$N\Omega$ dibaryon from lattice QCD near the physical point

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Abstract

The nucleon($N$)-Omega($\Omega$) system in the S-wave and spin-2 channel ($^5S_2$) is studied from the (2+1)-flavor lattice QCD with nearly physical quark masses ($m_\pi \simeq 146$ MeV and $m_K \simeq 525$ MeV). The time-dependent HAL QCD method is employed to convert the lattice QCD data of the two-baryon correlation function to the baryon-baryon potential and eventually to the scattering observables. The $N\Omega(^5S_2)$ potential is found to be attractive in all distances and to produce a quasi-bound state near unitarity: In this channel, the scattering length, the effective range and the binding energy from QCD alone read $a_0 = 5.30(0.44)(^{+0.16}_{-0.07})$ fm, $r_{\text{eff}} = 1.26(0.01)(^{+0.04}_{-0.01})$ fm, $B = 1.54(0.30)(^{+0.04}_{-0.01})$ MeV, respectively. Including the extra Coulomb attraction, the binding energy of $p\Omega^-(^5S_2)$ becomes $B_{p\Omega^-} = 2.46(0.34)(^{+0.04}_{-0.11})$ MeV. Such a spin-2 $p\Omega^-$ state could be searched through two-particle correlations in $p$-$p$, $p$-nucleus and nucleus-nucleus collisions.

Keywords: dibaryon, Lattice QCD, hyperon interaction

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1. Introduction

Quest for dibaryons is a long-standing experimental and theoretical challenge in hadron physics [1, 2]. Among various theoretical attempts to study dibaryons, one of the recent highlights is the (2+1)-flavor lattice QCD simulations near the physical point ($m_\pi \simeq 146$ MeV and $m_K \simeq 525$ MeV) by HAL QCD Collaboration. (For a recent summary, see Ref. [3].) This enables us to make model-independent investigations of the elusive $H$-dibaryon, originally proposed by the MIT bag model [4], on the basis of a coupled channel analysis of the lattice QCD data [5]. Also, the possible di-Omega ($\Omega\Omega$), originally proposed by the Skyrme model [6], has recently been examined in detail from the same lattice QCD data [7].

Another interesting candidate of the dibaryon is $N\Omega$ and ($uudss$ or $uddss$) in the $^5S_2$ channel. Since the Pauli exclusion does not operate among valence quarks and the color-magnetic interaction is attractive in the channel, it was predicted to be a resonance below the $N\Omega$ threshold in the constituent quark model [8, 9]. Moreover, $N\Omega(^5S_2)$ is expected to have relatively a small width since its strong decay into octet baryons such as $\Lambda\Xi$ and $\Sigma\Xi$, which must have orbital D-wave, would be kinematically suppressed. A pilot (2+1)-flavor lattice QCD simulations with a heavy pion mass ($m_\pi \simeq 875$ MeV) [10] suggests a short-range attraction between $N$ and $\Omega$ in the $^5S_2$ channel. Subsequently, theoretical studies on the $N\Omega$ system [11, 12, 13, 14, 15] as well as experimental measurements in relativistic heavy ion collisions [16] have been reported.

The purpose of this Letter is to study $N\Omega(^5S_2)$ on the basis of realistic (2+1)-flavor lattice QCD simulations near the physical point ($m_\pi \simeq 146$ MeV and $m_K \simeq 525$ MeV). As in the case of our previous pilot study [10], we employ the HAL QCD method [17, 18, 19] which allows us to extract the interaction between $N$ and $\Omega$ from the spacetime dependence of the two-baryon correlation function on the lattice.

