_{23} - V_{31})\Psi = 0 \tag{1}$$

is expressed as a sum of three components

$$\Psi = \Phi(12, 3) + \Phi(23, 1) + \Phi(31, 2). \tag{2}$$

The component $\Phi(12, 3)$ represents the state in which a pair of particles 1 and 2 are interacting in the final state, while the particle 3 is standing as the spectator. With (2), we decompose (1) into

$$(E - H_0 - V_{12})\Phi(12, 3) = V_{12}(\Phi(23, 1) + \Phi(31, 2)) \tag{3}$$

and two other equations which are obtained from (3) by cyclic permutations of 1, 2, and 3. Eq. (3) is called the Faddeev equation.

If three-nucleon forces

$$W_{123} = W_{12,3} + W_{23,1} + W_{31,2} \tag{4}$$

are acting in the three-nucleon system, we may extend the Faddeev equation (3) as

$$\begin{aligned}(E - H_0 - V_{12})\Phi(12, 3) = {} & V_{12}[\Phi(23, 1) + \Phi(31, 2)] \\ & + W_{12,3}[\Phi(12, 3) + \Phi(23, 1) + \Phi(31, 2)].\end{aligned} \tag{5}$$

Of course, there are other ways of extending (3) to accommodate three-nucleon potentials. In any way, if we solve the equation as exactly as possible, the way of accommodating the three-nucleon forces should not affect the calculated binding energy. Recently, the Los Alamos group [11] showed that if we perform the 34-channel calculations, the way of extending (3) to include (4) does not affect the result.

For simplicity, we use the following notations:

$$\begin{aligned}&\Phi \equiv \Phi(12, 3), \qquad \hat{Q}\Phi = \Phi(23, 1) + \Phi(31, 2), \\ &V \equiv V_{12}, \qquad\qquad U \equiv V\hat{Q} + W_{12,3}(1 + \hat{Q}).\end{aligned} \tag{6}$$

Eq. (5) reads then

$$(E - H_0 - V)\Phi = U\Phi. \tag{7}$$

Further, if we denote by G_0 the Green function

$$G_0 = \frac{1}{E - H_0 - V}, \tag{8}$$

Eq. (7) is expressed for a bound state as

$$\Phi = G_0 U \Phi. \tag{9}$$

This is the equation that we have to solve.

2.2 Equation for Binding Energy

Let $|F_0\rangle$ and $|f\rangle$ be functions which are regular at the origin of the relative coordinate for an interacting pair as well as of the coordinate for the spectator relative to the center of mass of the interacting pair. Both of these functions are chosen so that these functions decrease when three particles are parted from each other to a large distance. Except for these requirements, the choice of these functions is rather arbitrary.

Defining the potential U_1 by

$$U_1 = U - \frac{U|F_0\rangle\langle f|U}{\langle f|U|F_0\rangle}, \tag{10}$$

we use (10) in (9) to obtain

$$|\Phi\rangle = G_0\left[U_1|\Phi\rangle + \frac{U|F_0\rangle\langle f|U|\Phi\rangle}{\langle f|U|F_0\rangle}\right] \tag{11}$$

$$= \frac{1}{1 - G_0U_1} G_0U|F_0\rangle \frac{\langle f|U|\Phi\rangle}{\langle f|U|F_0\rangle}. \tag{12}$$

Further, we introduce two functions $|F_1\rangle$ and $|\Phi_1\rangle$ by

$$|F_1\rangle = G_0U|F_0\rangle \tag{13}$$

and

$$|\Phi_1\rangle = \frac{1}{1 - G_0U_1}|F_1\rangle = |F_1\rangle + G_0U_1|\Phi_1\rangle. \tag{14}$$

If we multiply $\langle f|U$ from the left of (12), we obtain

$$\langle f|U|\Phi_1\rangle - \langle f|U|F_0\rangle = 0. \tag{15}$$

This equality should be satisfied at a bound state. Namely, the solution of (15) determines the binding energy of the three-nucleon system.

