Exotic dibaryons with a heavy antiquark

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Abstract

We discuss the possible existence of exotic dibaryons with a heavy antiquark, being realized as three-body systems, $\bar{D}^{(*)}NN$ and $B^{(*)}NN$. These are genuinely exotic states with no quark-antiquark annihilation. We consider the heavy quark spin and chiral symmetries, and introduce the one pion exchange potential between a $\bar{D}^{(*)}$ ($B^{(*)}$) meson and a nucleon $N$. As for the $NN$ interaction, we employ the Argonne $v'_{8}$ potential. By solving the coupled-channel equations for $PNN$ and $P^*NN$ ($P^{(*)} = \bar{D}^{(*)}$ and $B^{(*)}$), we find bound states for $(I, J^{P}) = (1/2, 0^-)$ as well as resonant states for $(I, J^{P}) = (1/2, 1^-)$ both in $\bar{D}^{(*)}NN$ and $B^{(*)}NN$ systems. We also discuss the heavy quark limit, and find that the spin degeneracy is realized in the bound states with $(I, J^{P}) = (1/2, 0^-)$ and $(1/2, 1^-)$.

Keywords: Mesic nuclei, Heavy mesons, Heavy quark symmetry, One pion exchange potential

1. Introduction

Many exotic hadrons observed in experiments are considered to be loosely bound states of two hadrons, called hadronic molecules or hadronic composites [1]. The study of such configurations is important to establish interactions among hadrons as inputs of various hadronic phenomena, such as formation of bound/resonant states and decay processes in hadronic molecules.
Moreover, the hadron interactions provide us with useful information to study fundamental problems of QCD such as color confinement, chiral symmetry breaking and so on.

The hadron-nucleon interaction is the basic quantity, not only for hadronic molecules, but also for exotic nuclei. In fact, the hyperon-nucleon interaction determines the properties of a variety of hypernuclei \[2–4\]. Furthermore it is suggested by theoretical studies \[5–8\] that the attraction between a \( \bar{K} \) meson and a nucleon would lead to the formation of \( \bar{K} \)-mesic nuclei, which are also researched in experiments \[9, 10\]. In the heavy quark sector, there have been, not only analogous discussions, but also new approaches which are not accessible in light flavor sectors. There the heavy quark symmetry \[11–13\] becomes important. Under this symmetry, an interesting observation was made, namely there is a sufficiently strong attraction due to the tensor force of the one pion exchange at long distances between a \( \bar{D}^{(*)} (B^{(*)}) \) meson and a nucleon \( N \) leading to the \( \bar{D}^{(*)} N (B^{(*)}N) \) molecules around the thresholds \[14–18\]. Here, we note that \( \bar{D}^{(*)} \) stands for \( \bar{D} \) or \( \bar{D}^* \). A unique feature of such molecules is the exotic flavor structure of the minimum quark content \( \bar{Q}qqqq \), where \( \bar{Q} \) is a heavy antiquark and \( q \)'s are light quarks. Hence, the \( \bar{D}^{(*)} N (B^{(*)}N) \) molecules are genuinely exotic baryons, having no lower hadronic channels coupled by a strong decay.

In literatures, however, there are other theoretical studies suggesting that repulsions at short distances and weaker pion coupling may change the interaction less attractive or even repulsive, which makes no bound state \[19–22\]. Our strategy here is to determine the long range part of the interaction, namely the one pion exchange potential with the strong tensor force, with respecting the heavy quark symmetry.

