Physics Letters B 748 (2015) 463–466

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Doubly heavy baryons with chiral partner structure

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ARTICLE INFO

Article history:
Received 18 April 2015
Received in revised form 9 July 2015
Accepted 18 July 2015
Available online 22 July 2015
Editor: B. Grinstein

ABSTRACT

The spectrum and dominant strong decay properties of the doubly heavy baryons are revisited by using a chiral effective model with the chiral partner structure. By regarding the doubly heavy baryons in the ground states and light angular momentum, $j_l = 1/2$, sector of the first orbitally excited states as chiral partners, we estimate their mass splitting arising from the spontaneous breaking of chiral symmetry to be about 430 MeV for baryons including an unflavored light quark and about 350 MeV for that including a strange quark. We point out that, similar to the heavy–light meson sector, the intermultiplet decay from a baryon with negative parity to its chiral partner and a pion is determined by the mass splitting through the generalized Goldberger–Treiman relation. Furthermore, the isospin-violating decay of the Ω_{cc} baryon, $((1/2)^-, (3/2)^-)_s \to ((1/2)^+, (3/2)^+)_s + \pi^0$, through the η–π^0 mixing is the dominant decay channel of the doubly heavy baryons including a strange quark.

Heavy hadron spectroscopy has been drawing extensive interest since the last decade because a large amount of heavy hadrons are observed in particle colliders. It is reasonable to expect that the current and future scientific facilities such as LHCb and Belle II can observe more and heavier resonances, such as the doubly heavy baryons (DHBs) studied in this work.

The existence of DHBs is an immediate prediction of quantum chromodynamics (QCD). These baryons have been theoretically discussed for a long time [1–5]. Meanwhile, several experimental efforts have been made to detect such states [6–13], and the positive results from SELEX show that the mass of the doubly charmed baryons Ξ_{cc}^+ is about 3520 MeV [6–8]. In this paper, we investigate some properties of the DHBs based on the chiral partner structure. As a DHB contains only one light quark, unlike the light baryons, its chiral behavior is very simple.

The chiral partner structure of hadrons including a heavy quark has been studied by several groups. Nowak et al. in Ref. [14] and then Bardeen and Hill [15] proposed pioneering ideas in the heavy–light meson sector. In this situation, the heavy–light meson doublets in the heavy quark limit with quantum numbers $(1^-, 0^-)$ and $(1^+, 0^+)$ are regarded as chiral partners, and their mass splitting is induced by the dynamical breaking of the chiral symmetry so that the magnitude is approximately the constituent quark mass. This was confirmed by the spectrum of the relevant particles: $m_{D_0^*} - m_D \simeq m_{D_1} - m_{D^*} \simeq 450$ MeV is at the same order of $m_{D_{s0}}(2317) - m_{D_s} \simeq m_{D_{s1}}(2460) - m_{D_s^*} \simeq 350$ MeV (see, e.g., Refs. [16–18]). In the sector of heavy baryons including a heavy quark, the chiral partner structure is mainly accessed based on the bound-state approach [18]. However, in this sector, there are some disagreements about the chiral partner structure [19]. These disagreements might arise from the fact that the heavy baryons including one heavy quark contain two light quarks, which does not make their chiral properties quite as simple as the heavy–light meson sector and also the DHBs considered in this work.

Schematically, the quark contents of a DHB including the same heavy quarks can be written as QQq with Q and q being the heavy quark and light quark constituents, respectively. As the DHB is a colorless object, its two heavy quarks form an anti-color triplet [16]. Because the heavy quarks in the DHB have a large mass, requiring much larger energy to orbitally excite the heavy constituent quark than the light quark, it is reasonable to regard the two constituent quarks as a static compact object without orbital excitation. The constituent of the DHB can be denoted as $\bar{\mathbf{\Phi}}q$ with $\bar{\Phi}$ being the heavy quark component, which should be a bosonic quantity. Bearing such an intuitive scenario in mind, one can define the chiral partner structure similar to that in the heavy–light meson sector [14,15].