2.3 Method of Continued Fractions

Now, our task is to calculate the matrix element $\langle f|U|\Phi_1\rangle$ in (15). In this subsection, we show that this matrix element is calculated as a continued fraction. We note that (14) takes a form of the scattering equation with $|F_1\rangle$ as the initial state. Therefore, the method similar to ref. [10], which is designed for scattering problems, may be applied.

We define the potential U_2 by

$$U_2 = U_1 - \frac{U_1|F_1\rangle\langle f|U}{\langle f|U|F_1\rangle}. \tag{16}$$

In the first paper of ref. [10], we defined U_2 in another way. This is an example of the flexibility of MCF.

If we put (16) in (14), we obtain the following equality after some calculations:

$$\langle f|U|\Phi_1\rangle = \frac{\langle f|U|F_1\rangle^2}{\langle f|U|F_1\rangle - \langle f|U|\Phi_2\rangle}, \tag{17}$$

where Φ_2 is defined by

$$|\Phi_2\rangle = |F_2\rangle + G_0U_2|\Phi_2\rangle, \tag{18}$$

$$|F_2\rangle = G_0U_1|F_1\rangle. \tag{19}$$

In the course of deriving (17), we used a relation

$$G_0 U_2 |F_1\rangle = 0. \tag{20}$$

A similar manipulation is extended to any order n:

$$|F_n\rangle = G_0 U_{n-1} |F_{n-1}\rangle, \tag{21}$$

$$|\Phi_n\rangle = |F_n\rangle + G_0 U_n |\Phi_n\rangle, \tag{22}$$

$$U_n = U_{n-1} - \frac{U_{n-1}|F_{n-1}\rangle\langle f|U}{\langle f|U|F_{n-1}\rangle}. \tag{23}$$

Using these functions, we obtain continued fractions

$$\langle f|U|\Phi_n\rangle = \frac{\langle f|U|F_n\rangle^2}{\langle f|U|F_n\rangle - \langle f|U|\Phi_{n+1}\rangle}. \tag{24}$$

To get $\langle f|U|\Phi_1\rangle$ in (15), we start from (24) for some chosen number N and assume that

$$\Phi_N = F_N. \tag{25}$$

Then we calculate $\langle f|U|\Phi_N\rangle$, $\langle f|U|\Phi_{N-1}\rangle$ etc. by (24), until we calculate $\langle f|U|\Phi_1\rangle$.

2.4 Discussions for MCF

The functions $|F_0\rangle$ and $|f\rangle$ are arbitrary. As to $|f\rangle$, we have chosen the following form. Let $\vec{x}$ be the relative coordinate between the interacting pair 1 and 2, and $\vec{y}$ be the coordinate of the spectator 3 relative to the center of mass of the pair 1 and 2. We designate by q (p) momentum of the interacting pair (the spectator). For a triton of the binding energy $-|E|$, p and q satisfy the relationship

$$-|E| = (\hbar^2/M)q^2 + (3\hbar^2/4M)p^2, \tag{26}$$

where M is the nucleon mass. We take p as real, and q as pure imaginary; $q = i|q|$. Let $\phi_{1S}(|q|, x)$ [$\phi_{3S}(|q|, x)$ and $\phi_{3D}(|q|, x)$] be the normalized Sturm-Liouville function of the ${}^1S_0$ [${}^3S_1$ and ${}^3D_1$] two-body state for a given energy $-(\hbar^2/M)q^2$, multiplied by $(\lambda_q/(1-\lambda_q))^{1/2}$, where λ_q is the eigenvalue. We let $u_0(py)$ stand for the normalized plane wave in the s state:

$$u_0(py) = \sqrt{\frac{2}{\pi}}\, p \sin py/(py). \tag{27}$$

Denoting by χ the spin function of the spectator, we choose the function $|f\rangle$ to be given by

