Analysis of $\Lambda_c(2595)$, $\Lambda_c(2625)$, $\Lambda_b(5912)$, $\Lambda_b(5920)$ based on a chiral partner structure

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We construct an effective hadronic model including $\Lambda_c(2595)$, $\Lambda_c(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$ regarding them as chiral partners to $\Sigma_c(2455)$, $\Sigma_c(2520)$, $\Sigma_b$ and $\Sigma_b^*$, respectively, with respecting the chiral symmetry and heavy-quark spin-flavor symmetry. We determine the model parameters from the experimental data for relevant masses and decay widths of $\Sigma_c(2455)$ and $\Lambda_c(2595)$. Then, we study the decay widths of $\Lambda_c(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$. We find that, although the decay of $\Lambda_c(2595)$ is dominated by the resonant contribution through $\Sigma_c(2455)$, non-resonant contributions are important for $\Lambda_c(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$, which reflects the chiral partner structure. We also study the radiative decay widths of the baryons, and show that each of their widths is determined from the radiative decay width of their chiral partners.

I. INTRODUCTION

Chiral symmetry and its spontaneous breaking is one of the most important properties to understand the structures of hadrons including light quarks. The spontaneous chiral symmetry breaking is expected to generate a part of hadron masses and causes mass difference between chiral partners. We expect that the study of chiral partner structure will provide a clue for understanding the chiral symmetry and the heavy-quark spin-flavor symmetry. Determining model parameters from the experimental data for relevant masses and decay widths of $\Sigma_c(2455)$ and $\Lambda_c(2595)$, we study the decay widths of $\Lambda_c^*(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$.

This paper is organized as follows: In section II we study the chiral structure of single heavy baryons (SHBs). We construct an effective Lagrangian in section III. Sections IV and V are devoted to study the chiral symmetry and the heavy-quark spin-flavor symmetry. Determining model parameters from the experimental data for relevant masses and decay widths of $\Sigma_c(2455)$, $\Sigma_c^*(2520)$ and $\Lambda_c(2595)$, we study the decay widths of $\Lambda_c^*(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$.

II. CHIRAL STRUCTURE OF SINGLE HEAVY BARYONS

In this section, we study the chiral structure of single heavy baryons using interpolating quark fields.

First we consider interpolating field operators made from up or down quarks:

$$q_{L,R}^i, \quad (i = u, d),$$

where $L$ and $R$ denote left-handed and right-handed chirality, respectively. By using these, we can construct two combinations of diquarks carrying spin-zero which are expressed as

$$(q_{L}^i)^T C q_{L}^j, \quad (q_{R}^i)^T C q_{R}^j,$$

where $^T$ denote the transposition in the spinor space and $C = i\gamma^5\gamma^2$ is the charge conjugation matrix. When the relative angular momentum between two quarks are even, the indices $i$ and $j$ should be anti-symmetrized due to the Fermi statistics. In such a case, we can easily see that both of the above diquarks are chiral singlet. For clarifying this chiral structure, we introduce the following two diquarks which are chiral singlet:

$$\epsilon_{ij} (q_{L}^i)^T C q_{L}^j, \quad \epsilon_{ij} (q_{R}^i)^T C q_{R}^j,$$

where $\epsilon_{ij}$ is anti-symmetric tensor, $\epsilon_{ij} = -\epsilon_{ji}$, with $\epsilon_{ud} = 1$, and the summations over repeated indices are
understood. Since both of the above two diquarks are chiral singlet, two combination of them, which are parity eigenstates, are separately chiral singlet.

Now, let us introduce a field for the chiral-singlet light-quark cloud with $J^P = 0^+$ as

$$\Phi_+ = \epsilon_{ij} (q^i_L)^T C q^j_L + \epsilon_{ij} (q^i_R)^T C q^j_R ,$$

which belongs to $(1, 1)$ representation under $(SU(2)_L, SU(2)_R)$ symmetry. We construct a single heavy baryon by combining this light-quark cloud $(J^P = 0^+)$ to a heavy quark $Q$ ($Q = c, b$). The resultant baryon is a heavy-quark spin singlet, so that we identify it with the lightest $\Lambda_Q$ ($Q = c, b$):

$$\Lambda_Q \sim Q \Phi_+, \quad \Lambda_Q = (\Lambda^+_q, \Lambda^0_q) ,$$

which belongs to $(1, 1)$ representation under $(SU(2)_L, SU(2)_R)$ symmetry.

Next, we consider the following diquark:

$$[\Phi^\mu]_{ij} = [q^i_L C\gamma^\mu q^j_R]_{ij} = (q^i_L)^T C\gamma^\mu q^j_R ,$$

which belongs to $(2, 2)$ representation under $(SU(2)_L, SU(2)_R)$ symmetry. We can easily see that the following property is satisfied:

$$[q^i_R C\gamma^\mu q^j_L]_{ij} = - [q^i_L C\gamma^\mu q^j_R]_{ji} .$$

From these diquarks, we make two combinations of parity eigenstates:

$$[q^i_L C\gamma^\mu q^j_R]_{ij} + [q^j_R C\gamma^\mu q^i_L]_{ij} = [q^i_L C\gamma^\mu q^j_R]_{ij} = [\Phi^\mu]_{ij} ,$$

$$[q^i_L C\gamma^\mu q^j_R]_{ij} - [q^j_R C\gamma^\mu q^i_L]_{ij} = [q^i_L C\gamma^\mu q^j_R]_{ij} = [\Phi^\mu]_{ij} .$$

From the property in Eq. (7), one can easily check that the indices of the diquark with $J^P = 1^+$ is symmetric in the light-quark flavor space, and those of the one with $J^P = 1^-$ is anti-symmetric, i.e.