As the two heavy quarks in a DHB are antisymmetric in color space, they should have the total spin $J_Q = 1$ in s-wave; therefore, DHBs in the ground states can form a heavy quark doublet $D_{\mathbf{Q}}^{\mu}$

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http://dx.doi.org/10.1016/j.physletb.2015.07.046
0370-2693/

whose components have quantum numbers $J^P = \frac{1}{2}^+, \frac{3}{2}^+$. For the first orbital excitation with relative angular momentum between the light quark and heavy quark source $l = 1$, the light angular momentum could be $j_l = \frac{1}{2}, \frac{3}{2}$. Combing $j_l = \frac{1}{2}$, one can form another heavy quark doublet $N^\mu_{\mathbf{Q}}$ with quantum numbers $J^P = \frac{1}{2}^-, \frac{3}{2}^-$. We regard the doublets $D^\mu_{\mathbf{Q}}$ and $N^\mu_{\mathbf{Q}}$ as chiral partners, and DHBs constructed from $j_l = \frac{3}{2}$ can be regarded as chiral partners of some states from $l = 2$ baryons [16].

Similar to the heavy–light meson case, because the DHB contains only one light quark, the DHB doublets $D^\mu_{\mathbf{Q}}$ and $N^\mu_{\mathbf{Q}}$ can be written in the chiral basis by introducing the fields $D^\mu_{\mathbf{Q};L,R}$, which at the quark level are schematically written as $D^\mu_{\mathbf{Q};L,R} \sim \bar{\mathbf{\Phi}}^\mu q_{L,R}$. As the heavy quark component of the DHB is a boson, $D^\mu_{\mathbf{Q};L,R}$ should be Lorentz spinors and, under chiral transformation, transform as

$$D^\mu_{\mathbf{Q};L,R} \to g_{L,R} D^\mu_{\mathbf{Q};L,R}, \tag{1}$$

where $g_{L,R} \in SU(3)_{L,R}$. In terms of $D_{\mathbf{Q}}$ and $N_{\mathbf{Q}}$, one can write

$$D^\mu_{\mathbf{Q};L} = \frac{1}{\sqrt{2}}\left(D^\mu_{\mathbf{Q}} - iN^\mu_{\mathbf{Q}}\right), \quad D^\mu_{\mathbf{Q};R} = \frac{1}{\sqrt{2}}\left(D^\mu_{\mathbf{Q}} + iN^\mu_{\mathbf{Q}}\right), \tag{2}$$

which transform as $D^\mu_{\mathbf{Q};L,R} \leftrightarrow \gamma_0 D_{\mathbf{Q};\mu;R,L}$ under parity transformation and satisfy $\not{v} D^\mu_{\mathbf{Q};L,R} = D^\mu_{\mathbf{Q};L,R}$ and $v_\mu D^\mu_{\mathbf{Q};L,R} = 0$ for preserving the heavy quark symmetry and keeping the transversality. Further, for later convenience, following the procedure given in Ref. [20], the DHB doublets $D^\mu_{\mathbf{Q}}$ and $N^\mu_{\mathbf{Q}}$ in terms of the physical states are written as

$$D^\mu_{\mathbf{Q}} = \frac{1+\not{v}}{2}\Psi^\mu_{QQ} + \sqrt{\frac{1}{3}}\left(\gamma^\mu + v^\mu\right)\gamma^5 \frac{1+\not{v}}{2}\Psi_{QQ}, \quad N^\mu_{\mathbf{Q}} = \frac{1+\not{v}}{2}\Psi'^\mu_{QQ} + \sqrt{\frac{1}{3}}\left(\gamma^\mu + v^\mu\right)\gamma^5 \frac{1+\not{v}}{2}\Psi^*_{QQ}, \tag{3}$$

