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Magnetic moments of the lowest-lying singly heavy baryons

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A B S T R A C T

A light baryon is viewed as N_c valence quarks bound by meson mean fields in the large N_c limit. In much the same way a singly heavy baryon is regarded as $N_c - 1$ valence quarks bound by the same mean fields, which makes it possible to use the properties of light baryons to investigate those of the heavy baryons. A heavy quark being regarded as a static color source in the limit of the infinitely heavy quark mass, the magnetic moments of the heavy baryon are determined entirely by the chiral soliton consisting of a light-quark pair. The magnetic moments of the baryon sextet are obtained by using the parameters fixed in the light-baryon sector. In this mean-field approach, the numerical results of the magnetic moments of the baryon sextet with spin 3/2 are just 3/2 larger than those with spin 1/2. The magnetic moments of the bottom baryons are the same as those of the corresponding charmed baryons.

1. Introduction

Very recently, Ref. [1] showed that when the number of colors (N_c) goes to infinity singly heavy baryons can be described as $N_c - 1$ valence quarks bound by the meson mean fields that also have portrayed light baryons as N_c valence quarks bound by the same mean fields [2,3], being motivated by Diakonov [4]. The masses of the lowest-lying singly heavy baryons were well reproduced in both the charmed and bottom sectors, and the mass of the Ω_b was predicted within this framework. Using the method developed in Ref. [1], we were able to interpret two narrow Ω_c resonances as exotic baryons belonging to the anti-decapentaplet ($\overline{\mathbf{15}}$) [5] among five Ω_cs found by the LHCb Collaboration [6]. This mean-field approach is called the chiral quark-soliton model (χQSM) [7] (for a review, see Refs. [8,9] and references therein). Very recently, the model has also described successfully strong decays of heavy baryons [10] including those of the newly found two narrow Ω_cs.

The magnetic moments of the heavy baryons have been already investigated within various different approaches such as quark models [11–13], the MIT bag model [14], the quark potential model [15,16], the Skyrme models in bound-state approaches [17,18], a relativistic quark model [19], lattice QCD [20–22], heavy-baryon chiral perturbation theories [23,24], and QCD sum rules [25,26], and so on. Since there are no experimental data available yet, it is of great interest to compare the results with those from several different approaches. Since the baryon anti-triplet ($\overline{\mathbf{3}}$) consists of the light-quark pair with the total light-quark spin $J = 0$, the corresponding magnetic moment vanishes in the present mean-field approach with the infinitely heavy-quark mass limit considered. Thus, in the present work, we want to employ the χQSM to compute the magnetic moments of the lowest-lying singly heavy baryons, in particular, the baryon sextet ($\mathbf{6}$) with both spin $J' = 1/2$ and $J' = 3/2$. The magnetic moments of the light baryons were already studied within the χQSM [27,28]. A merit of this approach is that we can deal with light and heavy baryons on the same footing. All the dynamical parameters required for the present analysis were determined in Ref. [29] based on the experimental data on the magnetic moments of the baryon octet, we have no additional free parameter to handle for those of the heavy baryons. We obtain the results for the magnetic moments of the baryon sextet and compare them with those from other models and lattice QCD. The results turn out to be consistent with those from the other works, in particular, with those from Ref. [19]. Compared with the results from the lattice QCD [20–22], the present ones are consistently larger than them except for the Σ_c^{++} magnetic moment.

The structure of the present work is sketched as follows: In Section 2, we briefly review the general formalism of the χQSM in

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order to compute the magnetic moments of the heavy baryons. In Section 3, we show how to carry out the calculation of the magnetic moments within the present framework, using the dynamical parameters fixed in the light baryon sector. In Section 4, we present the numerical results of the magnetic moments of the heavy baryons, examining the effects of flavor $\mathrm{SU}(3)_{\mathrm{f}}$ breaking. We summarize the present work in the last Section.

2. General formalism

In the mean-field approach, a heavy baryon can be expressed by the correlation function of the $N_c - 1$ light-quark operators, while a heavy quark inside it is regarded as a static color source in the limit of the infinitely heavy quark mass ($m_Q \to \infty$). The heavy quark is required only to makes the heavy baryon a color singlet state. The electromagnetic current we now consider consists of both the light and heavy quark currents

$$J_\mu(x) = \bar{\psi}(x)\gamma_\mu \hat{\mathcal{Q}}\psi(x) + e_Q\,\bar{Q}\,\gamma_\mu Q\,, \tag{1}$$

where $\hat{\mathcal{Q}}$ denotes the charge operator of the light quarks in flavor SU(3) space, defined by

$$\hat{\mathcal{Q}} = \begin{pmatrix} \frac{2}{3} & 0 & 0 \\ 0 & -\frac{1}{3} & 0 \\ 0 & 0 & -\frac{1}{3} \end{pmatrix} = \frac{1}{2}\left(\lambda_3 + \frac{1}{\sqrt{3}}\lambda_8\right). \tag{2}$$

Here, λ_3 and λ_8 are the well-known flavor SU(3) Gell-Mann matrices. The e_Q in the second part of the electromagnetic current in Eq. (1) stands for the heavy-quark charge, which is given as $e_c = 2/3$ for the charm quark or as $e_b = -1/3$ for the bottom quark. The magnetic moment of a heavy quark is proportional to the inverse of the corresponding mass, i.e. $\boldsymbol{\mu} \sim (e_Q/m_Q)\boldsymbol{\sigma}$, so that it should be very small in comparison with the light-quark contributions. It plays an essential role only in describing the baryon anti-triplet, which is understandable, because the light-quark pair constitutes a spin-zero state. However, its effect is rather small when it comes to the baryon sextet. In the lattice QCD [20–22], it is known that the contribution of the heavy quark to the magnetic moments of the baryon sextet is approximately one order smaller than the light-quark contributions. Ref. [19] also examined the heavy-quark contribution separately and found that its effect is in general tiny on the magnetic moments of the baryon sextet.