$$[\Phi^\mu]_{ij} = \Phi^\mu_{ji} ,$$

$$[\Phi^\mu]_{ij} = - \Phi^\mu_{ji} .$$

From this we can easily see that, when the chiral symmetry is spontaneously broken into the isospin symmetry, $\Phi^{\mu}_{(3)}$ is the isospin diquark with $J^P = 1^+$, and $\Phi^{\mu}_{(1)}$ is the isosinglet diquark with $J^P = 1^-$. The diquark $\Phi^\mu$ combined with a heavy quark makes a set of heavy-quark doublets of single heavy baryons (SHBs) with $1/2^-$ and $3/2^-$ as

$$S^\mu_Q \sim Q \Phi^\mu ,$$

where $S^\mu_Q$ denotes the field for the set of SHBs. The $S^\mu_Q$ includes iso-triplet SHBs and iso-singlet SHBs as

$$\Sigma^\mu_Q(1/2^+) , \Sigma^\mu_Q(3/2^+) \sim Q \Phi^\mu_{(3)} ,$$

$$\Lambda_{Q1}(1/2^-) , \Lambda^*_{Q1}(3/2^-) \sim Q \Phi^\mu_{(1)} ,$$

where we omitted the index $\mu$ in the left hand sides. It should be stressed that, since both $\Phi^\mu_{(3)}$ and $\Phi^\mu_{(1)}$ are included in one chiral multiplet $\Phi^\mu$, the heavy quark multiplet of $(\Lambda_{Q1}(1/2^-), \Lambda^*_{Q1}(3/2^-))$ is the chiral partner to that of $(\Sigma^0_Q(1/2^+), \Sigma^0_Q(3/2^+))$. In the present work, we identify $(\Sigma^0_Q(1/2^+), \Sigma^0_Q(3/2^+))$ with the lightest iso-triplet single-heavy baryons with positive parity, and $(\Lambda_{Q1}(1/2^-), \Lambda^*_{Q1}(3/2^-))$ with the lightest iso-singlet ones with negative parity:

$$(\Sigma^+_c, \Sigma^+_b) = (\Sigma^0_c(2455; 1/2^+), \Sigma^0_c(2520; 3/2^+)) ,$$

$$(\Lambda_{c1}, \Lambda^*_{c1}) = (\Lambda_c(2595; 1/2^-), \Lambda_c(2625; 3/2^-)) ,$$

$$(\Sigma^+_b, \Sigma^+_b) = (\Sigma^0_b(1/2^+), \Sigma^0_b(3/2^+)) ,$$

$$(\Lambda_{b1}, \Lambda^*_{b1}) = (\Lambda_b(5912; 1/2^-), \Lambda_b(5920; 3/2^-)) .$$

### III. EFFECTIVE LAGRANGIAN

In this section we construct an effective Lagrangian for the relevant single heavy baryons (SHBs) based on the heavy-quark spin-flavor symmetry and the chiral symmetry. We use the field $\Lambda_Q$ for expressing the SHBs belonging to the chiral singlet in Eq. (4). For expressing the SHBs belonging to chiral $(2, 2)$ representations we introduce the field $S^\mu_Q$ in Eq. (10) which transforms as

$$S^\mu_Q \rightarrow g_S S^\mu_Q g_T^Q ,$$

As we discussed in the previous section, we assume that the fields include the iso-triplet SHBs with positive parity and the iso-singlet SHBs with negative parity as chiral partners to each others. They are embedded into the field $S^\mu_Q$ as

$$S^\mu_Q = \tilde{\Sigma}^\mu_Q + \tilde{\Lambda}^\mu_{Q1} ,$$

where $\tilde{\Sigma}^\mu_Q$ and $\tilde{\Lambda}^\mu_{Q1}$ include the iso-triplet and iso-singlet fields, respectively as

$$\tilde{\Sigma}^\mu_Q = \left( \begin{array}{c} \Sigma^I=1\mu \\ \Sigma^I=0\mu \\ \Sigma^I=0\mu \\ \Sigma^I=1\mu \end{array} \right) ,$$

$$\tilde{\Lambda}^\mu_{Q1} = \left( \begin{array}{c} 0 \\ \frac{1}{\sqrt{2}} \Lambda^\mu_{Q1} \\ \frac{1}{\sqrt{2}} \Lambda^\mu_{Q1} \\ 0 \end{array} \right) .$$

These $\tilde{\Sigma}^\mu_Q$ and $\tilde{\Lambda}^\mu_{Q1}$ are decomposed into spin-3/2 baryon fields and spin-1/2 fields as

$$\Sigma^\mu_Q = \Sigma^\mu_Q - \frac{1}{\sqrt{3}} (\gamma^\mu + v^\mu) \gamma_5 \Sigma^Q ,$$

$$\Lambda^\mu_{Q1} = \Lambda^\mu_{Q1} - \frac{1}{\sqrt{3}} (\gamma^\mu + v^\mu) \gamma_5 \Lambda_{Q1} ,$$

where $\Sigma^\mu_Q$ and $\Lambda^\mu_{Q1}$ denote the spin-3/2 baryon fields, and $\Sigma^Q$ and $\Lambda_{Q1}$ the spin-1/2 fields, respectively. We
note that the parity transformation of the $S_Q^\mu$ field is given by

$$S_Q^\mu \mapsto -\gamma^0 S_Q^{\mu T},$$

(19)

where $T$ denotes the transposition of the $2 \times 2$ matrix in the light-quark flavor space, and that the Dirac conjugate is defined as

$$\bar{S}_Q = S_Q^{\mu t} \gamma^0.$$  

(20)

We introduce a $2 \times 2$ matrix field $M$ for scalar and pseudoscalar mesons including a light quark and a light anti-quark, which belongs to the $(2, 2)$ representation under the chiral SU(2)$_L \times$SU(2)$_R$ symmetry. The transformation properties of $M$ under the chiral symmetry and the parity are given by

$$M \overset{Ch.}{\mapsto} g_L M g_R^\dagger, \hspace{1cm} (21)$$

$$M \overset{P}{\mapsto} M^\dagger.$$  

(22)

We assume that the effective Lagrangian terms for $M$ are constructed in such a way that the $M$ has a vacuum expectation value (VEV) which breaks the chiral symmetry spontaneously, and the VEV is proportional to the pion decay constant $f_\pi$:

$$\langle M \rangle = f_\pi \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}.$$  

(23)

In the following, for studying the decays of the single heavy baryons with emitting pions, we parameterize the field $M$ as

$$M = f_\pi U,$$  

(24)

where

$$U = e^{\frac{2i}{f_\pi} \pi},$$  

(25)

with $\pi$ being the $2 \times 2$ matrix field including pions as

$$\pi = \frac{1}{2} \begin{pmatrix} \pi^0 & \sqrt{2}\pi^+ \\ \sqrt{2}\pi^- & -\pi^0 \end{pmatrix}.$$  

(26)