The possible attraction between a \( \bar{D}^{(*)} (B^{(*)}) \) meson and a nucleon motivates us to explore the few-body problems in exotic nuclei with heavy quarks, because it is naturally expected that the binding energy becomes larger as the baryon number increases. In the light flavor sector, it has been studied that the hadron-nucleon interaction gives us rich phenomena in few-body systems such as the impurity effects, e.g. shrinking of the wave functions due to glue-like effects in hypernuclei \[2\] and possible high density states in \( \bar{K} \)-mesic nuclei \[6\], which have never been realized in normal nuclei. In the heavy flavor sector, however, few-body systems of \( \bar{D}^{(*)} (B^{(*)}) \) nuclei with a few baryon numbers have not been investigated so far in the literature, though there have been several works for \( \bar{D}^{(*)} (B^{(*)}) \) mesons in nuclear matter with infinite volume \[23–34\] (see Ref. \[33\] for a summary of recent results).
and in atomic nuclei such as $^{12}$C, $^{40}$Ca and $^{208}$Pb with larger baryon numbers \[23, 32, 33\]. The few-body systems would be more likely to be produced in experiments in hadron colliders rather than the nuclei with middle and large baryon numbers.

In this paper, we study the mass spectrum of $\bar{D}^*(N)N$ and $B^*(N)N$ bound and/or resonant states of dibaryons (baryon number two) with a heavy antiquark for the first time. We emphasize that $\bar{D}^*(N)N$ and $B^*(N)N$ are unique as few-body systems, because the counterparts of $KNN$ in the strangeness sector do not exist due to the repulsive interaction between a $K$ meson and a nucleon. Moreover, we can reveal the role of the heavy quark spin symmetry, as a fundamental symmetry of QCD, in multi-hadron systems with a heavy quark, which does not appear in light flavor systems.

This paper is organized as follows. In the next section, we briefly summarize the interaction of $\bar{D}^*(N)N$ and $B^*(N)N$, where we emphasize the role of the one pion exchange potential which has been investigated in Refs. \[14–16\]. In Sec. 3 we show the method to solve the $\bar{D}^*(N)N$ and $B^*(N)N$ systems with appropriate three-body wave functions. In Sec. 4 the numerical results are shown, where we study bound and resonant states both for $\bar{D}^*(N)N$ and $B^*(N)N$ with quantum numbers $J^P = 0^-$ and $1^-$, and $I = 1/2$. We also discuss the spin degeneracy of the three-body systems in the heavy quark mass limit. In the last section, we summarize the present work.

2. Interactions

Let us start with the discussions of the basic interaction for the $P(N)$ ($P(N) = \bar{D}(N)$ and $B(N)$). As we have emphasized in Refs. \[14–16\] for the system with a heavy meson with a heavy antiquark, the one pion exchange potential (OPEP) is the basic ingredient to provide a strong attraction. The existence of the OPEP is a robust consequence of chiral symmetry in the presence of a light quark in the $P(N)$ meson, while its importance in the $P(N)$ interaction is supported by the heavy quark symmetry of QCD \[11–13\]. The relevance of the heavy quark symmetry is in the spin degeneracy of $P$ and $P^*$ mesons. It is related to the mass degeneracy between $P$ and $P^*$ as experimental data shows the small mass differences, $m_{\bar{D}^*} - m_{\bar{D}} \sim 140$ MeV and $m_{B^*} - m_B \sim 45$ MeV. Thanks to these small mass splittings, $P$ and $P^*$ mesons mix dynamically through the couplings of $PN - P^*N$ and $P^*N - P^*N$ caused by the OPEP. We note that the OPEP does not exist in $PN - PN$ only, because the $PP\pi$ vertex cannot exist from the parity.
conservation. The interaction for $PN - PN$ is effectively supplied from the mixing processes like $PN \rightarrow P^*N \rightarrow PN$. Thus, the OPEP together with the heavy quark spin symmetry provides a unique dynamics in the systems of a heavy meson.