In principle, one could consider the heavy-quark effects on the magnetic moments as done in the quark models. It would give an overall constant contribution to the magnetic moments of the baryon sextet such as $-e_Q/6m_Q$ [19], which is parametrically very small. However, if one wants to consider the heavy-quark contribution within the present formalism consistently, one should go beyond the mean-field approximation. This is yet a difficult task, since we do not know proper nonperturbative interactions between the light and heavy quarks. Thus, we want to restrict ourselves to the light-quark contribution from the mean-field approximation, so we will ignore in the present work that from the heavy quark current in the limit of $m_Q \to \infty$.

Hence, we will deal with the first term of Eq. (1) when we compute the magnetic moments of heavy baryons by considering the following baryon matrix elements:

$$\langle B_Q|\bar{\psi}(x)\gamma_\mu \hat{\mathcal{Q}}\psi(x)|B_Q\rangle. \tag{3}$$

Since we have ignored the heavy-quark contributions, we obtain the same results for both the charmed and bottom baryons. So, we will mainly focus on the magnetic moments of the charmed baryon sextet in the present work.

The general expressions for the magnetic moments of light baryons have been constructed already in previous works [27–30]. We will extend the formalism for those of heavy baryons in this work. Taking into account the rotational $1/N_c$ and linear m_s corrections, we are able to write the collective operator for the magnetic moments as

$$\hat{\mu} = \hat{\mu}^{(0)} + \hat{\mu}^{(1)}, \tag{4}$$

where $\hat{\mu}^{(0)}$ and $\hat{\mu}^{(1)}$ represent the leading and rotational $1/N_c$ contributions, and the linear m_s corrections respectively

$$\begin{aligned} \hat{\mu}^{(0)} =&\; w_1 D^{(8)}_{\mathcal{Q}3} + w_2 d_{pq3} D^{(8)}_{\mathcal{Q}p}\cdot \hat{J}_q + \frac{w_3}{\sqrt{3}} D^{(8)}_{\mathcal{Q}8}\hat{J}_3, \\ \hat{\mu}^{(1)} =&\; \frac{w_4}{\sqrt{3}} d_{pq3} D^{(8)}_{\mathcal{Q}p} D^{(8)}_{8q} + w_5\left(D^{(8)}_{\mathcal{Q}3}D^{(8)}_{88} + D^{(8)}_{\mathcal{Q}8}D^{(8)}_{83}\right) \\ &+ w_6\left(D^{(8)}_{\mathcal{Q}3}D^{(8)}_{88} - D^{(8)}_{\mathcal{Q}8}D^{(8)}_{83}\right). \end{aligned} \tag{5}$$

The indices of symmetric tensor d_{pq3} run over $p = 4, \cdots, 7$. $\hat{J}_3$ and $\hat{J}_p$ denote the third and the pth components of the spin operator acting on the soliton. $D^{(\nu)}_{ab}(R)$ stand for the SU(3) Wigner matrices in the representation ν, which arise from the quantization of the soliton. $D^{(8)}_{\mathcal{Q}3}$ is defined by the combination of the SU(3) Wigner D functions

$$D^{(8)}_{\mathcal{Q}3} = \frac{1}{2}\left(D^{(8)}_{33} + \frac{1}{\sqrt{3}}D^{(8)}_{83}\right), \tag{6}$$

which is obtained from the SU(3) rotation of the electromagnetic octet current. The coefficients w_i in Eq. (5) encode a concrete dynamics of the chiral soliton and are independent of baryons involved. In fact, w_1 includes the leading-order contribution, a part of the rotational $1/N_c$ corrections, and linear m_s corrections, whereas w_2 and w_3 represent the rest of the rotational $1/N_c$ corrections. The m_s dependent term in w_1 is not explicitly involved in the breaking of flavor SU(3) symmetry. Thus, we will treat w_1 as if it had contained the SU(3) symmetric part, when the magnetic moments are computed. On the other hand, w_4, w_5, and w_6 are indeed the SU(3) symmetry breaking terms. Yet another m_s corrections will come from the collective wave functions, which we will discuss soon. In principle, w_i can be computed within a specific chiral solitonic model such as the χQSM [27,30].

We want to emphasize that the structure of Eq. (5) is rather *model-independent* and is deeply rooted in the hedgehog Ansatz or hedgehog symmetry. Since we consider the embedding of the SU(2) soliton into SU(3) [3], which keeps the hedgehog symmetry preserved, we have $\mathrm{SU}(2)_T \times \mathrm{U}(1)_Y$ symmetry. So, the structure of the collective operator is determined by the $\mathrm{SU}(2)_T \times \mathrm{U}(1)_Y$ invariant tensors

$$\begin{aligned} &d_{abc} = \frac{1}{4}\mathrm{tr}(\lambda_a\{\lambda_b,\lambda_c\}), \quad S_{ab3} = \sqrt{\frac{1}{3}}(\delta_{a3}\delta_{b8} + \delta_{b3}\delta_{a8}), \\ &F_{ab3} = \sqrt{\frac{1}{3}}(\delta_{a3}\delta_{b8} - \delta_{b3}\delta_{a8}). \end{aligned} \tag{7}$$

In this respect, we will determine w_i by using the experimental data on the magnetic moments of the baryon octet as done in Refs. [28,29,31], instead of relying on a specific model. We will briefly show how to fix w_i, using the experimental data in the next Section.

To obtain the magnetic moments of the heavy baryons, the operator $\hat{\mu}$ in Eq. (4) needs to be sandwiched between heavy baryon states. Since we consider the linear m_s corrections perturbatively,