Now, let us write down an effective Lagrangian including the baryon fields $\Lambda_Q$ and $S_Q^\mu$ together with the meson field $M$, based on the heavy-quark spin-flavor symmetry and the chiral symmetry. We do not include the terms including more than square of $M$ field or more than two derivatives. A possible Lagrangian is given by

$$\mathcal{L}_Q = - \text{tr}(S_Q^\mu (v \cdot iD - \Delta_Q) S_Q^\mu + \tilde{\Lambda}_Q (v \cdot iD) \Lambda_Q$$

$$+ \frac{g_1}{2f_\pi} \text{tr} \left( S_Q^\mu M^\dagger S_Q^\mu v + S_Q^{\mu T} v M^\dagger S_Q^{\mu T} \right)$$

$$- \frac{g_2}{2f_\pi} \text{tr} \left( S_Q^\mu M^\dagger S_Q^\mu v + S_Q^{\mu T} v M^\dagger S_Q^{\mu T} \right)$$

$$- \frac{i}{4f_\pi} \text{tr} \left( S_Q^{\mu u} M v \cdot \partial M S_Q^\mu + S_Q^{\mu u} v \cdot \partial M S_Q^\mu \right)$$

$$- \frac{i}{4f_\pi} \text{tr} \left( S_Q^{\mu u} M v \cdot \partial M S_Q^\mu + S_Q^{\mu u} v \cdot \partial M S_Q^\mu \right)$$

$$+ \frac{h_1}{2f_\pi} \text{tr} \left( S_Q^{\mu u} M v \cdot \partial M S_Q^\mu + S_Q^{\mu u} v \cdot \partial M S_Q^\mu \right)$$

$$+ \frac{h_2}{2f_\pi} \text{tr} \left( S_Q^{\mu u} M v \cdot \partial M S_Q^\mu + S_Q^{\mu u} v \cdot \partial M S_Q^\mu \right)$$

$$+ \frac{g_3}{2\sqrt{2}f_\pi} \tilde{\Lambda}_Q \text{tr} \left( \partial^\mu M S_Q^\mu v^2 - \partial^\mu M S_Q^\mu v^2 \right) \text{h.c.,}$$

(27)

where $m_{\Lambda_Q}$ ($Q = c, b$) are the masses of $\Lambda_c(2286)$ and $\Lambda_b$ in the ground state, $\Delta_Q$ provides the difference between the chiral invariant masses of $(\Sigma_Q, \Lambda_Q)$ chiral multiplet and the chiral singlet $\Lambda_Q$ with heavy-quark flavor violation included. $g_i$ ($i = 1, 2, 3$), $g_2^u$, $h_1^u$, $h_1^R$ and $h_2$ are dimensionless coupling constants. Note that we included $g_2^u$-term to incorporate the heavy-flavor violation needed for explaining the mass differences of charm and bottom sectors. Although we can add heavy-quark flavor violation terms corresponding to $g_1$-term, such contributions are absorbed into the definition of $\Delta_Q$. We expect that heavy-quark flavor violating corrections to other terms are small.

**IV. Masses and $\Sigma_Q^{(*)} \rightarrow \Lambda_Q \pi$ Decays**

In this section, we determine the coupling constants $g_2$ and $g_3^u$ from masses of relevant heavy baryons, and $g_3$.

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1 Here we adopt the normalization of $f_\pi = 92.4$ MeV.
from $\Sigma^0(2595) \to \Lambda_c \pi$ decays. Then we make predictions of $\Sigma^0(2595) \to \Lambda_c \pi$ decays.

When the chiral symmetry is spontaneously broken, the light meson field $M$ acquires its vacuum expectation value as in Eq. (23). Then the masses of $\Sigma^0_Q$ and $\Lambda^0_Q$ are expressed as

$$m(\Sigma^0_Q) = m_{\Lambda_Q} + \Delta_Q + g_1 f_\pi - \frac{g_2^Q}{2} f_\pi,$$

$$m(\Lambda^0_Q) = m_{\Lambda_Q} + \Delta_Q + g_1 f_\pi + \frac{g_2^Q}{2} f_\pi,$$

where $g_2^Q$ is

$$g_2^Q = g_2 + g_2^D \frac{f_\pi}{m_{\Lambda_Q}}.$$

In the present analysis, we assume that the heavy-quark multiplet of $(\Lambda_{c1}, \Lambda_{c2}) = (\Lambda_{c}(2595); J^P = 1/2^-), \Lambda_{c}(2625; 3/2^-)$) is the chiral partner to the multiplet of $(\Sigma_c, \Sigma_c^*) = (\Sigma_c(2455; 1/2^+), \Sigma_c(2520; 3/2^+))$, and that $(\Lambda_{c1}, \Lambda_{c2}) = (\Lambda_c(5912; 1/2^-), \Lambda_c(5920; 3/2^-))$ to $(\Sigma_b, \Sigma_b^*) = (\Sigma_b(1/2^+), \Sigma_b(3/2^+))$. We list experimental data of their masses and full decay widths in Table I. We determine the values of the coupling constants $g_2$ from the mass differences $\Delta M_Q$ of chiral partners with spin-1/2 and spin-3/2 as

$$\Delta M_Q^{(1/2, \text{exp})} = M_{\Lambda_{c1}} - M_{\Sigma_c},$$

$$\Delta M_Q^{(3/2, \text{exp})} = M_{\Lambda_{c2}} - M_{\Sigma_c^*}, (Q = c, b),$$

where $M_{\Lambda_{cQ}}$ and $M_{\Sigma_{cQ}}$ are given by taking the isospin average of relevant masses. Using the values listed in Table I we obtain

$$\frac{\Delta M_Q^{(1/2, \text{exp})}}{f_\pi} = 1.19,$$

$$\frac{\Delta M_Q^{(3/2, \text{exp})}}{f_\pi} = 1.53$$

for charm sector. By taking the spin average of these values, we determine the center value of $g_2^c$ as

$$g_2^c = \frac{1}{3} \left( \frac{\Delta M_Q^{(1/2, \text{exp})}}{f_\pi} + 2 \frac{\Delta M_Q^{(3/2, \text{exp})}}{f_\pi} \right) = 1.30 \text{.}$$

By taking the violation of heavy-spin symmetry, we evaluate the error as

$$g_2^c = 1.30^{+1.53+1.30}_{-1.19-1.30} = 1.30^{+0.23}_{-0.11} \text{.}$$

Similarly, $g_2^b$ is evaluated as

$$g_2^b = 0.980^{+0.069}_{-0.046} \text{.}$$

Let us determine the value of the coupling constant $g_3$ from $\Sigma_c \to \Lambda_c \pi$ decays. We use the experimental values of the full widths of $\Sigma_c^0(2455; 1/2^+), \Sigma_c^0(2520; 3/2^+), \Sigma_c^0(2520; 3/2^+)$ and $\Sigma_c^0(2520; 3/2^+)$ with assuming that the one-pion decay is dominant decay mode for each particle. We first calculate four values of the effective couplings $g_3(\Sigma_c^0(2455; 1/2^+), \Sigma_c^0(2520; 3/2^+)$ and $\Sigma_c^0(2520; 3/2^+)$) with assuming that the one-pion decay is dominant decay mode for each particle. We then calculate the isospin average for $J^P = 1/2^+$ and $3/2^+$ separately, we obtain

$$g_3^{(1/2)} = \frac{g_3(\Sigma_c^0 \to \Lambda_c^+ \pi^-) + g_3(\Sigma_c^0 \to \Lambda_c^+ \pi^-)}{2} = 0.673 \text{,}$$

$$g_3^{(3/2)} = \frac{g_3(\Sigma_c^0 \to \Lambda_c^+ \pi^-) + g_3(\Sigma_c^0 \to \Lambda_c^+ \pi^-)}{2} = 0.695 \text{.}$$