This paper is organized as follows. In Sec. 2, we introduce the HAL QCD method to extract the hadron interaction from lattice QCD. In Sec. 3, we summarize the setup of our lattice QCD simulations near the physical point. In Sec. 4, we analyze the $N\Omega$ system in $^5S_2$ channel in detail. Sec. 5 is devoted to summary and concluding remarks.
2. HAL QCD method

Let us consider the \( N\Omega(5S_2) \) characterized by the following two-baryon correlation function,

\[
C_{N\Omega}(\vec{r}, t) = \frac{1}{24} \sum_{\mathcal{R} \in \mathcal{O}} \sum_{\vec{x}} P^{(s=2)}_{\alpha\beta,\ell}(0) N_{\alpha}(\mathcal{R}[\vec{r}] + \vec{x}, t) \Omega_{\beta,\ell}(\vec{x}, t) \mathcal{J}_{N\Omega}(0) \vert 0 \rangle, \tag{1}
\]

with \( \mathcal{J}_{N\Omega} \) being the wall-type quark source. The interpolating operators for the nucleon and the \( \Omega \)-baryon are

\[
N_{\alpha}(x) = \epsilon_{abc}(u^a T(x) C\gamma_5 d^b(x)) q_c^\alpha(x), \quad \Omega_{\beta,\ell}(x) = \epsilon_{abc}s_a^\beta(x)(s^b T(x) C\gamma_\ell s^c(x)), \tag{2}
\]

where \( \alpha \) and \( \beta \) are Dirac indices, \( \ell \) is a spatial label of gamma matrices, \( a, b, c \) are the color indices and \( C \equiv \gamma_4 \gamma_2 \) and Dirac indices are restricted to the upper two components. The summation over the cubic group element \( \mathcal{R} \in \mathcal{O} \) leads to a projection onto the S-wave state\(^1\). On the other hand, the projection operator onto the spin-2 state is chosen to be

\[
P^{(s=2)}_{\alpha\beta,\ell} = \left[ \delta_{\alpha,+1/2} \delta_{\beta,+1/2} (\delta_{\ell,2} + i \delta_{\ell,1}) + \delta_{\alpha,-1/2} \delta_{\beta,-1/2} (\delta_{\ell,2} - i \delta_{\ell,1}) \right] / 4. \]

Here we neglect the coupling of \( N\Omega(5S_2) \) to the D-wave octet-octet channels, \( \Lambda\Xi \) and \( \Sigma\Xi \), below the \( N\Omega \) threshold. In the present paper, we assume the coupling between the S-wave and the D-wave is kinematically suppressed and focus only on the single channel\(^2\).

It is convenient to define the following ratio which we call the “\( R \)-correlator”,

\[
R_{N\Omega}(\vec{r}, t) \equiv \frac{C_{N\Omega}(\vec{r}, t)}{C_{N}(t) C_{\Omega}(t)}, \tag{3}
\]

where \( C_N(t) \) and \( C_{\Omega}(t) \) are single-baryon correlators. Below the inelastic threshold, \( R_{N\Omega}(\vec{r}, t) \) can be shown to satisfy the integro-differential equation

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\(^1\)Strictly speaking, this operation projects onto the \( A_1^+ \) state which contains not only \( L = 0 \) state but also \( L = 4, 6, \cdots \) states in the continuum theory.

\(^2\)A recent phenomenological study with such S-wave – D-wave coupling\([14]\) indicates that the contribution from the D-wave channels to the volume integral of the \( N\Omega(5S_2) \) interaction is \( \sim 1\% \) and is negligible. Nevertheless, further analysis with the coupled-channel HAL QCD method\([21]\) would be necessary in the future to validate our assumption. Note that such an assumption is not justified for \( N\Omega(5S_1) \) which can decay into S-wave octet-baryons.
with a non-local and energy-independent kernel $U(\vec{r}, \vec{r}')$ \cite{19},

$$
\left[ -\frac{\partial}{\partial t} + \frac{1 + 3\delta^2}{8m} \frac{\partial^2}{\partial t^2} + \mathcal{O}(\delta^2 \partial_t^3) \right] R(\vec{r}, t) = H_0 R(\vec{r}, t) + \int U(\vec{r}, \vec{r}') R(\vec{r}', t) d\vec{r}',
$$

(4)

with $H_0 \equiv -\nabla^2/2m$, the reduced mass $m \equiv (m_Nm_\Omega)/(m_N + m_\Omega)$ and the asymmetry parameter $\delta \equiv (m_N - m_\Omega)/(m_N + m_\Omega)$.