which is the same as that for the heavy baryons including one heavy quark [21,22], and $\Psi^{(\prime)\mu}_{QQ}$ is the spin-$\frac{3}{2}$ Rarita–Schwinger field. In Eq. (3), Ψ_{QQ}, Ψ^*_{QQ}, Ψ^μ_{QQ}, and Ψ'^μ_{QQ} denote the DHBs with spin parity $\frac{1}{2}^+, \frac{1}{2}^-, \frac{3}{2}^+$, and $\frac{3}{2}^-$, respectively. One can easily check that these spinors satisfy $\not{v} D^\mu_{\mathbf{Q}} = D^\mu_{\mathbf{Q}}$ and $\not{v} N^\mu_{\mathbf{Q}} = N^\mu_{\mathbf{Q}}$. The intrinsic parity behavior was imposed:

$$\mathrm{P}: \ \Psi^\mu_{QQ} \to -\gamma_0 \Psi_{QQ,\mu}, \ \ \Psi_{QQ} \to \gamma_0 \Psi_{QQ}, \quad \Psi'^\mu_{QQ} \to \gamma_0 \Psi'_{QQ,\mu}, \ \ \Psi^*_{QQ} \to -\gamma_0 \Psi^*_{QQ}. \tag{4}$$

When the heavy quark in the DHB is a c quark and the light quark is a u, d, or s quark, the DHB field, for example, Ψ_{QQ} stands for Ξ^{++}_{cc}, Ξ^{+}_{cc}, and Ω^{+}_{cc}, respectively.

Now, the chiral effective theory of DHBs can be constructed in the chiral basis. It is noted that the quark–diquark symmetry [23] relates the doubly heavy baryons with the heavy mesons having the same brown muck [24]. For relating the parameters based on the quark–diquark symmetry, an effective Lagrangian is first written for the heavy–light mesons with the chiral partner structure by introducing chiral fields $\mathcal{H}_{L,R}$ [14,15]. These chiral fields relate to the heavy–light meson doublets H and G with quantum numbers $(0^-, 1^-)$ and $(0^+, 1^+)$, respectively, through

$$\mathcal{H}_R = \frac{1}{\sqrt{2}}\left[G - iH\gamma_5\right], \quad \mathcal{H}_L = \frac{1}{\sqrt{2}}\left[G + iH\gamma_5\right], \tag{5}$$

where G and H are heavy–light meson fields with positive and negative parity, respectively. In terms of the physical states, they are expressed as follows:

$$H = \frac{1+\not{v}}{2}\left[D^{*\mu}\gamma_\mu + iD\gamma_5\right], \quad G = \frac{1+\not{v}}{2}\left[-D'^\mu_1\gamma_\mu\gamma_5 + D^*_0\right]. \tag{6}$$

It should be noticed that because the heavy component is a heavy quark and the light component is a light antiquark in the heavy–light meson fields, not the chiral fields $\mathcal{H}_{L,R}$ but their conjugates $\bar{\mathcal{H}}_{L,R} \equiv \gamma_0 \mathcal{H}_{L,R}\gamma_0$ transform as the chiral quark fields $q_{L,R}$ under chiral transformation, that is, the same as Eq. (1). Here, we consider only the terms that survive in the heavy quark limit, including the terms up to one derivative. For the light mesons, we consider the chiral field M, which transforms as $M \to g_L M g_R^\dagger$ under chiral transformation. The effective Lagrangian is written as [25,26]

$$\begin{aligned}\mathcal{L}_{\mathrm{M}} = {} & \mathrm{tr}\left[\mathcal{H}_L(iv\cdot\partial)\bar{\mathcal{H}}_L\right] + \mathrm{tr}[\mathcal{H}_R(iv\cdot\partial)\bar{\mathcal{H}}_R] \\ & - \Delta\,\mathrm{tr}\left[\mathcal{H}_L\bar{\mathcal{H}}_L + \mathcal{H}_R\bar{\mathcal{H}}_R\right] \\ & - \frac{1}{2}g_\pi\,\mathrm{tr}\left[\mathcal{H}_L M\bar{\mathcal{H}}_R + \mathcal{H}_R M^\dagger\bar{\mathcal{H}}_L\right] \\ & + i\frac{g_A}{f_\pi}\mathrm{tr}\left[\mathcal{H}_L\gamma_5\gamma^\mu\partial_\mu M\bar{\mathcal{H}}_R - \mathcal{H}_R\gamma_5\gamma^\mu\partial_\mu M^\dagger\bar{\mathcal{H}}_L\right],\end{aligned} \tag{7}$$