The spin average of the above values are calculated as

$$g_3 = \frac{1}{3} \left( g_3^{(1/2)} + 2g_3^{(3/2)} \right) = 0.688 \text{.}$$

We include the systematic error of the spin average and the statistical error of the experimental data as

$$g_3 = g_3^{(1/2)} + g_3^{(3/2)} + \text{stat. e.}$$

$$- g_3^{(1/2)} - g_3^{(3/2)} - \text{stat. e.}$$

to obtain

$$g_3 = 0.688^{+0.013}_{-0.025} \text{.}$$

We determine the values of the coupling constants $g_2^Q$ ($Q = c, b$) from the mass differences $\Delta M_Q$ of chiral partners in the following way: First, we separately evaluate the mass differences of chiral partners with spin-1/2 and spin-3/2 as
We summarize the estimated values of \( g_2^Q \) and \( g_3 \) in Table II. Using the estimated value of \( g_3 \), we calculate the decay widths of \( \Sigma_c^{(*)} \rightarrow \Lambda Q \pi \) as shown in Table III. This shows that the obtained widths of \( \Sigma_c(JP = 1/2^+) \)

| decay modes | our model [MeV] | expt. [MeV] |
|-------------|-----------------|-------------|
| \( \Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+ \) | 1.96^{+0.97}_{-1.14} | 1.89^{+0.99}_{-1.18} |
| \( \Sigma_c^+ \rightarrow \Lambda_c^0 \pi^0 \) | 2.28^{+0.69}_{-0.12} | < 4.6 |
| \( \Sigma_c^0 \rightarrow \Lambda_c^{-} \pi^- \) | 1.94^{+0.07}_{-0.14} | 1.83^{+0.11}_{-0.29} |
| \( \Sigma_c^{++} \rightarrow \Lambda_c^0 \pi^0 \) | 7.41^{+0.42}_{-0.53} | 4.9^{+0.7}_{-0.8} |
| \( \Sigma_c^{++} \rightarrow \Lambda_c^0 \pi^- \) | 14.1^{+1.1}_{-0.5} | 15.3^{+0.6}_{-0.4} |
| \( \Sigma_c^0 \rightarrow \Lambda_c^{-} \pi^0 \) | 7.25^{+0.27}_{-0.31} | 4.9^{+0.7}_{-0.8} |
| \( \Sigma_c^0 \rightarrow \Lambda_c^{-} \pi^- \) | 11.0^{+0.4}_{-0.3} | 11.5^{+0.2}_{-0.1} |
| \( \Sigma_c^{(*)} \rightarrow \Lambda_Q^{(*)} \pi^0 \) | 12.3^{+0.5}_{-0.5} | 7.5^{+2.2}_{-1.8} |
| \( \Sigma_c^{(*)} \rightarrow \Lambda_Q^{(*)} \pi^- \) | 7.09^{+0.27}_{-0.31} | 4.9^{+0.7}_{-0.8} |

and \( \Sigma_c^*(3/2^+) \) are consistent with each other even though we used common coupling constant \( g_3 \). This implies that heavy quark spin violation between them is small. Furthermore, the predicted widths of \( \Sigma_c^0 \) and \( \Sigma_c^{(*)} \) obtained by the common \( g_3 \) coupling for charm and bottom sectors are consistent with experiments. This indicates that the violation of the heavy-quark flavor symmetry is small at this moment, but precise determination of them by future experiments might require the inclusion of heavy-quark flavor violation. We should note that the predicted widths of \( \Sigma_c^{(*)+} \rightarrow \Lambda_c^+ \pi^0 \) are larger than those of their iso-spin partners since the phase space is larger due to the smallness of the mass of \( \pi^0 \).

V. \( \Lambda_Q^{(*)} \rightarrow \Lambda_Q \pi \pi \) DECAYS

In this section, we consider \( \Lambda_Q^{(*)} \rightarrow \Lambda_Q \pi \pi \) decays. In Fig. 1, we plot the relevant diagrams for \( \Lambda_Q^{(*)} \rightarrow \Lambda_Q \pi \pi \) in our model. In the diagrams (a), (b) and (d), \( \Sigma_Q^{(*)} \) appear as intermediate states, while in the diagram (c) and (e), \( \Lambda_Q^{(*)} \) and \( \Lambda_Q \) couple to two pions directly. It should be noticed that, due to the chiral partner structure, the coupling constant in (c) and (e) is equivalent to the \( \Sigma_Q^{(*)} \rightarrow \Lambda_Q \pi \) coupling in (a), (b) and (d). Then, it is not suitable to drop the contributions in (c) and (e). Actually, as we will show below, They are not negligible for \( \Lambda_{c1}^{(*)} \) and \( \Lambda_{b1}^{(*)} \) decays.

![Figure 1](image-url)

From the diagrams in Fig. 1, the amplitude of \( \Lambda_Q \rightarrow \Lambda_Q \pi \pi \) decays is calculated as

\[
\mathcal{M} = - \frac{g_3}{\sqrt{3} f_2^2} \left( p_2^+ - p_3^+ \right) \bar{u}(p_1, t) \left( \gamma_\mu + \frac{P_\mu}{M} \right) \gamma_5 u_1(P, s) \\
- \frac{g_3}{\sqrt{3} f_2} \left\{ (p_2 + (\vec{h}_1 + i \vec{h}_2) \frac{E_2(p_2)}{f_2}) S_f^+(q) p_3^0 \bar{u}(p_1, t) \left( \gamma_\mu + \frac{q_\mu}{m^-} \right) \gamma_5 (m^0 + q) u_1(P, s) \\
- \frac{g_3}{\sqrt{3} f_2} \left\{ (p_2 + (\vec{h}_1 + i \vec{h}_2) \frac{E_2(p_3)}{f_2}) S_f^0(k) p_3^0 \bar{u}(p_1, t) \left( \gamma_\mu + \frac{k_\mu}{m^0} \right) \gamma_5 (m^0 + k) u_1(P, s) \right\}
\]