The OPEP between a $P^{(s)}$ meson and a nucleon $N$ can be described from the heavy meson effective theory. The OPEPs for $PN - P^*N$ and $P^*N - P^*N$ are given by

$$V_{PN-P^*N}(r) = -\frac{g_\pi g_{\pi NN}}{\sqrt{2}m_N f_\pi} \frac{1}{3} \left[ \varepsilon^\dagger \cdot \vec{\sigma} C(r) + S_\epsilon T(r) \right] \vec{T}_P \cdot \vec{T}_N, \quad (1)$$

$$V_{P^*N-P^*N}(r) = \frac{g_\pi g_{\pi NN}}{\sqrt{2}m_N f_\pi} \frac{1}{3} \left[ \vec{S} \cdot \vec{\sigma} C(r) + S_\epsilon T(r) \right] \vec{T}_P \cdot \vec{T}_N, \quad (2)$$

as a sum of the central and tensor forces, $C(r)$ and $T(r)$, where $m_N = 940$ MeV and $f_\pi = 132$ MeV are the mass of the nucleon and the pion decay constant, respectively. In Eqs. (1) and (2), $\varepsilon (\varepsilon^\dagger)$ is the polarization vector of the incoming (outgoing) $P^*$, $\vec{S}$ is the spin-one operator of $P^*$, and $S_\epsilon (S_\epsilon)$ is the tensor operator $S_\epsilon (S) = 3(\vec{\sigma} \cdot \vec{r})(\vec{\sigma} \cdot \vec{r}) - (\vec{\sigma} \cdot \vec{r}) (\vec{\sigma} \cdot \vec{r})$ with $\vec{r} = \vec{r}/r$ and $r = |\vec{r}|$ for $\vec{S} = \varepsilon (\varepsilon^\dagger)$, where $\varepsilon$ is the relative position vector between $P^{(s)}$ and $N$. $\vec{\sigma}$ are Pauli matrices acting on nucleon spin. $\vec{T}_P (\vec{T}_N)$ are isospin operators for $P^{(s)} (N)$. The coupling constant for $P^{(s)}P^*\pi$ vertex is given by $g_\pi = 0.59$ which is determined by the experimental value of the decay width of $D^* \rightarrow D\pi$. Here we use $g_\pi$ with the same value for Eqs. (1) and (2) from the heavy quark spin symmetry. The coupling constant $g_{\pi NN}/4\pi = 13.6$ for $NN\pi$ vertex is given in Ref. [35]. The functions $C(r)$ and $T(r)$ are

$$C(r) = \int \frac{d^3 \vec{q}}{(2\pi)^3} \frac{m_\pi^2}{\vec{q}^2 + m_\pi^2} e^{i\vec{q} \cdot \vec{r}} F(\vec{q}), \quad (3)$$

$$S_\epsilon (\vec{r}) T(r) = \int \frac{d^3 \vec{q}}{(2\pi)^3} \frac{-\vec{q}^2}{\vec{q}^2 + m_\pi^2} S_\epsilon (\vec{q}) e^{i\vec{q} \cdot \vec{r}} F(\vec{q}), \quad (4)$$

with $\vec{q} = \vec{q}/|\vec{q}|$, where the dipole-type form factor $F(\vec{q}) = (\Lambda_D^2 - m_\pi^2)/(\Lambda_D^2 - |\vec{q}|^2)(\Lambda_N^2 + |\vec{q}|^2)$ with cutoff parameters $\Lambda_D$ and $\Lambda_N$ is introduced for spatially extended hadrons. From a quark model estimation, we use $\Lambda_D = 1.35\Lambda_N$ for $D^{(*)}$ meson, $\Lambda_B = 1.29\Lambda_N$ for $B^{(*)}$ meson, and $\Lambda_{PQ} = 1.12\Lambda_N$ for a $P^{(*)}$ meson defined as a meson ($\vec{Q}q_{\text{spin,0}(1)}$) having an infinitely heavy quark mass, as discussed in Refs. [14,16]. The cutoff $\Lambda_N = 830$ MeV is determined to reproduce the binding energy of the deuteron, 2.22 MeV.
when only the OPEP is used as the \( NN \) potential. As we discussed in Ref. \[15\] and further verified, the resulting OPEP provides a reasonable one-parameter approximation to the low energy properties of the \( NN \) system in the \( ^3S_1 - ^3D_1 \) channel, including the deuteron, as summarized in Table 1.