$$|f\rangle = \int_0^{p_M} dp\, u_0(py)\{\phi_{1S}(|q|, x)[{}^1S_0(\hat{x}) \otimes \chi]_{J=1/2} + \phi_{3S}(|q|, x)[{}^3S_1(\hat{x}) \otimes \chi]_{J=1/2} + \phi_{3D}(|q|, x)[{}^3D_1(\hat{x}) \otimes \chi]_{J=1/2}\}, \tag{28}$$

where ${}^1S_0(\hat{x})$ etc. denote the spin-angular function of the interacting pair and J represents the total spin of triton. In (28), p_M is a cut-off momentum which is chosen judiciously. However, since the function $|f\rangle$ is a trial function, we need not to be nervous in choosing p_M.

As for $|F_0\rangle$, we have taken the following function,

$$|F_0\rangle = G_0 V\hat{Q}|f\rangle. \tag{29}$$

For any perturbation method to converge very quickly, the starting function should be chosen so that it has an important character of the solution. This general requirement is fulfilled by the above choice of $|F_0\rangle$. Since the operator $G_0 V\hat{Q}$ in (29) is a part of the kernel of (9), the function $|F_0\rangle$ and hence $|F_1\rangle$, which is the starting function of MCF, should be very similar in its behavior to the true wave function. Especially, due to the presence of the permutation operator $\hat{Q}$ in (29) and U in (9), the functions $|F_0\rangle$ and hence $|F_1\rangle$ have a node which is characteristic of the Faddeev component for a soft core potential [12]. If we started the calculation of the continued fraction by such a function without a node as (28), the convergence would be very slow.

In Sect. 3, we shall demonstrate the result of calculations taking

$$|f\rangle = (28) \tag{30}$$

in all equations. However, since $|f\rangle$ is arbitrary, some other choice is of course possible. For instance, taking (30) up to (15) and

$$|f\rangle = |F_n\rangle \tag{31}$$

for $n \geqslant 2$ in (23) including (16), we made the order-by-order comparison of the left-hand side of (15) for RSC5 taking $E = 7.031$ MeV. In the case (I) [(II)] of Table 2, the result of the choice (30) [(31)] is demonstrated. From this table, we see that the choice (31) is better than (30) for the order-by-order convergence. However, the choice (31) takes more time than the choice (30), because for (31)

Table 2. Order-by-order comparison of convergence, used (30) for $|f\rangle$ (Case I) and (31) for $|f\rangle$ (Case II). Diff(m) denotes the value of the right-hand side of (15) when the continued fraction (24) starts from m [$=N$ in (25)]. This table illustrates the calculation for RSC5 ($E = 7.031$ MeV)

	Diff(m)	
m	(I)	(II)
1	−0.02529	−0.22963
2	−0.00416	−0.12590
3	0.01555	−0.00512
4	0.00163	0.00465
5	0.00039	−0.00168
6	−0.00292	−0.00044
7	−0.00012	−0.00008
8	−0.00006	0.00002
9	−0.00068	−0.00004
10	−0.00001	0.00000
11	0.00000	0.00000

we have to calculate $\langle F_i|U|\Phi_j\rangle$, $N-1 \geqslant j \geqslant i+1$, thus one more loop than the choice (30) [although, in practice, the difference of the computational time is not significant for a very fast computer].