The central potential in the leading-order (LO) analysis under the derivative expansion

$$
U(\vec{r}, \vec{r}') = \sum_n V_n(\vec{r}) \nabla^n \delta(\vec{r} - \vec{r}')
$$

is given by

$$
V_C(r) = -\frac{H_0 R_{N\Omega}(\vec{r}, t)}{R_{N\Omega}(\vec{r}, t)} - \frac{(\partial/\partial t) R_{N\Omega}(\vec{r}, t)}{R_{N\Omega}(\vec{r}, t)} + \frac{1 + 3\delta^2}{8m} \frac{(\partial^2/\partial t^2) R_{N\Omega}(\vec{r}, t)}{R_{N\Omega}(\vec{r}, t)},
$$

(5)

up to $\mathcal{O}(\delta^2 \partial_t^3)$-terms in the right hand side. Spatial and temporal derivatives on the lattice at $(\vec{r}, t)$ are calculated in central difference scheme using nearest neighbour points. If $R_{N\Omega}(\vec{r}, t)$ is dominated by a single state (ground state) at large $t$, each term in the r.h.s. of Eq. (5) should have no $t$-dependence. Such a ground state saturation, however, is not necessary to obtain $V_C(r)$ in the time-dependent HAL QCD method as long as $R_{N\Omega}(\vec{r}, t)$ is dominated by the elastic states. In general, each term in the r.h.s. receives $t$-dependence which provides “signal” instead of “noise” for $V_C(r)$. (If there remains residual $t$-dependence in $V_C(r)$, the next-to-leading order of the derivative expansion must be taken into account \cite{20}.) This is why the data at moderate values of $t \sim 1$ fm are sufficient to extract the baryon-baryon interaction in HAL QCD method.\cite{4}

3. Lattice Setup

Gauge configurations are generated by using the (2+1)-flavor lattice QCD with the Iwasaki gauge action at $\beta = 1.82$ and the non-perturbatively $\mathcal{O}(a)$-improved Wilson quark action with the six APE stout smearing with the

3Good convergence of this derivative expansion at low energies for the point-sink scheme is demonstrated in Ref. \cite{20}.

4This is in sharp contrast to the so-called “finite volume method”. It requires strict ground state saturation, so that very large value of $t > 10$ fm is necessary. For such large $t$, however, no signal can be obtained due to the explosion of statistical errors. See \cite{22,23,24} for explicit demonstration of this fact.
smearing parameter \( \rho = 0.1 \) at nearly physical quark masses \((m_\pi \simeq 146 \text{ MeV}
and m_K \simeq 525 \text{ MeV})\). The lattice cutoff is \( a^{-1} \simeq 2.333 \text{ GeV} \) \((a \simeq 0.0846 \text{ fm})\) and the lattice volume \( L^4 \) is 96\(^4\), corresponding to \( La \simeq 8.1 \text{ fm} \). This is sufficiently large volume to accommodate two baryons. We employ the wall-type quark source with the Coulomb gauge fixing. The periodic (Dirichlet) boundary condition for the spatial (temporal) direction is imposed for quarks. The quark propagators are obtained by using the domain-decomposed solver [26, 27, 28, 29], and the unified contraction algorithm is employed to calculate the correlation functions [30].

The forward and backward propagations are averaged and the hypercubic symmetry on the lattice (4 rotations) are utilized for each configuration. 414 configurations are available by picking up one per five trajectories: For 207 configurations which are separated by ten trajectories, 48 source locations are used, while 24 source locations are used for the rest (207 configurations), and the total number of measurements read 119,232. The statistical errors are estimated by the jackknife method with 20 samples (bin size 5,952 measurements). We have checked that the bin size dependence is small by comparing the result with 40 samples (bin size 2,880 measurements). The fit to the effective mass in the range \( 12 \leq t/a \leq 17 \) for \( N \) and \( 17 \leq t/a \leq 22 \) for \( \Omega \) lead to \( m_N = 954.7(2.7) \text{ MeV} \) and \( m_\Omega = 1711.5(1.0) \text{ MeV} \). These values are about 2\% heavier than physical values due to a slight difference of the present quark masses from the physical point.