Table 1: Low energy properties of the \( NN \) system, the binding energy \( E_B \), relative distance of proton and neutron \( \langle r^2 \rangle^{1/2} \) and quadrupole moment \( Q_d \) of the deuteron, and \( NN \) scattering length \( a \) and effective range \( r_e \). The predictions obtained by the OPEP \[15\] and AV8', and experimental values (Exp.) summarized in Ref. \[36\] are compared.

|       | \( E_B \) [MeV] | \( \langle r^2 \rangle^{1/2} \) [fm] | \( Q_d \) [fm] | \( a \) [fm] | \( r_e \) [fm] |
|-------|-----------------|---------------------------------|----------------|---------|-------------|
| OPEP  | 2.22            | 3.7                             | 0.24           | 5.27    | 1.50        |
| AV8'  | 2.23            | 3.9                             | 0.27           | 5.39    | 1.75        |
| Exp.  | 2.22            | 3.9                             | 0.29           | 5.42    | 1.76        |

We note that, although the OPEP has been known to play an essential role in the nucleon-nucleon interaction, attention has not been paid in the meson-nucleon interaction. In fact, the mixing effect is much suppressed in the light meson sector like \((\pi, \rho)\) and \((K, K^*)\), where the mass splittings are much larger than in the heavy meson sector, \( m_\rho - m_\pi \sim 600 \text{ MeV} \) and \( m_{K^*} - m_K \sim 400 \text{ MeV} \), and therefore the OPEP plays only a minor role. We expect that the nuclear systems containing \( P^{(*)} \) mesons, in accordance with the heavy quark spin symmetry, revive the importance of the role of the OPEP in the hadron-nuclear systems.

In the present work, we employ the OPEP for the interaction of \( P^{(*)}N \), while we could have a short-range interaction as provided by the \( \rho \) and \( \omega \) meson exchanges, or even by the quark exchange model \[18, 22, 37, 38\]. Although these interactions are less determined than the OPEP, it turns out that they are almost irrelevant for \( P^{(*)}N \) and \( P^{(*)}NN \) systems. In fact, we have numerically verified in the previous study that the two-body matrix elements of the OPEP are much larger than those of the short-range interaction of \( \rho \) and \( \omega \) exchanges; the latter are typically less than ten percent \[15\] of the former, an analogous situation in the deuteron. This argument justifies

\footnote{In Ref. \[15\], the quantity \( \langle r^2 \rangle^{1/2} \) should have been called a relative distance as we do in this article, rather than the radius of the deuteron. Besides this misleadingness, actual calculations were done properly.}
the use of the OPEP for the $P^{(*)}N$.

For the nucleon-nucleon interaction, though the OPEP is reasonably good for the description of low energy properties, we employ a more realistic potential of Argonne $v'_{8}$ (AV8′) \(^{39}\). It is written as

\[
v'_{8}(r) = \sum_{p=1}^{8} v_{p}(r)O^{p},
\]

formed by a sum of 8 operators $O^{p=1,\ldots,8}$; the central operators $[1, (\vec{\sigma}_{1} \cdot \vec{\sigma}_{2})] \otimes [1, (\vec{\tau}_{1} \cdot \vec{\tau}_{2})]$, the tensor operators $S_{12} \otimes [1, (\vec{\tau}_{1} \cdot \vec{\tau}_{2})]$, and the $LS$ operators $(\vec{L} \cdot \vec{S}) \otimes [1, (\vec{\tau}_{1} \cdot \vec{\tau}_{2})]$. The function $v_{p}(r)$ is given in Ref. \(^{39}\). The AV8′ potential is simpler than the more elaborated one of the Argonne $v_{18}$ (AV18) potential \(^{36}\), which is the reason that we employ the former in the present study. The AV8′ is realistic because it reproduces $NN$ phase shifts and deuteron properties. When applied to three or more nucleon systems, however, the AV8′ provides slightly more attraction than the AV18 potential. In the $\bar{D}^{(*)}NN$ systems, however, there are only two nucleons with the binding energy similar to that of the deuteron, and therefore, the AV8′ gives essentially the same results as the AV18.

