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Radiative Decays of Excited Λ_Q Baryons in the Bound State Picture

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Abstract

It is shown that, in the bound state picture, the $\Lambda_c(2593) \to \Lambda_c\gamma$ and $\Lambda_c(2625) \to \Lambda_c\gamma$ decays are severely suppressed. On the other hand, for their bottom counterparts, which are predicted to have masses 5900 and 5926 MeV respectively, may have significant radiative branching ratio. In particular, the $\Lambda_b(5926) \to \Lambda_b\gamma$ mode may even dominate over the strong decay mode. The isospin allowed $\Lambda_Q^{**} \to \Sigma_Q\gamma$ mode is expected to be small.

In Ref. [1–4], the bound state picture of heavy baryons was developed. This model is motivated by the large N_c limit, in which the ground state light baryons (nucleon N and Delta Δ) just appear as topological solitons of the Goldstone fields. As a result, the interactions of light baryons and heavy mesons can be studied under chiral perturbation theory. It turns out that, under chiral $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$, the only stable bound states are those with $K = I + s_\ell = 0$. These states can be identified with the low-lying heavy baryons Λ_Q and $\Sigma_Q^{(*)}$, with $I = s_\ell = 0$ and 1 respectively. Then one can deduce the decay properties of the bound state from those of the constituents. For example, the $\Sigma_c \to \Lambda_c\pi$ decay was studied in Ref. [1]. The decay rate is determined by the axial current coupling g_3 defined in Ref. [5].

$$\Gamma(\Sigma_c^{(*)} \to \Lambda_c\pi) = \frac{g_3^2}{6\pi}\frac{|\vec{p}_\pi|^3}{f^2}. \tag{1}$$

In Ref. [1], it was shown that

$$g_3 = g_{soliton} + g_{heavy} \tag{2}$$

where

$$g_{soliton} = \sqrt{\tfrac{3}{2}}g_A, \tag{3a}$$

$$g_{heavy} = -\sqrt{\tfrac{1}{6}}g. \tag{3b}$$

Here $g_A = 1.25$ is the nucleon axial current coupling and g is the heavy meson axial current coupling constant defined in Ref. [7–9]. It is also possible to investigate other decay modes like $\Sigma_c \to \Lambda_c\gamma$ [6] in this picture.

In this paper, we are going to study the electric dipole decay of the orbitally excited Λ_Q, denoted by Λ_Q^{**} [1] in this article. The $\Lambda_c(2625)$ and $\Lambda_c(2593)$ doublet, seen by ARGUS [10],

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[1]There is no universal nomenclature for these excited Λ_c baryons. In Ref. [13] they are denoted by $\Lambda_{c1}^{(*)}$, while J. Körner call them Λ_c^{**} in analogy to D^{**} in the meson sector. The experimentists usually just call them by their masses, i.e., $\Lambda_c(2593)$ and $\Lambda_c(2625)$. The choice of nomenclature in this article just reflects the preference of the author.

CLEO [11] and E687 [12], decay dominantly through the channels $\Lambda_c(2625) \to \Lambda_c\pi\pi$ and $\Lambda_c(2593) \to \Sigma_c\pi$. The absence pf the $\Lambda_c\pi$ mode indicates that these are $I = 0$ states, i.e., excited Λ_c's rather than Σ_c's. Theoretically, these states have been studied in Ref. [13] under chiral perturbation theory and heavy quark symmetry. It is noted that the $\Lambda_c(2593) \to \Sigma_c\pi$ channel is severely suppressed by phase space ($M_{\Lambda_c(2593)} - M_{\Sigma_c} = 141$ MeV). The isospin allowed $\Lambda_c(2625) \to \Sigma_c^*\pi$ channel is simply kinematically forbidden ($M_{\Lambda_c(2625)} - M_{\Sigma_c^*} = 105$ MeV) and have to go through a three-body decay. As a result, it is possible that their electromagnetic decay to Λ_c or $\Sigma_c^{(*)}$ may have a significant branching ratio. However, there are undetermined coupling constants in the lagrangian in Ref. [13] and no definite predictions are given on the actual decay widths.

If these excited Λ_Q's are P-wave states as commonly believed, their parities will be opposite to that of the ground states Λ_Q and Σ_Q. As a result, they can decay through electric dipole transitions to these ground states. The decay width is given by

$$\Gamma(\Lambda_Q^{**} \to \Lambda_Q\gamma) = \tfrac{4}{3}|\vec{p}_\gamma|^3|\vec{\mathbf{P}}|^2, \tag{4}$$

where $\vec{\mathbf{P}}$ is the electric dipole moment between the initial and final state. This width has the order of magnitude of

$$\Gamma(\Lambda_Q^{**} \to \Lambda_Q\gamma) \sim |\vec{p}_\gamma|^3\alpha r^2, \tag{5}$$

where $\alpha \sim 1/137$ is the fine structure constant and r is the typical size of the heavy baryons. Take $r \sim \Lambda_{\rm QCD}^{-1} = (300\ \text{MeV})^{-1}$ and $|\vec{p}_\gamma| \sim 300$ MeV, we have the estimate