4. Spin-2 \( N\Omega \) potential

Shown in Fig. 1 is the \( R \)-correlator defined by Eq. (3) in the range \( t/a = 10 - 15 \), which are rescaled by the value of \( r = 3 \text{ fm} \). At large \( r \), the \( R \)-correlator approaches a constant. This implies that \( V_C(r) \) in Eq. (5) becomes a constant at long distance. At small \( r \), the \( R \)-correlator increases with the second-order derivative in \( r \) being always positive, which implies that there is an attractive potential at short distances. The weak \( t \)-dependence at small \( r \) indicates contributions from the elastic scattering states. As mentioned before, this \( t \)-dependence provides signal instead of noise.

To extract \( V_C(r) \) from the \( R \)-correlator, we choose \( t/a = 11 - 14 \) in order to reduce the systematic uncertainties. For smaller values of \( t \), the inelastic

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\(^5\)Due to the presence of time derivatives up to \( \mathcal{O}(\partial_t^2) \), the actual lattice data used in our analysis are in the interval \( 10 \leq t/a \leq 15 \).
contribution starts to appear so that $V_C(r)$ remains non-vanishing even for large $r$. For larger values of $t$, it is difficult to control the systematic uncertainties of the fitting of the potential due to the large statistical errors. Note that we take relatively larger values of $t/a$ to make accurate determination of $m_N$ and $m_\Omega$, whose values agree with the effective masses at $t/a = 12$ in 1%.

In Fig. 2, $V_C(r)$ as well as its breakdown into different components are shown for $t/a = 12$ as an example. First of all, $V_C(r)$ (red squares) is attractive everywhere. This is qualitatively consistent with the result in our pilot study with heavy pion mass ($m_\pi \simeq 875$ MeV) [10]. Also, we found that the $H_0$-term (blue circles) is dominant, yet the $\partial/\partial t$-term (green triangles) gives non-negligible $r$-dependent contribution. On the other hand, the $\partial^2/\partial t^2$-term (orange diamonds) is consistent with zero.

We summarize the central potential $V_C(r)$ in Fig. 3(a) for $t/a = 11 - 14$. These potentials are consistent with each other within statistical errors, which is an indirect evidence of the small coupling with the D-wave octet-octet states below the $N\Omega$ threshold in the spin-2 channel as suggested by [14]. (Such a stability of the potential in the same range of $t$ is not found for the spin-1 $N\Omega$ system which can couple to the S-wave octet-octet states below threshold.)

To obtain observables such as the scattering phase shifts and binding
energy, we fit the lattice QCD potential by Gaussian + (Yukawa)\(^2\) with a form factor [10]:

\[
V_{\text{fit}}(r) = b_1 e^{-b_2 r^2} + b_3 \left(1 - e^{-b_4 r^2}\right)^n \left(\frac{e^{-m_\pi r}}{r}\right)^2.
\]

(6)

The (Yukawa)\(^2\) form at long distance is motivated by the picture of two-pion exchange between \(N\) and \(\Omega\) with an OZI violating vertex [14]. The pion mass in Eq. (6) is taken from our simulation, \(m_\pi = 146\) MeV. After trying both \(n = 1\) and \(2\) in the form factor, we found that only \(n = 1\) can reproduce the short distance behavior of the lattice potential, so that we will focus on the \(n = 1\) case below. The results of the fit and the corresponding parameters are summarized in Fig. 3(b,c,d,e) and Table 1, respectively.

Shown in Fig. 4 (Left) is the S-wave scattering phase shift \(\delta_0\) as a function of the kinetic energy. The values of \(k \cot \delta_0\) are also shown in Fig. 4 (Right). These results for \(t/a = 11, 12, 13\) and \(14\) are consistent with each other within the statistical errors. In the \(k \to 0\) limit, the phase shift approaches to 180\(^\circ\), and the scattering length \(a_0 \equiv -\lim_{k \to 0} \tan \delta_0/k\), becomes positive. This implies that the existence of a quasi-bound state of \(N\Omega\) in the \(^5S_2\) channel.