3. Three-body wave functions

The total Hamiltonian is given by

\[
H = T + V_{P^{(*)}N} + V_{NN},
\]

where $T$ is the kinetic term, and $V_{P^{(*)}N}$ ($V_{NN}$) is the interaction potential between a $P^{(*)}$ meson and a nucleon (between two nucleons) as introduced above. We note that the heavy quark spin symmetry is respected in the $V_{P^{(*)}N}$ potential. The small violation of the heavy quark spin symmetry is taken into account in the mass splitting between $P$ and $P^{*}$. We investigate $\bar{D}^{(*)}NN$ and $B^{(*)}NN$ with $J^{P} = 0^{-}$ and $1^{-}$ and $I = 1/2$ (total angular momentum $J$, parity $P$ and total isospin $I$). We also consider $P^{(*)}NN$ in the heavy quark limit.

In order to express the three-body wave functions, we employ the Gaussian expansion method \(^{40}\) which is one of the powerful methods to solve few-body calculations and has been utilized to investigate bound and scattering states of such systems in hadron and nuclear physics. In the study, we do
not solve the Faddeev equation which gives accurate solutions for the three-body systems. However, it has been discussed that the Gaussian method works well and that their results are equivalent to those by the Faddeev method, when the good convergence of the eigenenergy is obtained [40].

The three-body wave function is described as a sum of the rearrange channel amplitudes \((c = 1, 2)\) as functions of the Jacobi coordinates shown in Fig. 1:

\[
\Psi_{JM} = \sum_{c=1}^{2} \sum_{n l_1, N l_2, L, s_{12} S, l_{12} I_{12}} C_{nl_1, N l_2, L, s_{12} S, l_{12} I_{12}}^{(c)} A \left\{ \left[ \phi_{nl_1 m_1}^{(c)}(\vec{r}_c) \psi_{NLm_2}^{(c)}(\vec{R}_c) \right]_{L} \times \left[ [\chi_{s_1} \chi_{s_2}]_{s_{12}} \chi_{s_3} \right]_{S} \right\}_{JM} \left[ [\eta_I \eta_I]_{I_{12}} \eta_I \right]_{I}. \tag{7}
\]

\(A\) is the anti-symmetrization operator for exchange between two nucleons. \(l_1\) and \(l_2\) stand for the relative orbital angular momenta associated with the coordinates \(\vec{r}_c\) and \(\vec{R}_c\), respectively. \(L\) is the total orbital angular momentum of the three-body system. \(\chi_{s_i} (\eta_i)\) with \(i = 1, 2, 3\) is the spin (isospin) function of the particle with the spin \(s_i\) (isospin \(I_i\)). \(s_{12} (I_{12})\) is the spin (isospin) of two particles combined by the relative coordinate \(\vec{r}_c\), and \(S\) is the total spin of the three-body system. The functions \(\phi_{nl_{1} m_{1}}(\vec{r})\) and \(\psi_{NLm_{2}}(\vec{R})\) are expressed in terms of the Gaussian functions [40] as

\[
\phi_{nl_{1} m_{1}}(\vec{r}) = \sqrt{\frac{2}{\Gamma(l_1 + 3/2)b_n^3}} \left( \frac{r}{b_n} \right)^{l_1} \exp \left( -\frac{r^2}{2b_n^2} \right) Y_{l_1 m_1}(\hat{r}), \tag{8}
\]

\[
\psi_{NLm_{2}}(\vec{R}) = \sqrt{\frac{2}{\Gamma(l_2 + 3/2)B_N^3}} \left( \frac{R}{B_N} \right)^{l_2} \exp \left( -\frac{R^2}{2B_N^2} \right) Y_{l_2 m_2}(\hat{R}). \tag{9}
\]