3 Triton Binding Energy

Using the technique given in Sect. 2, we solved the Faddeev equation taking TRS, PARIS and AV as the two-nucleon potentials and TM as the three-nucleon potential for 5, 18, 26 and 34 channels. We show the calculated results without (with) including TM in Table 3 (Table 4). Also, we add, for comparison, the results from RSC, URG and UG without TM in Table 3, and RSC with TM in Table 4 for 5- and 18-channel calculations. The results of the Los Alamos group [11] are shown in parenthesis. [For 26-channel calculations, the states which are taken into account are different for both groups. The twenty-six states of the Los Alamos group include all channels in Table 1, provided that eight states with odd L (odd l) are deleted below state number 19.] In our calculation, the dipole form with cutoff mass $\Lambda = 800$ MeV is taken as the πNN form factor in TM. Although the method of calculation by the Los Alamos group and by us are completely different, the agreement of the results in Table 3 is impressive, which shows that the binding energy obtained from realistic two-nucleon potentials is at most 7.7 MeV. Since TM is quite singular at small distances, a difference of treatments causes some difference in the result. Nevertheless, the results from Los Alamos almost agree with ours in Table 4. Thus, we conclude that if we include the TM three-nucleon potential, we get overbinding for any realistic two-nucleon potential. Here, we record some numbers used in our calculations: $x_{\max} = 8$ fm, 40 mesh points for $0 \leqslant x \leqslant 8$ fm, and $p_{\max} = 6.03825$ fm^{-1}.

In ref. [3], we have pointed out that the calculated binding energy depends strongly on the assumed πNN form factor and the cutoff mass. Table 5 lists the

Table 3. Triton binding energy for some realistic two-nucleon potentials without TM three-body force

Number of channels	RSC	URG	UG	TRS	PARIS	AV
5	7.03 (7.02)	7.48	7.40	7.46	7.48	7.45 (7.44)
18	7.24 (7.23)	7.50	7.44	7.49	7.56	7.58 (7.57)
26				7.55	7.63	7.67
34				7.55	7.64	7.68 (7.67)

Table 4. Triton binding energy for some realistic two-nucleon potentials with TM

Number of channels	RSC + TM	TRS + TM	PARIS + TM	AV + TM
5	7.56 (7.55)	8.62	8.27	8.18 (8.26)
18	9.11 (8.93)	10.08	9.49	9.64 (9.49)
26		9.62	9.06	9.16
34	(8.86)	9.71	9.18	9.29 (9.36)

Table 5. Λ-dependence of (2NP + TM)34

Potential	2NP	2NP + TM Λ = 700	800
TRS	7.55	8.47	9.71
PARIS	7.64	8.32	9.18
AV	7.68	8.42	9.29

34-channel results (2NP + TM)34, with the dipole form factor for Λ = 700 and 800 MeV. This table shows that if we take TM as a phenomenological three-nucleon potential with Λ as a parameter, it is possible to find a value of Λ near 700 MeV, that fits the triton binding energy, for each of realistic 2NP + TM.

4 Summary and Conclusion

The method of continued fractions [10], which was originally intended to handle scattering problems, is extended to treat a bound state. It is applied to the Faddeev equation with two and three nucleon potentials. By this method, we can get the binding energy in a very short time.

Using this method, we have solved the Faddeev equation for the triton with the super-soft core, Paris and Argonne two-nucleon and the Tucson-Melbourne three-nucleon potentials for 34 channels.

The binding energy obtained from these three realistic potentials is at most 7.7 MeV, if we omit the three-nucleon potential (Table 3). If we solve the Faddeev equation including the Tucson-Melbourne three-nucleon potential, the triton is overbound for all of these realistic potentials (Table 4), for the dipole πNN form factor with a cutoff mass of Λ = 800 MeV. If we take Λ = 700 MeV, we obtain the triton binding energy which almost agrees with the experimental value: (TRS + TM)34, 8.47 MeV; (PARIS + TM)34, 8.32 MeV; (AV + TM)34, 8.42 MeV.

At this moment, we are assuming that the Tucson-Melbourne potential is the only three-nucleon potential, neglecting all other kinds of three-nucleon potentials. Therefore, the value of Λ should be taken as a parameter that represents all other effects, not only the two-pion exchange three-nucleon process. Nevertheless, it is interesting to see that a value of Λ which is not so strange can reproduce the triton binding energy.

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