where Δ provides the mass shift to both G and H in the same direction. After a suitable choice of the potential sector of the light meson Lagrangian, which will not be specified here, the chiral symmetry in the Nambu–Goldstone phase can be realized. In such a case, after the spontaneous breaking of the chiral symmetry, the meson field M can be replaced by $\bar{M} + \tilde{M}$ with $\bar{M} = \mathrm{diag}(v, v, v_3)$ being the vacuum expectation value of the chiral field in the isospin limit, which corresponds to the quark condensate, and $\tilde{M}$ being the fluctuation fields. Then, this g_π term provides the mass difference between G and H as

$$\Delta M_i = m_{G,i} - m_{H,i} = g_\pi \bar{M}_{ii}, \tag{8}$$

where the sub-indices i stand for the light flavor with $i = 1, 2$, and 3 representing the u, d, and s quarks, respectively. Here, $v = f_\pi = 92.4$ MeV is used, so that $g_\pi = 4.65$ is obtained from $\Delta M_{u,d} = 430$ MeV. Note that the g_π term also gives the interaction for the pionic transition between G and H. The relation between these two quantities are known as the generalized Goldberger–Treiman relation [14,15]. On the other hand, the g_A term gives the interaction of the pionic transition within G or H. The value of g_A is determined from the experimental value of $D^* \to D + \pi$ decay as $g_A = 0.56$ (see, e.g., Ref. [25]).

Now, the effective Lagrangian for the doubly heavy baryons can be considered. As stated earlier, the quark–diquark symmetry relates the Lagrangian to the above Lagrangian for the heavy mesons. The resultant effective Lagrangian is expressed as

$$\begin{aligned}\mathcal{L}_{\mathrm{B}} = {} & \bar{D}^\mu_{\mathbf{Q};L} iv\cdot\partial D_{\mathbf{Q};\mu;L} + \bar{D}^\mu_{\mathbf{Q};R} iv\cdot\partial D_{\mathbf{Q};\mu;R} \\ & - \Delta\left(\bar{D}^\mu_{\mathbf{Q};L} D_{\mathbf{Q};\mu;L} + \bar{D}^\mu_{\mathbf{Q};R} D_{\mathbf{Q};\mu;R}\right) \\ & - \frac{1}{2}g_\pi\left(\bar{D}^\mu_{\mathbf{Q};L} M D_{\mathbf{Q};\mu;R} + \bar{D}^\mu_{\mathbf{Q};R} M^\dagger D_{\mathbf{Q};\mu;L}\right) \\ & + \frac{ig_A}{f_\pi}\left[\bar{D}^\mu_{\mathbf{Q};L}\gamma_5\gamma^\nu\partial_\nu M D_{\mathbf{Q};\mu;R}\right. \\ & \qquad \left. + \bar{D}^\mu_{\mathbf{Q};R}\gamma_5\gamma^\nu\partial_\nu M^\dagger D_{\mathbf{Q};\mu;L}\right].\end{aligned} \tag{9}$$

By substituting (2) into the Lagrangian (9) and considering the spontaneous chiral symmetry breaking, one obtains the Lagrangian

$$\begin{aligned}\mathcal{L}_{\mathrm{B}} = {} & \bar{D}^\mu_{\mathbf{Q}} iv\cdot\partial D_{\mathbf{Q};\mu} + \bar{N}^\mu_{\mathbf{Q}} iv\cdot\partial N_{\mathbf{Q};\mu} - \Delta\left(\bar{D}^\mu_{\mathbf{Q}} D_{\mathbf{Q};\mu} + \bar{N}^\mu_{\mathbf{Q}} N_{\mathbf{Q};\mu}\right) \\ & - \frac{1}{2}g_\pi\left(\bar{D}^\mu_{\mathbf{Q}}\bar{M} D_{\mathbf{Q};\mu} - \bar{N}^\mu_{\mathbf{Q}}\bar{M} N_{\mathbf{Q};\mu}\right)\end{aligned}$$