The Gaussian ranges \(b_n\) and \(B_N\) are given by the form of geometric series as

\[
b_n = b_1 a^{n-1}, \quad B_N = B_1 A^{N-1}. \tag{10}
\]

For the sum of Eq. (7), we include all possible coupled channels to obtain solutions with sufficiently good accuracy. For instance, we include orbital angular momentum of \(l_1, l_2 \leq 2\). Furthermore, we consider two independent isospin states to form the total isospin \(I = 1/2\). For instance, we include the NN subsystems of \(I = 0\) and \(1\) which are combined with the \(\bar{D}^{(*)}\) meson of \(I = 1/2\) for the total \(I = 1/2\).
Figure 1: Jacobi coordinates of the $P^{(*)}NN$ systems.

By diagonalizing the total Hamiltonian using the three-body bases introduced above, we obtain eigenenergies and coefficients $C_{nl_1,Nl_2,S,l_12,l_2}^{(c)}$. We also calculate the poles for resonances as complex eigenvalues by using the complex scaling method \cite{41–44}.

4. Numerical Results

Let us present the results of $\bar{D}^{(*)}NN$ and $B^{(*)}NN$ for $J^P = 0^-$. We obtain bound states both of $\bar{D}^{(*)}NN$ and $B^{(*)}NN$ with energy levels shown in Fig. 2. The bound state of $\bar{D}^{(*)}NN$, whose binding energy is $-5.2$ MeV, locates below the threshold of $\bar{D}N(1/2^-) + N$. Here $\bar{D}N(1/2^-)$ is the bound state of $\bar{D}^{(*)}$ and $N$ with binding energy $-1.6$ MeV for $J^P = 1/2^-$ and $I = 0$ as discussed in Refs. \cite{14, 15}. Therefore, the three-body state of $\bar{D}^{(*)}NN$ is more bound than the two-body state of $\bar{D}^{(*)}N$, as naturally expected. We also find the $B^{(*)}NN$ state with the binding energy $-26.2$ MeV. The $B^{(*)}NN$ state is more bound than the $\bar{D}^{(*)}NN$ state, because the mixing effect between $P^NN$ and $P^{*}NN$ is enhanced, when $P$ and $P^*$ mesons become more degenerate.

Let us investigate how the bound states are formed. For this purpose, we analyze various components of interaction matrix elements. In Table 2, we summarize the expectation values of the potentials, $V_{P_{N-P^*N}}$, $V_{P^*_{N-P^*N}}$ and $V_{NN}$, sandwiched by the obtained wave functions. In $\bar{D}^{(*)}NN$, we find that the tensor force of $V_{\bar{D}N-D^*N}$ mixing $\bar{D}NN$ and $D^*NN$ is the dominant contribution. In contrast, $V_{\bar{D}^*N-D^*N}$ is very small. Thus, the tensor force of $V_{\bar{D}N-D^*N}$ is a driving force giving bound states in $\bar{D}^{(*)}NN$. We note that the strong tensor force also provides a dominant attraction in the two-body $\bar{D}^{(*)}N$ systems \cite{14, 15}. The same result is also applied to $B^{(*)}NN$. 

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Figure 2: Energy levels of $\bar{D}^{(*)}NN$, $B^{(*)}NN$ and $P_{Q}^{(*)}NN$ with $I = 1/2$ and $J^P = 0^-$ and $1^-$ (solid lines). The complex energies for resonances are given as $E_{re} - i\Gamma/2$, where $E_{re}$ is a resonance energy and $\Gamma/2$ is a half decay width. Thresholds (subthresholds) are denoted by dashed (dashed-dotted) lines.