where \( P \) is the initial momentum of \( \Lambda_Q, p_1 \) the momentum of \( \Lambda_Q, p_2 \) and \( p_3 \) are the momenta of pions, and \( k \) and \( q \)
the momenta of intermediate $\Sigma_Q$s. $S_f$ is the propagator for the intermediate $\Sigma_Q$s given by

$$S_f(k) = \frac{1}{m_{\Sigma_Q}^2 - k^2 + i m_{\Sigma_Q} \Gamma_{\Sigma_Q}},$$

(42)

where $m_{\Sigma_Q}$ and $\Gamma_{\Sigma_Q}$ are the mass and decay width of intermediate $\Sigma_Q$. We used isospin-averaged values of masses and decay widths in the present analysis. Similarly, the amplitude $\Lambda_{Q1} \to \Lambda_Q \pi^0 \pi^0$ decays is

$$\mathcal{M} = - \frac{g_3}{\sqrt{3} f_\pi} \left( g_2^0 \left( \rho_{\Sigma} + m_{\Sigma} \right) \frac{p_{\Sigma}^2}{f_\pi} \right) \gamma_5 u_1 (P, s)$$

$$- \frac{g_3}{\sqrt{3} f_\pi} \left( g_2 + (h_1 + i h_2) \right) E_{\Sigma} (p_{\Sigma}^3) \frac{p_{\Sigma}^2}{f_\pi} S^f_1 (k) \gamma_5 u_1 (p_1, t) \left( \gamma_\mu + \frac{k_\mu}{m^+} \right) \gamma_5 (m^+ + k) u_1 (P, s).$$

(43)

We determine the relation between the values of $h_1$ and $h_2$ from the full width of $\Lambda_{c1}$ ($\Lambda_c(2595)$). Taking into account the errors of $g_2$, $g_3$ and the total width with $\Lambda_{c1}$, we determine the allowed range of $h_1^1$ and $h_2$ as shown in Fig. 2. Using these values we calculate the two-pion decay widths of $\Lambda_c(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$, which are summarized in Table IV.

![FIG. 2. Allowed range of $h_1^1$ and $h_2$ shown by purple area.](image)

The predicted decay width of $\Lambda_c(2625)$ is consistent with predictions of a quark model in Ref. [13, 14]. We note that the predicted decay widths of $\Lambda_b(5912)$ and $\Lambda_b(5920)$ are extremely tiny due to the phase space suppression. As we will show in the next section, the radiative decay widths for $\Lambda_b(5912)$ and $\Lambda_b(5920)$ are comparable with or even larger than the hadronic decay widths.

Using four sets of parameters in Table V, we study contributions of intermediate states to $\Lambda_{Q1} \to \Lambda_Q \pi \pi$ decays.

![TABLE IV. Predicted widths of $\Lambda_{Q1} \to \Lambda_Q \pi \pi$ decays.](image)

![TABLE V. Four typical parameter sets determined from $\Lambda_{c1} \to \Lambda_c \pi \pi$ decay width.](image)
which are shown in Tables VI-IX. Table VI shows the contributions of intermediate states to $\Lambda_c(2595; 1/2^-) \to \Lambda_c \pi \pi$

| decay mode | intermediate states set 1 [keV] | set 2 [keV] | set 3 [keV] | set 4 [keV] |
|------------|---------------------------------|-------------|-------------|-------------|
| $\Lambda_c(2595; 1/2^-) \to \Lambda_c^+ \pi^+ \pi^-$ | NR(c) | 4.10 | 4.10 | 4.10 | 4.10 |
| $\Sigma_c^{++}(b)$ | 344 | 408 | 438 | 302 |
| $\Sigma_c^0(a)$ | 390 | 466 | 497 | 344 |
| $\Sigma_c^{++}(b) & \Sigma_c^0(a)$ | 15.7 | 26.4 | 21.5 | 18.2 |
| NR(c) & $\Sigma_c^{++}(b)$ | 42.7 | -49.1 | 36.3 | -21.1 |
| NR(c) & $\Sigma_c^0(a)$ | 44.3 | -51.5 | 38.1 | -22.6 |

In this section, we consider radiative decays of the heavy baryons. The relevant Lagrangian is given by

$$\mathcal{L}_{\text{rad}} = \frac{r_1}{F} \text{tr} \left( \tilde{S}_Q Q_{\text{light}} S_{Q}^\nu + \tilde{S}_Q^T Q_{\text{light}} S_{Q}^{\nu T} \right) F_{\mu\nu} + \frac{r_2}{F} \text{tr} \left( \tilde{S}_Q Q_{\text{light}} S_{Q}^\nu - \tilde{S}_Q^T Q_{\text{light}} S_{Q}^{\nu T} \right) \tilde{F}_{\mu\nu}$$

and $\Lambda_b(5920)$ are comparable to resonant contributions, partly because the threshold for $\Sigma_c \pi$ decays are not open. As we stressed, the coupling constant $g_3$ in Figs. I(c) and (e) is fixed from $\Sigma_c \to \Lambda_c \pi$ decay based on the chiral partner structure. Then, the experimental check of non-resonant contributions will give a clue to the chiral symmetry structure.

VI. RADIATIVE DECAYS

$$\mathcal{L}_{\text{rad}} = \frac{r_1}{F} \text{tr} \left( \tilde{S}_Q Q_{\text{light}} S_{Q}^\nu + \tilde{S}_Q^T Q_{\text{light}} S_{Q}^{\nu T} \right) F_{\mu\nu} + \frac{r_2}{F} \text{tr} \left( \tilde{S}_Q Q_{\text{light}} S_{Q}^\nu - \tilde{S}_Q^T Q_{\text{light}} S_{Q}^{\nu T} \right) \tilde{F}_{\mu\nu}$$

In this section, we consider radiative decays of the heavy baryons. The relevant Lagrangian is given by

$$\mathcal{L}_{\text{rad}} = \frac{r_1}{F} \text{tr} \left( \tilde{S}_Q Q_{\text{light}} S_{Q}^\nu + \tilde{S}_Q^T Q_{\text{light}} S_{Q}^{\nu T} \right) F_{\mu\nu} + \frac{r_2}{F} \text{tr} \left( \tilde{S}_Q Q_{\text{light}} S_{Q}^\nu - \tilde{S}_Q^T Q_{\text{light}} S_{Q}^{\nu T} \right) \tilde{F}_{\mu\nu}$$

(44)
where \( F_{\mu\nu} \) is the field strength of the photon and \( \tilde{F}_{\mu\nu} \) is its dual tensor: \( \tilde{F}_{\mu\nu} = (1/2)\epsilon_{\mu\nu\rho\sigma}F^{\rho\sigma} \). \( r_i \) \( (i = 1, \ldots ,4) \) are dimensionless constants, and \( F \) is a constant with dimension one. In this analysis, we take \( F = 350 \) MeV following Ref. \( [15] \). We note that the values of the constants \( r_i \) are of order one based on quark models \( [15] \).