$$\Gamma(\Lambda_Q^{**} \to \Lambda_Q\gamma) \sim 2\ \text{MeV}. \tag{6}$$

Compare with the estimates of Ref. [13],

$$\Gamma(\Lambda_c(2593) \to \Sigma_c\pi) \sim 10\ \text{MeV}, \tag{7a}$$

$$\Gamma(\Lambda_c(2625) \to \Lambda_c\pi\pi) \sim 0.1\ \text{MeV}, \tag{7b}$$

it is evident that the E1 decays are very probably observable, especially for the $s = \frac{3}{2}$ state. Hence it is of interest to calculate the electric dipole moments in the bound state picture.

In the bound state model, Λ_c^{**}'s are bound states of heavy mesons $D^{(*)}$ to nucleons N in the first excited state of a simple harmonic potential [4].

$$|\Lambda_c^{**}\rangle = \frac{1}{\sqrt{2}}(|D^{(*)+}n^0\rangle + |D^{(*)0}p^+\rangle) \tag{8}$$

Denoting the position of the chiral soliton and the heavy meson as $\vec{r}_1$ and $\vec{r}_2$, and $\vec{r} = \vec{r}_1 - \vec{r}_2$ their relative position, the electric dipole moment can be written as,

$$\vec{\mathbf{P}} = \sum_{k=1,2} \int d^3r\ \phi_{1s}(\vec{r})\ q_k \vec{r}_k\ \phi_{2p}(\vec{r}), \tag{9}$$

where the $\phi(\vec{r})$'s the the simple harmonic wave functions, and the q_k are the electric charges carried by the chiral soliton and the heavy meson respectively. For the $D^{(*)+} - n^0$ system, $q_1 = 0$, $q_2 = e$ and the dipole moment is

$$\begin{aligned} \vec{\mathbf{P}} &= e \int d^3r\ \phi_{1s}(\vec{r})\ \vec{r}_2\ \phi_{2p}(\vec{r}) \\ &= -e \frac{M_N}{M_N + M_D} \int d^3\ \phi_{1s}(\vec{r})\ \vec{r}\ \phi_{2p}(\vec{r}), \end{aligned} \tag{10}$$

and for the $D^{(*)0} - p^+$ system,

$$\begin{aligned} \vec{\mathbf{P}} &= e \int d^3r\ \phi_{1s}(\vec{r})\ \vec{r}_1\ \phi_{2p}(\vec{r}) \\ &= e \frac{M_D}{M_N + M_D} \int d^3\ \phi_{1s}(\vec{r})\ \vec{r}\ \phi_{2p}(\vec{r}). \end{aligned} \tag{11}$$

Hence, with

$$r_c = |\vec{r}_c| = |\langle 1s|\vec{r}|2p\rangle| = |\int d^3r\ \phi_{1s}(\vec{r})\ \vec{r}\ \phi_{2p}(\vec{r})|, \tag{12}$$

we sum up the two contributions above and get,

$$|\vec{\mathbf{P}}| = \frac{er_c}{2} \frac{M_D - M_N}{M_D + M_N}. \tag{13}$$

Note the there is a cancellation between the $D^{(*)+} - n^0$ and the $D^{(*)0} - p^+$ contribution. Physically it is because the electric dipole vector is pointing from the soliton to the heavy meson in the former case, and is reversed in the latter. The cancellation is exact if $M_D = M_N$. In the real world, $M_D/M_N \sim 2$ and this kinematical suppession factor is about $\frac{1}{3}$.

It is necessary to extract the parameters of the simple harmonic binding potential to calculate r_c. The reduced mass of the system is,

$$\mu_c = \left(\frac{1}{M_N} + \frac{1}{M_D}\right)^{-1} = 625 \text{ MeV} \tag{14}$$

The spin-averaged excitation energy is,

$$\omega_c = \tfrac{1}{6}(4(340) + 2(308)) \text{ MeV} = 329 \text{ MeV}. \tag{15}$$

Hence

$$r_c = \langle 1s|r|2p\rangle = (2\mu_c\omega_c)^{-1/2} = (641 \text{ MeV})^{-1}, \tag{16}$$

Putting all the pieces together, we have

$$\Gamma(\Lambda_c(2593) \to \Lambda_c\gamma) = 0.016 \text{ MeV}, \tag{17a}$$

$$\Gamma(\Lambda_2(2625) \to \Lambda_c\gamma) = 0.021 \text{ MeV}, \tag{17b}$$

The small width is mainly due to the kinematical suppression factor. Given the small width, it would be quite difficult to observe this radiative decay for $\Lambda_c(2593)$. On the other hand, the radiative $\Lambda_c(2625)$, may have a significant branching ratio.