For $V_{NN}$, one may expect that the tensor force causing the $^3S_1 - ^3D_1$ mixing could be the most dominant one, because it is a driving force giving a deuteron $d$. However, the tensor force in $V_{NN}$ is almost irrelevant in the present systems. In fact, it is shown in Table 2 that the tensor force is suppressed, while the central force is rather dominant. This is reasonable because $d$ does not exist in the main component of $\bar{D}NN$ due to limited combinations of quantum numbers. It may exist in the $\bar{D}^*NN$ component, but the amplitude of the $\bar{D}^*NN$ is small due to the excess of mass of about 140 MeV ($\sim m_{D^*} - m_D$). Thus, the $NN$ interaction provides only a weak attraction. Therefore, the tensor force mixing $\bar{D}N$ and $\bar{D}^*N$ gives the most dominant term of attraction in the $\bar{D}^{(*)}NN$ state. The same behavior is also found in the $B^{(*)}NN$ state as shown in Table 2.

The quantum number $0^-$ of the $\bar{D}^{(*)}NN$ state may be investigated in experiments through two particle correlations. The $D^{(*)}NN$ state can decay into $K + \pi$ (or $\pi\pi$) + $N + N$ by the weak decay of a $D$ meson. The absence of $d$ in the final $NN$ state will be an important signal suggesting $0^-$. In
Table 2: Expectation values of central, tensor and LS forces of the $\bar{D}^{(*)}N (B^{(*)}N)$ and $NN$ potentials in the bound state of $\bar{D}^{(*)}NN (B^{(*)}NN)$. All values are in units of MeV.

| $\bar{D}^{(*)}NN$ | $\langle V_{\bar{D}N-N\bar{N}} \rangle$ | $\langle V_{\bar{D}N-N\bar{N}} \rangle$ | $\langle V_{NN} \rangle$ |
|-------------------|--------------------------------|--------------------------------|----------------|
| Central           | −2.3                          | −0.1                          | −9.5           |
| Tensor            | −47.1                         | 0.7                           | −0.2           |
| LS                | −                              | −                              | −0.03          |

| $B^{(*)}NN$ | $\langle V_{BN-B\bar{N}} \rangle$ | $\langle V_{B\bar{N}-B\bar{N}} \rangle$ | $\langle V_{NN} \rangle$ |
|-------------|--------------------------------|--------------------------------|----------------|
| Central     | −6.5                          | 0.3                           | −11.6          |
| Tensor      | −92.0                         | −2.7                          | −1.0           |
| LS          | −                              | −                              | −0.1           |

contrast, if $d$ might be observed, the quantum number would be $1^-$. It will be also the case for $B^{(*)}NN$.

In scattering states, we find resonances for $J^P = 1^-$ and $I = 1/2$ as shown in Fig. 2. The resonance energy for $\bar{D}^{(*)}NN$ is 111.2 MeV measured from the threshold of $\bar{D}NN$. The decay width is 18.6 MeV. We note that there are open channels of the $\bar{D}NN$ and $\bar{D} + d$ scattering states below the resonance, and of the $\bar{D}^* + d$ and $\bar{D}N(3/2^-) + N$ scattering states above the resonance. Here $\bar{D}N(3/2^-)$ is a Feshbach resonance of $\bar{D}^{(*)}$ and $N$ with $J^P = 3/2^-$ and $I = 0$, which was found in Ref. [15]. Those scattering states are included in the present calculation. We obtain a resonance also for $B^{(*)}NN$ with much smaller resonance energy and decay width, 6.8 MeV and 0.4 MeV, respectively. The mechanism of formation of the resonances is interesting. When we ignore the $\bar{D}NN$ channel and consider only the $\bar{D}^*NN$ channel, we obtain a bound state of $\bar{D}^*NN$. Hence, the $\bar{D}^*NN$ channel predominates. Therefore, the obtained resonance is a Feshbach resonance for the three-body system, as in the case of the two-body $\bar{D}N(3/2^-)$ system [15]. These features also hold for $B^{(*)}NN$.

From the above analysis, we see that the many features of the two-body $P^{(*)}N$ system survive in the three-body $P^{(*)}NN$ system, because the $P^{(*)}N$ interaction is the dominant force which determines the main properties of the system rather than the $NN$ interaction.