Table 1
Spectrum of the doubly charmed baryons and the partial widths of one-pion intermultiplet transitions. Here, $m_{\Xi_{cc}} = 3520$ MeV [6], $m_{\Omega_{cc}} = 3678$ MeV [5], and $m_{\pi^\pm} = m_{\pi^0} \simeq 140$ MeV are taken as input. Other partial widths of intermultiplet transitions can be obtained using the isospin relation.

Spectrum	Prediction (MeV)	Decay channel	Partial width (MeV)
$m_{\Xi^*_{cc}}$	3950	$\Xi^{*++}_{cc} \to \Xi^{++}_{cc} + \pi^0$	331
$m_{\Xi^\mu_{cc}}$	3625	$\Xi^{*++}_{cc} \to \Xi^{+}_{cc} + \pi^+$	662
$m_{\Xi'^\mu_{cc}}$	4055	$\Xi'^{++}_{cc} \to \Xi^{++\mu}_{cc} + \pi^0$	332
$m_{\Omega^*_{cc}}$	4028	$\Xi'^{++}_{cc} \to \Xi^{+\mu}_{cc} + \pi^+$	663
$m_{\Omega^\mu_{cc}}$	3783	$\Omega^{*+}_{cc} \to \Omega^{+}_{cc} + \pi^0$	20×10^{-3}
$m_{\Omega'^\mu_{cc}}$	4133	$\Omega'^{+\mu}_{cc} \to \Omega^{+\mu}_{cc} + \pi^0$	20×10^{-3}

$$
\begin{aligned}
&- \frac{1}{2} g_\pi \left(\bar{D}^\mu_{\mathbf{Q}} S D_{\mathbf{Q};\mu} - \bar{N}^\mu_{\mathbf{Q}} S N_{\mathbf{Q};\mu} \right) \\
&+ \frac{1}{2} g_\pi \left(\bar{D}^\mu_{\mathbf{Q}} \Phi N_{\mathbf{Q};\mu} + \bar{N}^\mu_{\mathbf{Q}} \Phi D_{\mathbf{Q};\mu} \right) \\
&- \frac{g_A}{f_\pi} \Big[\bar{D}^\mu_{\mathbf{Q}} \gamma_5 \gamma_\nu \partial_\nu \Phi D_{\mathbf{Q};\mu} - \bar{N}^\mu_{\mathbf{Q}} \gamma_5 \gamma_\nu \partial_\nu \Phi N_{\mathbf{Q};\mu} \\
&\qquad + \bar{D}^\mu_{\mathbf{Q}} \gamma_5 \gamma_\nu \partial_\nu S N_{\mathbf{Q};\mu} + \bar{N}^\mu_{\mathbf{Q}} \gamma_5 \gamma_\nu \partial_\nu S D_{\mathbf{Q};\mu} \Big] , \qquad (10)
\end{aligned}
$$

where S and Φ are defined as $M = S + i\Phi = S + 2i\left(\pi^a T^a\right)$ with π^a being the pion fields and $\mathrm{tr}\,(T_a T_b) = (1/2)\delta^{ab}$.

In the case of heavy mesons, the Δ term shifts the masses of the DHBs in the same direction, and the g_π term provides the mass difference between the chiral partners as

$$\Delta M_{B;i} = m_{N_{\mathbf{Q},i}} - m_{D_{\mathbf{Q},i}} = g_\pi \bar{M}_{ii}, \qquad (11)$$

which is exactly the same as that for the heavy–light mesons in Eq. (8). Then, the mass difference for the non-strange doubly heavy baryon is determined as

$$m_{N_{\mathbf{Q},q}} - m_{D_{\mathbf{Q},q}} = 430 \text{ MeV}. \qquad (12)$$

When the quantum numbers of the Ξ^+_{cc} observed in Ref. [6] are identified as $\frac{1}{2}^+$, the mass of the state Ξ^{*+}_{cc} can be estimated to be 3950 MeV.