\(^6\)Here, the sign of the scattering length is defined to be opposite to that in [10].
Figure 3: (a) The central potential $V_C(r)$ of the $N\Omega(5S_2)$ system at $t/a = 11$ (blue up-pointing triangles), 12 (red squares), 13 (green circles) and 14 (black down-pointing triangles). (b) The result of the fitting of $V_C(r)$ (red circles) at $t/a = 11$ by using $V_{fit}(r)$ in Eq. (6). The black dotted (orange solid) line denotes the first (second) term in Eq. (6), and the blue dashed line is the sum of two terms. (c), (d) and (e) are the cases of $t/a = 12$, 13 and 14, respectively.
Table 1: The fitting parameters in Eq. (6) in physical unit with the statistical errors.

| $t/a$ | 11   | 12   | 13   | 14   |
|------|------|------|------|------|
| $b_1$ [MeV] | -306.5(5.5) | -313.0(5.3) | -316.7(9.4) | -296(18) |
| $b_2$ [fm$^{-2}$] | 73.9(4.4) | 81.7(5.4) | 81.9(8.4) | 64(16) |
| $b_3$ [MeV.fm$^2$] | -266(32) | -252(27) | -237(43) | -272(109) |
| $b_4$ [fm$^{-2}$] | 0.78(11) | 0.85(10) | 0.91(18) | 0.76(34) |

Figure 4: (Left) The S-wave scattering phase shifts $\delta_0$ as a function of the kinetic energy, $k^2/2m$. (Right) $k \cot \delta_0/m_\pi$ as a function of $(k/m_\pi)^2$.

The effective range expansion (ERE) of the phase shifts up to the next-leading-order (NLO) reads

$$k \cot \delta_0 = -\frac{1}{a_0} + \frac{1}{2} r_{\text{eff}} k^2 + O(k^4)$$  \hspace{1cm} (7)

with $r_{\text{eff}}$ being the effective range. The ERE parameters ($a_0$, $r_{\text{eff}}$) obtained from our phase shifts are found to be

$$a_0 = 5.30(0.44)(^{+0.16}_{-0.01}) \text{ fm}, \quad r_{\text{eff}} = 1.26(0.01)(^{+0.02}_{-0.01}) \text{ fm},$$  \hspace{1cm} (8)

where the central values and the statistical errors are estimated at $t/a = 12$, while the systematic errors in the last parentheses are estimated from the central values for $t/a = 11, 13$ and 14.

In Fig. 5 the ratio $r_{\text{eff}}/a_0$ as a function of $r_{\text{eff}}$ for $N\Omega(^5S_2)$ is plotted together with the experimental values for $NN(^3S_1)$ (deuteron) and $NN(^1S_0)$ (di-neutron) as well as lattice QCD value for $\Omega\Omega(^1S_0)$ (di-Omega) \cite{7}. Small
Figure 5: The ratio of the effective range $r_{\text{eff}}$ and the scattering length $a_0$ as a function of $r_{\text{eff}}$ for $N\Omega(^5S_2)$ (red circle) and $\Omega\Omega(^1S_0)$ [7] (blue diamond) on the lattice, as well as for $NN(^3S_1)$ (purple up-pointing triangle) and $NN(^1S_0)$ (green down-pointing triangle) [31] in experiments.

values of $|r_{\text{eff}}/a_0|$ in all these cases indicate that these systems are located close to the unitary limit.

The binding energy $B$ and the root mean square distance ($\sqrt{\langle r^2 \rangle}$) of $N\Omega(^5S_2)$ are obtained by solving the Schrödinger equation with the potential fitted to our lattice results:

$$B = 1.54(0.30)(^{+0.04}_{-0.10}) \text{ MeV}, \quad \sqrt{\langle r^2 \rangle} = 3.77(0.31)(^{+0.11}_{-0.01}) \text{ fm.} \quad (9)$$

Although the $N\Omega$ is attractive everywhere, the binding energy is as small as $\sim 1 \text{ MeV}$ due to the short range nature of the potential. Accordingly, the root mean square distance is comparable to the scattering length, indicating that the system is loosely bound like the deuteron and the di-Omega.