Let us first study the electromagnetic intramultiplet transitions governed by the \( r_1 \)-term in Eq. \( (44) \). Let \( B^* \) denotes the decaying baryon with spin-3/2 \( (B^* = \Lambda_{Q1}; \Sigma_Q), \) and \( B \) the daughter baryon with spin-1/2 \( (B = \Lambda_{Q1}; \Sigma_Q) \). Then the radiative decay width is given by

\[
\Gamma_{B^* \rightarrow B\gamma} = C_{B^*,B\gamma}^2 \frac{16\alpha r_i^2 m_B}{9F^2} m_{B^*}E_\gamma^3
\]  

(45)

where \( \alpha \) is the electromagnetic fine structure constant, \( E_\gamma \) is the photon energy and \( C_{B^*,B\gamma} \) is the Clebsh-Gordon constant given by

\[
C_{\Sigma_Q^+; \Sigma_Q^+; \gamma} = C_{\Sigma_Q^+; \Sigma_Q^+; \gamma} = \frac{2}{3} ,
\]

\[
C_{\Sigma_Q^+; \Sigma_Q^+; \gamma} = C_{\Sigma_Q^0; \Sigma_Q^0; \gamma} = \frac{1}{6} ,
\]

\[
C_{\Sigma_Q^0; \Sigma_Q^0; \gamma} = C_{\Sigma_Q^0; \Sigma_Q^0; \gamma} = \frac{1}{3} ,
\]

\[
C_{\Lambda_{Q1}^+; \Lambda_{Q1}^+; \gamma} = C_{\Lambda_{Q1}^+; \Lambda_{Q1}^+; \gamma} = \frac{1}{6} .
\]  

(46)

From this, one naively expects that ratios of radiative decay widths are determined from the squares of these constants as

\[
C_{\Sigma_Q^+; \Sigma_Q^+; \gamma}^2 : C_{\Sigma_Q^+; \Sigma_Q^+; \gamma}^2 : C_{\Sigma_Q^0; \Sigma_Q^0; \gamma}^2 : C_{\Lambda_{Q1}^+; \Lambda_{Q1}^+; \gamma}^2 = 16 : 1 : 4 : 1 .
\]  

(47)

In Table \( \text{X} \) we show our predictions on the decay widths of \( \Lambda_{Q1} \rightarrow \Lambda_{Q1}\gamma \) and \( \Sigma_Q \rightarrow \Sigma_Q\gamma \) comparing with the predictions in Refs. \( [15] \) and \( [16] \). The predicted values for \( \Sigma_Q \rightarrow \Sigma_Q\gamma \) decay widths are consistent with the ratio in Eq. \( (47) \), while the values for \( \Lambda_{Q1} \rightarrow \Lambda_{Q1}\gamma \) decay...
widths are much smaller than the ratio in Eq. (47). This is because the mass differences between \( \Lambda_{Q_1} \) and \( \Lambda_{Q_1} \) are quite small governing huge phase space suppression. We note that our predictions are consistent with the predictions in Refs. [15] and [16].

TABLE XI. Radiative decay widths of \( \Lambda_{Q_1} \rightarrow \Lambda_{Q_1} \gamma \) and \( \Sigma_{Q} \rightarrow \Sigma_{Q} \gamma \) in unit of keV. The values in the row indicated by “Predictions” are our predicted values, where \( r_1 \) is an undetermined parameter of \( O(1) \). For comparison, we list predictions in Refs. [15] and [16].

| decay mode          | Predictions  | Ref. [15] | Ref. [16] |
|---------------------|--------------|-----------|-----------|
| \( \Sigma_{c\gamma} \rightarrow \Sigma_{c\gamma} \gamma \) | 12 \( r_1^2 \) | - | 11.6 |
| \( \Sigma_{c\gamma} \rightarrow \Sigma_{c\gamma} \gamma \) | 0.75 \( r_1^2 \) | - | 0.85 |
| \( \Sigma_{c\gamma} \rightarrow \Sigma_{c\gamma} \gamma \) | 3.1 \( r_1^2 \) | - | 2.92 |
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 0.13 \( r_1^2 \) | 0.107 \( c_R \) | - |
| \( \Sigma_{b\gamma} \rightarrow \Sigma_{b\gamma} \gamma \) | 0.42 \( r_1^2 \) | - | 0.60 |
| \( \Sigma_{b\gamma} \rightarrow \Sigma_{b\gamma} \gamma \) | 0.024 \( r_1^2 \) | - | 0.02 |
| \( \Sigma_{b\gamma} \rightarrow \Sigma_{b\gamma} \gamma \) | 0.089 \( r_1^2 \) | - | 0.06 |
| \( \Lambda_{b\gamma} \rightarrow \Lambda_{b\gamma} \gamma \) | 0.0013 \( r_1^2 \) | - | - |

We next study the \( \Lambda_{Q_1} \rightarrow \Sigma_{Q} \gamma \) decays which concern the \( r_2 \)-term. The decay widths are expressed as

\[
\Gamma_{\Lambda_{Q_1} \rightarrow \Sigma_{Q} \gamma} = \frac{16\alpha r_2^2 m_{\Sigma_{Q}}}{9 F^2} \, m_{\Lambda_{Q_1}} E_3^3, \\
\Gamma_{\Lambda_{Q_1} \rightarrow \Sigma_{Q} \gamma} = \frac{8\alpha r_2^2 m_{\Sigma_{Q}}}{9 F^2} \, m_{\Lambda_{Q_1}} E_3^3, \\
\Gamma_{\Lambda_{Q_1} \rightarrow \Sigma_{Q} \gamma} = \frac{4\alpha r_2^2 m_{\Sigma_{Q}}}{9 F^2} \, m_{\Lambda_{Q_1}} E_3^3, \\
\Gamma_{\Lambda_{Q_1} \rightarrow \Sigma_{Q} \gamma} = \frac{20\alpha r_2^2 m_{\Sigma_{Q}}}{9 F^2} \, m_{\Lambda_{Q_1}} E_3^3. \\
\]

In Table XI, we show our predictions comparing with those in Ref. [15].