Finally, we consider $P_Q^{(*)}NN$ systems in the heavy quark limit, where $P_Q$ and $P_Q^*$ are exactly degenerate in mass. Interestingly, we find the bound states both for $J^P = 0^-$ and $1^-$ with the same binding energy $−38.5$ MeV.
measured from $P_Q^{(s)}NN$ threshold. Those numerical results indicate that they are degenerate. Comparing three cases of $\bar{c}$, $\bar{b}$ and $\bar{Q}$ with $m_Q \to \infty$, we see that the mass splitting between the three-body systems with $J^P = 0^−$ and $1^−$ decreases as the mass of the heavy quark increases, and finally those two states become degenerate in the heavy quark limit, as shown in Fig. 2.

The degeneracy in $P_Q^{(s)}NN$ systems is not accidental. Generally, the spin degeneracy is a feature in the heavy quark limit, not only for a single hadron, but also for multi-hadrons, when a single heavy quark is contained. In QCD, the spin-dependent interaction of the heavy quark is suppressed by the inverse of the heavy quark mass. Thus, there should exist spin doublets (singlets) forming (non-)degenerate states with $j \neq 0$ ($j = 0$) in the heavy quark limit. Here $j$ is the quantum number of the contained light components (the brown muck [13] or the spin-complex [15]). The spin degeneracy is shown, not only from the heavy quark effective theory [46, 47], but also from the heavy hadron effective theory [15]. In the present discussion from the heavy meson effective theory (Eqs. (1) and (2)), we have confirmed that the ground state of three-body $P_Q^{(s)}NN$ systems exhibits the spin doublet having $0^−$ and $1^−$. Both of them contain a common light component with total angular momentum and parity $j^P = 1/2^+$ and isospin $I = 1/2$.

5. Summary

In this paper, we have explored the possible existence of genuinely exotic dibaryons, $\bar{D}^{(s)}NN$ and $\bar{B}^{(s)}NN$ for charm and bottom sectors, and $P_Q^{(s)}NN$ in the heavy quark limit. The OPEP introduced by chiral symmetry and the heavy quark spin symmetry was employed between a heavy meson and a nucleon. As for the $NN$ interaction, we employed the Argonne $v_8'$ potential. By solving the coupled-channel equations for $PN$ and $P^*NN$, we have obtained bound states with $J^P = 0^−$ and Feshbach resonances with $J^P = 1^−$ for $I = 1/2$ both in $\bar{D}^{(s)}NN$ and $\bar{B}^{(s)}NN$. The tensor force of the OPEP mixing $PN$ and $P^*N$ plays an important role to produce a strong attraction. For the $P_Q^{(s)}NN$ systems in the heavy quark limit, we have obtained bound states of both $J^P = 0^−$ and $1^−$, which are the spin degenerate states containing the brown muck or the spin-complex with $j^P = 1/2^+$ and $I = 1/2$. This is the first study to show the degeneracy in the few-body system with a heavy antiquark. Hence the bound states with $J^P = 0^−$ and the resonance with $J^P = 1^−$ in the actual charm and bottom sectors have the common origin of the spin doublet from the heavy quark limit.
The $\bar{D}(^{(*)}NN$ and $B(^{(*)}NN$ states can be searched in experiments in hadron colliders. The productions of the exotic hadrons will be studied in relativistic heavy ion collisions in RHIC and LHC [48]. Furthermore, the search for the $\bar{D}(^{(*)}NN$ would be also carried out in J-PARC and GSI-FAIR.

Acknowledgments

The authors would like to thank Dr. Y. Kikuchi for valuable discussions and fruitful suggestions. This work is supported in part by Grant-in-Aid for Scientific Research on Priority Areas “Elucidation of New Hadrons with a Variety of Flavors(E01: 21105006)” (S. Y. and A. H.) from the Ministry of Education, Culture, Sports, Science and Technology of Japan, and Grant-in-Aid for “JSPS Fellows(24-3518)” (Y. Y.) from Japan Society for the Promotion of Science.

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