In our pilot study [10], we found $B = 18.9(5.0)(^{+12.1}_{-1.8}) \text{ MeV}$ for heavy pion mass $m_\pi = 875 \text{ MeV}$. The larger magnitude of $B$ than the present result in Eq. (9) originates partly from the heavy masses of $N$ and $\Omega$ in [10] which

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The values in the fm unit are $(a_0, r_{\text{eff}})_{NN(^3S_1)} = (5.4112(15), 1.7463(19))$, $(a_0, r_{\text{eff}})_{NN(^1S_0)} = (-23.7148(43), 2.750(18))$ from the experiment [31], and $(a_0, r_{\text{eff}})_{\Omega\Omega(^1S_0)} = (4.6(6)(^{+1.2}_{-0.5}), 1.27(3)(^{+0.08}_{-0.03}))$ from the lattice QCD calculation [7].
reduce the kinetic energy and thus increase the binding energy. Another reason is that the short-range attraction for heavy pion is relatively stronger.

So far, we have not considered extra attraction in the $p\Omega^-$ system due to Coulomb attraction. By taking into account the correction $V_C(r) \to V_C(r) - \alpha/r$ with $\alpha \equiv e^2/(4\pi) = 1/137.036$, we obtain the observables,

$$B_{p\Omega^-} = 2.46(0.34)(^{+0.04}_{-0.11}) \text{ MeV, } \sqrt{\langle r^2 \rangle}_{p\Omega^-} = 3.24(0.19)(^{+0.06}_{-0.00}) \text{ fm.}$$

These results for $p\Omega^-(^5S_2)$ are summarized in Fig. 6 together with $n\Omega^-(^3S_2)$ without Coulomb correction.

Before ending this section, let us briefly discuss other possible systematic errors in Eqs. (8), (9) and (10). The first one is the finite volume effect whose typical error would be $\exp(-2m_\pi(L/2)) \approx \exp(-6) \approx 0.25\%$ and is much smaller than the statistical errors in our simulation. The second one is the finite cutoff effect, which is also expected to be small assuming the naive order estimate $(\Lambda a)^2 \leq 1\%$ with the non-perturbative $O(a)$ improvement. The third systematic error is due to the slightly heavy hadron masses ($m_\pi = 146$ MeV, $m_N = 955$ MeV and $m_\Omega = 1712$ MeV). By using the
same parameter set for $t/a = 12$ in Table 1 with $m_\pi = 146$ MeV kept fixed but with physical baryon masses ($m_p = 938$ MeV and $m_\Omega^- = 1672$ MeV), we find less binding than Eq. (10) as expected: $B_{p\Omega^-} \simeq 2.18(32)$ MeV and $\sqrt{\langle r^2 \rangle_{p\Omega^-}} \simeq 3.45(22)$ fm. On the other hand, if we additionally employ $m_\pi^+ = 140$ MeV for the potential (see Eq. (6)), we find more binding than Eq. (10) due to smaller pion mass: $B_{p\Omega^-} \simeq 3.00(39)$ MeV and $\sqrt{\langle r^2 \rangle_{p\Omega^-}} \simeq 3.01(16)$ fm.

5. Summary

In this paper, we have studied the $N$-$\Omega$ system in the $^5S_2$ channel, which is one of the promising candidates for quasi-stable dibaryon, from the (2+1)-flavor lattice QCD simulations with nearly physical quark masses ($m_\pi \simeq 146$ MeV and $m_K \simeq 525$ MeV). The $N$-$\Omega$ central potential in the $^5S_2$ channel obtained by the time-dependent HAL QCD method is found to be attractive in all distances. The scattering length and the effective range obtained by solving the Schrödinger equation using the resultant potential show that $N\Omega(^5S_2)$ is close to unitarity similar to the cases of the deuteron ($pn$) and di-Omega ($\Omega\Omega$). The binding energy of $p\Omega^-$ without (with) the Coulomb attraction is about 1.5 MeV (2.5 MeV), which indicates the existence of a shallow quasi-bound state below the $N\Omega$ threshold. In our simulation, we did not find a signature of the strong coupling between $N\Omega(^5S_2)$ and $\Lambda\Xi$ or $\Sigma\Xi$ in the D-wave state.

The $N\Omega(^5S_2)$ in the unitary regime can be studied in the two-particle correlation measurements in $p$-$p$ and $p$-nucleus and nucleus-nucleus collisions as suggested theoretically in [12] and experimentally reported by the STAR Collaboration at RHIC [16]. Phenomenological analyses along this line on the basis of the results in the present paper will be reported elsewhere [32].

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