The electric dipole moment for the $\Lambda_b^{**} \to \Lambda_b\gamma$ decay do not suffer the same kinematic suppression. The Λ_b^{**} is a linear combination of the $B^{(*)0}n^0$ bound state and the $B^{(*)-}p^+$ bound state.

$$|\Lambda_b^{**}\rangle = \tfrac{1}{\sqrt{2}}(|B^{(*)0}n^0\rangle + |B^{(*)-}p^+\rangle) \tag{18}$$

The electric dipole moment of the $B^{(*)0}n^0$ system vanishes identically as both the constituents are electrically neutral, while the $B^{(*)-}p^+$ system has a nonvanishing electric dipole moment.

To calculate this electric dipole moment, it is again necessary to extract the parameters of the simple harmonic binding potential. The reduced mass μ_b for the Λ_b^{**} system is

$$\mu_b = \left(\frac{1}{M_N} + \frac{1}{M_B}\right)^{-1} = 796 \text{ MeV}. \tag{19}$$

One can predict the frequency of the harmonic potential ω_b, which is just the excitation energies of Λ_b^{**} over the ground state Λ_b, from the counterpart in the charmed system. In Ref. [4], it was shown that

$$\frac{M_{\Lambda_b^{**}}}{M_{\Lambda_c^{**}}} = 1 - \frac{1}{2}\left(\frac{M_N}{M_D} - \frac{M_N}{M_B}\right), \tag{20}$$

which predicts the masses of the excited Λ_b's at $\Lambda_b(5900)$ and $\Lambda_b(5926)$ (258 and 285 MeV above Λ_b respectively). Hence

$$\omega_b = \tfrac{1}{6}(4(285) + 2(258)) \text{ MeV} = 276 \text{ MeV}. \tag{21}$$

As a result,

$$r_b = \langle 1s|r|2p\rangle = (2\mu_b\omega_b)^{-1/2} = (662 \text{ MeV})^{-1}, \tag{22}$$

The electric dipole moment is

$$\begin{aligned}
\vec{\mathbf{P}} &= \tfrac{1}{2}\textstyle\sum_{k=1,2}\int d^3r\ \phi_{1s}(\vec{r})\ q_k\vec{r}_k\ \phi_{2p}(\vec{r}) \\
&= \tfrac{e}{2}\int d^3r\ \phi_{1s}(\vec{r})\ \left(\tfrac{M_B}{M_B+M_N}\vec{r} - \tfrac{-M_N}{M_B+M_N}\vec{r}\right)\ \phi_{2p}(\vec{r}) \\
&= \tfrac{e}{2}\int d^3r\ \phi_{1s}(\vec{r})\ \vec{r}\ \phi_{2p}(\vec{r})\ \ = \tfrac{e}{2}\vec{r}_b,
\end{aligned} \tag{23}$$

which is expected as the electric dipole moment of a neutral system. Finally putting all the pieces together, we end up with

$$\Gamma(\Lambda_b(5900) \to \Lambda_b\gamma) = 0.090 \text{ MeV}, \tag{24a}$$

$$\Gamma(\Lambda_b(5926) \to \Lambda_b\gamma) = 0.119 \text{ MeV}, \tag{24b}$$

which is still smaller than the naive estimate above (essentially because the quarks are actually fractionally charged which effectively weakens the electromagnetic coupling) but yet measurable, and may even dominate for $\Lambda_b(5926)$. It is expected that this width (or equivalently the branching ratio) will be measured within the next several years.

It should also be noted that, in this picture, the isospin allowed $\Lambda_Q^{**} \to \Sigma_Q \gamma$ modes do not happen in leading order. Physically it means that electric dipole transitions do not couple to spins of particles. Since Λ-type and Σ-type heavy baryons have different spin structures (light valence quark spins are in a singlet state in the former and in a triplet state in the latter case), the electric dipole between them vanish identically. Thus this mode must goes through a magnetic quadruapole transition, which is severely suppressed.

The formalism described in this article can also be applied to the electric dipole transitions of other orbitally excited heavy baryons. For example, the orbitally excited $\Sigma_c^{(*)}$'s can decay to $\Sigma_c^{(*)}\gamma$. In the bound state picture, the $\Sigma_c^{(*)}$ baryons are linear combinations of $D^{(*)}$–N and $D^{(*)} - \Delta$ bound states. The calculations of the relevant electric dipole moments are similar in this case, except one has to take the finite Δ–N mass different into account. These orbitally excited $\Sigma_c^{(*)}$'s, however, can decay strongly to the ground state Λ_c and $\Sigma_c^{(*)}$ and these decays are not phase space suppressed. Hence the electromagnetic decays are expected to be unimportant in these cases.

In conclusion, we have shown that the bound state picture predicts that the $\Lambda_c^{**} \to \Lambda_c \gamma$ decay is severely suppressed. On the other hand, the $\Lambda_b^{**} \to \Lambda_b \gamma$ is unsuppressed and may have a significant branching ratio, especially for the $s = \frac{3}{2}$ state. The $\Lambda_Q^{**} \to \Sigma_Q \gamma$ decay mode is expected to be small. It is hopeful that these statements will be tested by experiments in the coming years.

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