Next, the intermultiplet one-pion decays of the DHBs in the isospin symmetry limit are considered. The relevant partial widths are expressed as

$$
\begin{aligned}
\Gamma\left(\Xi^{*++}_{cc} \to \Xi^{++}_{cc} + \pi^0\right) &= \Gamma\left(\Xi'^{++\mu}_{cc} \to \Xi^{++\mu}_{cc} + \pi^0\right) \\
&= \frac{(\Delta M_{B;u,d})^2}{8\pi f_\pi^2} |p_\pi| , \qquad (13)
\end{aligned}
$$

where $|p_\pi|$ is the three-momentum of π in the rest frame of the decaying DHB. Other partial widths of different possible charged states can be obtained by using the isospin relation. Our numerical results are given in Table 1.

Next, the DHBs including a strange quark are considered. In such a case, by using the spectrum of the heavy–light meson including a strange quark, one predicts [16–18]

$$m_{N_{\mathbf{Q},s}} - m_{D_{\mathbf{Q},s}} = m_{G_s} - m_{H_s} = 350 \text{ MeV}. \qquad (14)$$

In this sector, due to the conservation of isospin, one might naively expect the dominant transition channel of $\Omega^{*+}_{cc,s}$ to be $\Omega^{*+}_{cc,s} \to \Omega^{+}_{cc,s} + \eta$. However, as the mass splitting (350 MeV) is smaller than the eta meson mass, $m_\eta = 548$ MeV, this channel is forbidden for kinetic reasons, and the dominant channel should be $\Omega^{*+}_{cc,s} \to \Omega^{+}_{cc,s} + \pi^0$ arising from the η–π^0 mixing. The partial decay widths are expressed as

$$
\begin{aligned}
\Gamma\left(\Omega^{*+}_{cc} \to \Omega^{+}_{cc} + \pi^0\right) &= \Gamma\left(\Omega'^{\mu+}_{cc} \to \Omega^{\mu+}_{cc} + \pi^0\right) \\
&= \frac{(\Delta M_{B;s})^2}{2\pi f_\pi^2} \Delta^2_{\pi^0\eta} |p_\pi|, \qquad (15)
\end{aligned}
$$

Table 2
Spectrum of the doubly bottom baryons and the partial widths of one-pion intermultiplet transitions.

Spectrum	Prediction (MeV)	Decay channel	Partial width (MeV)
$m_{\Xi^*_{bb}}$	10,580	$\Xi^{*0}_{bb} \to \Xi^{0}_{bb} + \pi^0$	343
$m_{\Xi^\mu_{bb}}$	10,184	$\Xi^{*0}_{bb} \to \Xi^{-}_{bb} + \pi^+$	686
$m_{\Xi'^\mu_{bb}}$	10,614	$\Xi'^{0}_{bb} \to \Xi^{0\mu}_{bb} + \pi^0$	343
$m_{\Omega^*_{bb}}$	10,658	$\Xi'^{0}_{bb} \to \Xi^{-\mu}_{b} + \pi^+$	686
$m_{\Omega^\mu_{bb}}$	10,342	$\Omega^{*-}_{bb} \to \Omega^{-}_{bb} + \pi^0$	20×10^{-3}
$m_{\Omega'^\mu_{bb}}$	10,692	$\Omega'^{-\mu}_{bb} \to \Omega^{-\mu}_{bb} + \pi^0$	20×10^{-3}

where $\Delta_{\pi^0\eta} = -5.32 \times 10^{-3}$ is the magnitude of the η–π^0 mixing estimated in Ref. [27] based on the two-mixing angle scheme (see, e.g., Ref. [28] and the references therein). As the magnitude of the isospin breaking η–π^0 mixing is very small, the partial width of decay $\Omega^{*+}_{cc,s} \to \Omega^{+}_{cc,s} + \pi^0$ is small. This situation is very similar to that in the heavy–light meson system in which $D_{s0}(2317)$ is regarded as the chiral partner of D_s and the dominant decay channel of the former is the isospin-violating process $D_{s0}(2317) \to D_s + \pi^0$.