TABLE XII. Radiative decay widths of \( \Lambda_{Q_1} \rightarrow \Lambda_{Q_1} \gamma \) in unit of keV. The values in the row indicated by “Predictions” are our predicted values, where \( r_3 \) is an undetermined parameter of \( O(1) \). For comparison, we list predictions in Ref. [15].

| decay mode   | Predictions  | Ref. [15] |
|--------------|--------------|-----------|
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 25 \( r_3^2 \) | 43 \( r_3^2 \) |
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 110 \( r_3^2 \) | 164 |
| \( \Sigma_{b\gamma} \rightarrow \Sigma_{b\gamma} \gamma \) | 74 \( r_3^2 \) | 893 |
| \( \Sigma_{b\gamma} \rightarrow \Sigma_{b\gamma} \gamma \) | 99 \( r_3^2 \) | 288 |

The width of \( \Sigma_{Q} \rightarrow \Sigma_{Q} \gamma \) decay via the \( r_4 \)-term is given by

\[
\Gamma_{\Sigma_{Q} \rightarrow \Sigma_{Q} \gamma} = \frac{8\alpha r_4^2 f^2}{3 F^4} \, m_{\Sigma_{Q}} E_3^3, \\
\]

and the predicted values are shown in Table XIII.

TABLE XIII. Radiative decay widths of \( \Sigma_{Q} \rightarrow \Sigma_{Q} \gamma \) in unit of keV. The values in the row indicated by “Predictions” are our predicted values, where \( r_4 \) is an undetermined parameter of \( O(1) \). For comparison, we list predictions in Ref. [16].

| decay mode   | Predictions  |
|--------------|--------------|
| \( \Sigma_{c\gamma} \rightarrow \Sigma_{c\gamma} \gamma \) | 127 \( c_{RS} \) |
| \( \Sigma_{c\gamma} \rightarrow \Sigma_{c\gamma} \gamma \) | 58 \( c_{RS} \) |
| \( \Sigma_{c\gamma} \rightarrow \Sigma_{c\gamma} \gamma \) | 54 \( c_{RS} \) |
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 98 \( r_2^2 \) |
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 35 \( r_2^2 \) |
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 31 \( r_2^2 \) |
| \( \Lambda_{c\gamma} \rightarrow \Lambda_{c\gamma} \gamma \) | 81 \( r_2^2 \) |

The \( r_3 \)-term generates the \( \Lambda_{Q_1} \rightarrow \Lambda_{Q1} \gamma \) decay, the width of which is expressed as

\[
\Gamma_{\Lambda_{Q_1} \rightarrow \Lambda_{Q1} \gamma} = \frac{8\alpha r_3^2 f^2}{27 F^4} \, m_{\Lambda_{Q1}} E_3^3. \\
\]

VII. A SUMMARY AND DISCUSSIONS

We constructed an effective hadronic model regarding \( \Lambda_{Q_1} = \{ \Lambda_c(2595, J^P = 1/2^-) \}, \Lambda_b(5912, 1/2^-) \) and \( \Lambda_{Q_1} = \{ \Lambda_c^*(2625, 3/2^-), \Lambda_b^*(5920, 3/2^-) \} \) as chiral partners to \( \Lambda_{Q1} \rightarrow \{ \Sigma_{Q}(2455, 1/2^-), \Sigma_b(1/2^-) \} \) and \( \Lambda_{Q1} \rightarrow \{ \Sigma_c^*(2520, 3/2^+), \Sigma_b^*(3/2^+) \} \), respectively, based on the chiral symmetry and heavy-quark spin-flavor symmetry. We determined the model parameters from the experimental data for relevant masses and decay widths of \( \Sigma_{Q}(2455, 1/2^-), \Sigma_b^*(2520, 3/2^+) \) and \( \Lambda_b(5912, 1/2^-) \). Then, we studied the decay widths of \( \Lambda_{Q1}^*(2625), \Lambda_b(5912) \) and \( \Lambda_b(5920) \). We showed that the coupling constant for non-resonant contributions depicted in Figs. (c) and (c) is fixed from the \( \Sigma_c^* \) decays reflecting the chiral partner.
structure. As a result, the decay of $\Lambda_c(2595)$ is dominated by the resonant contribution through $\Sigma_c(2455)$ depicted in Fig. 1(d), since the threshold of $\Lambda_c(2595) \to \Sigma^+_c(2595)\pi^0$ decay is open. We found non-resonant contributions depicted in Figs. 1(c) and (e) are comparable to resonant contributions for $\Lambda_c(2625)$, $\Lambda_b(5912)$ and $\Lambda_b(5920)$, partly because the threshold for $\Sigma_Q\pi$ decays are not open. Our result indicates that studying non-resonant contributions will give a clue to understand the chiral partner structure for single heavy baryons.

We also studied the radiative decays of $\Sigma^*_Q$ and $\Lambda^*_Q$ using effective interaction Lagrangians in Eq. (44). We showed that there is a relation among $\Sigma^*_Q \to \Sigma_Q\gamma$ and $\Lambda^*_Q \to \Lambda_Q\gamma$ decays reflecting the chiral partner structure, which can be checked in future experiments.

In Table XIV, we summarize our predictions of the decay widths of single heavy baryons. We expect that comparison of these predictions with experimental data will give some clues to understand the chiral structure of single heavy baryons. We note that, since the hadronic decays of $\Lambda_b(5912)$ and $\Lambda^*_b(5920)$ are suppressed by the small phase space factors, radiative decays may be dominant modes.

Several comments are in order. The present model does not include the decay $\Lambda^*_Q \to \Lambda_Q\pi\pi$ via $\Sigma^*_Q$, which needs two pions in the D-wave in the heavy quark limit. We expect that such decays are suppressed compared with the decays having two pions in the S-wave. Similarly, we expect that the decay $\Lambda^*_Q \to \Lambda_Q\pi\pi$ via $\Sigma_Q$ is also suppressed.

It is interesting to extend the present model including only two flavors to the one with the strange quark in addition based on the chiral $SU(3)_L \times SU(3)_R$ symmetry. In the case, the flavor 3 representation including $\Lambda_Q$ becomes the chiral partner to the flavor 6 representation including $\Sigma_Q$. We leave the analysis in a future publication.

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TABLE XIV. Predicted decay widths of single heavy baryons (SHBs). The row indicated by “Our model” shows the predictions of the present analysis. The row indicated by “exp.” shows the experimental values for the full width of the relevant SHBs, in which “<” implies no experimental data.

| SHB $J^P$ | decay modes | Our model [MeV] | exp. [MeV] |
|-----------|-------------|-----------------|------------|
| $\Sigma^0_{b^+}$ 1/2$^+$ $\Lambda_c\pi^+$ | $\Lambda_c\pi^+$ | $1.96^{+0.07}_{-0.14}$ | $1.89^{+0.09}_{-0.18}$ |
| $\Sigma^+_c$ 1/2$^+$ $\Lambda_c\pi^0$ | $\Lambda_c\pi^0$ | $2.28^{+0.09}_{-0.17}$ | $< 4.6$ |
| | $\Lambda_c\gamma$ | $0.043 r_4^2$ | |
| $\Sigma^0_{c^+}$ 1/2$^+$ $\Lambda_c\pi^-$ | $\Lambda_c\pi^-$ | $1.94^{+0.07}_{-0.14}$ | $1.89^{+0.11}_{-0.19}$ |
| $\Sigma^{++}_{c^+}$ 3/2$^+$ $\Lambda_c\pi^+$ | $\Lambda_c\pi^+$ | $14.7^{+0.6}_{-1.1}$ | $14.78^{+0.30}_{-0.40}$ |
| | $\Sigma^{++}_{c^+}\gamma$ | $0.012 r_4^2$ | |
| $\Sigma^{++}_{c^+}$ 3/2$^+$ $\Lambda_c\pi^0$ | $\Lambda_c\pi^0$ | $15.3^{+0.6}_{-1.1}$ | $< 17$ |
| | $\Sigma^{++}_{c^+}\gamma$ | $0.75 r_4^2 \times 10^{-3}$ | $< 17$ |
| | $\Lambda_c\gamma$ | $0.11 r_4^2$ | |
| $\Sigma^{0}_{c^0}$ 3/2$^+$ $\Lambda_c\pi^-$ | $\Lambda_c\pi^-$ | $14.7^{+0.6}_{-1.1}$ | $15.3^{+0.4}_{-0.5}$ |
| | $\Sigma^{0}_{c^0}\gamma$ | $3.1 r_4^2 \times 10^{-3}$ | $< 17$ |
| $\Lambda_c 1/2^-$ $\Lambda_c\pi^+\pi^-$ | $\Lambda_c\pi^+\pi^-$ | $0.562-1.09$ | $0.562-1.09$ |
| | $\Lambda_c\pi^0\pi^0$ | $1.23-2.31$ | $0.562-1.09$ |
| | $\Sigma^{++}_{c^+}\gamma$ | $0.25 r_4^2$ | $< 0.97$ |
| | $\Lambda_c\gamma$ | $0.25 r_4^2$ | |
| $\Lambda_{c^1} 3/2^-$ $\Lambda_c\pi^+\pi^-$ | $\Lambda_c\pi^+\pi^-$ | $0.0618-0.507$ | $0.0618-0.507$ |
| | $\Lambda_c\pi^0\pi^0$ | $0.0431-0.226$ | $0.0618-0.507$ |
| | $\Lambda_{c^1}\gamma$ | $0.13 r_4^2 \times 10^{-3}$ | $< 0.97$ |
| | $\Sigma^{+}_{c^1}\gamma$ | $0.12 r_4^2$ | |
| | $\Sigma^{++}_{c^1}\gamma$ | $0.16 r_4^2$ | |
| | $\Lambda_{c^1}\gamma$ | $0.035 r_4^2$ | |
| $\Sigma^{+}_{b^+}$ 1/2$^+$ $\Lambda_b\pi^+$ | $\Lambda_b\pi^+$ | $6.14^{+0.23}_{-0.45}$ | $9.7^{+0.8}_{-2.8} \pm 1.1$ |
| | $\Sigma^{+}_{b^0}\gamma$ | $7.27^{+0.27}_{-0.53}$ | $9.7^{+0.8}_{-2.8} \pm 1.1$ |
| | $\Lambda_b\gamma$ | $0.074 r_4^2$ | |
| $\Sigma^{0}_{b^0}$ 1/2$^+$ $\Lambda_b\pi^0$ | $\Lambda_b\pi^0$ | $7.02^{+0.27}_{-0.51}$ | $4.9^{+0.4}_{-2.3} \pm 1.1$ |
| | $\Sigma^{0}_{b^0}\gamma$ | $11.0^{+0.4}_{-0.8}$ | $4.9^{+0.4}_{-2.3} \pm 1.1$ |
| | $\Lambda_b\gamma$ | $0.42 r_4^2 \times 10^{-3}$ | $11.5^{+0.7}_{-2.2} \pm 1.5$ |
| $\Sigma^{+}_{b^0}$ 3/2$^+$ $\Lambda_b\pi^+$ | $\Lambda_b\pi^+$ | $12.3^{+0.5}_{-0.9}$ | $12.3^{+0.5}_{-0.9}$ |
| | $\Sigma^{+}_{b^0}\gamma$ | $0.024 r_4^2 \times 10^{-3}$ | $< 0.97$ |
| | $\Lambda_b\gamma$ | $0.074 r_4^2$ | |
| $\Sigma^{0}_{b^0}$ 3/2$^+$ $\Lambda_b\pi^0$ | $\Lambda_b\pi^0$ | $11.9^{+0.4}_{-0.9}$ | $11.9^{+0.4}_{-0.9}$ |
| | $\Sigma^{0}_{b^0}\gamma$ | $0.089 r_4^2 \times 10^{-3}$ | $7.5^{+0.2}_{-0.8} \pm 0.9$ |
| | $\Lambda_b\gamma$ | $0.099 r_4^2$ | |
| $\Lambda_{b^0} 1/2^-$ $\Lambda_b\pi^+\pi^-$ | $\Lambda_b\pi^+\pi^-$ | $(0.67-4.4) \times 10^{-3}$ | $0.67-4.4 \times 10^{-3}$ |
| | $\Lambda_b\pi^0\pi^0$ | $(1.4-6.0) \times 10^{-3}$ | $< 0.66$ |
| | $\Sigma^{0}_{b^0}\gamma$ | $0.098 r_4^2$ | |
| | $\Sigma^{0}_{b^0}\gamma$ | $0.025 r_4^2$ | |
| | $\Lambda_b\gamma$ | $0.027 r_4^2$ | |
| $\Lambda_{b^0} 3/2^-$ $\Lambda_b\pi^+\pi^-$ | $\Lambda_b\pi^+\pi^-$ | $(0.75-13) \times 10^{-3}$ | $0.75-13 \times 10^{-3}$ |
| | $\Lambda_b\pi^0\pi^0$ | $(2.2-12) \times 10^{-3}$ | $< 0.63$ |
| | $\Lambda_{b^0}\gamma$ | $0.0013 r_4^2 \times 10^{-3}$ | $< 0.63$ |
| | $\Sigma^{0}_{b^0}\gamma$ | $0.031 r_4^2$ | |
| | $\Sigma^{0}_{b^0}\gamma$ | $0.081 r_4^2$ | |
| | $\Lambda_b\gamma$ | $0.029 r_4^2$ | |