Further, a comment is made on the mass splitting of the baryons in a doublet that is beyond the scope of the Lagrangian (9) constructed here. The result obtained in Ref. [29] is as follows:

$$m_{\Psi^{(\prime)\mu}_{cc}} - m_{\Psi^{(*)}_{cc}} = \frac{3}{4}(m_{D^*} - m_D) \simeq 105 \text{ MeV}, \qquad (16)$$

which is smaller than the pion mass. Therefore, in contrast to the heavy–light meson sector, the one-pion intramultiplet decays are forbidden in the DHB sector for kinetic reasons. Note that it is reasonable to expect 60% correction from $\mathcal{O}(1/m_c)$ to result (16) for doubly charmed baryons [24] so that the one-pion decay channel can open for the intermultiplet decay.

For the DHBs constituted by a pair of bottom quarks, the mass differences of the chiral partners of the non-strange and strange baryons are the same as (12) and (14), respectively, and the one-pion transitions between chiral partners have the same expressions as (13) and (15). To obtain the absolute values of the masses and transition widths, $m_{\Xi_{bb}} = 10{,}150$ MeV is considered, which is the average of the central values obtained in Refs. [3,4]. The mass of Ω_{bb} is estimated as $m_{\Omega_{bb}} \simeq m_{\Xi_{bb}} + m_{\Omega_{cc}} - m_{\Xi_{cc}} \simeq 10{,}308$ MeV, which is in agreement with the calculation of Ref. [4]. Our numerical results are summarized in Table 2. For the bottom baryon sector, the mass splitting of the baryons in a doublet can be displayed as follows:

$$m_{\Psi^{(\prime)\mu}_{bb}} - m_{\Psi^{(*)}_{bb}} = \frac{3}{4}(m_{B^*} - m_B) \simeq 34 \text{ MeV}, \qquad (17)$$

which agrees with the theoretical estimations [3,4].

The DHBs with quark content $QQ'q$ are discussed elsewhere, because the spin combination of two different heavy quarks is more complicated than the QQq case.

Note that in this letter only the chiral partners of the ground states are discussed. As in the heavy–light meson sector [30], the discussions proposed here can be easily extended to the excited states.

As the mass splitting between the chiral partners arise from the spontaneous breaking of chiral symmetry, a particularly relevant problem is the effect of extreme condition on this splitting in QCD. From the findings in the heavy–light meson sector [26,31], it is reasonable to expect that the magnitude of the mass splitting will be reduced in hot/dense matter. Such a scenario can be tested in a future scientific facility.

Finally, it must be stressed that the present work mainly concerns the spectrum and dominant strong decay channels of the DHBs. Some other quantities such as the weak transitions of the DHBs through changing a heavy flavor are also interesting phenomenologically. These physics will be reported elsewhere.

In summary, the spectrum and the dominant strong decay properties of the DHBs based on the chiral dynamics were studied. The mass spitting between the lowest lying DHBs and their chiral partners is estimated to be about 450 MeV for the non-strange DHBs and 350 MeV for the strange DHBs. Moreover, it was predicted that, due to kinetic reasons, the dominant decay channel of the parity odd strange DHB is an isospin-violating process, therefore resulting in a small partial width.

Acknowledgments

The work of Y.-L.M. was supported in part by the National Science Foundation of China (NSFC) under Grant No. 11475071 and the Seeds Funding of Jilin University. M.H. was supported in part by the JSPS Grant-in-Aid for Scientific Research (S) No. 22224003 and (c) No. 24540266.

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