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Ω_c excited states with heavy-quark spin symmetry

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We study the $C = 1$, $S = -2$, $I = 0$ sector, where five excited Ω_c states have been recently observed by the LHCb Collaboration. We start from a recently developed unitarized baryon-meson model that takes, as bare baryon-meson interaction, an extended Weinberg-Tomozawa kernel consistent with both chiral and heavy-quark spin symmetries. This $SU(6) \times$ HQSS scheme leads to a successful description of the observed lowest-lying odd parity charmed $\Lambda_c(2595)$ and $\Lambda_c(2625)$ states, and bottomed $\Lambda_b(5912)$ and $\Lambda_b(5920)$ resonances. Within this model, five odd-parity Ω_c states are dynamically generated, but with masses below 3 GeV, not allowing for an identification with the observed LHCb resonances. We revise this model and explore two different scenarios for the renormalization scheme, that is, using a modified common energy scale to perform the subtractions or utilizing a common ultraviolet cutoff to render finite the ultraviolet divergent loop functions in all channels. In both cases, we show that some (at least three) of the dynamically generated states can be identified with the experimental Ω_c, while having odd parity and $J = 1/2$ or $J = 3/2$. Two of these states turn out to be part of the same $SU(6) \times$ HQSS multiplets as the charmed and bottomed Λ baryons.

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1. Introduction

The LHCb Collaboration [1] recently reported the existence of five excited Ω_c states, by means of analyzing the $\Xi_c^+ K^-$ decay in pp collisions. With masses ranging between 3 and 3.1 GeV, four of these states have been also found by the Belle Collaboration [2]. This discovery has triggered a large activity to determine their nature within quark models, QCD sum-rule schemes, quark-soliton models, lattice QCD or molecular models. The final goal is to answer to the long-standing question whether these states can be explained within the quark model picture and/or these are molecules that are dynamically generated via hadron-hadron scattering.

Within molecular models, there have been previous predictions on excited Ω_c states [3, 4, 5, 6]. In view of the new discoveries, in Ref. [7] two Ω_c resonant states at 3050 MeV and 3090 MeV were generated with spin-parity $J^P = 1/2^-$, reproducing two of the experimental states. Also, within an extended local hidden gauge approach that incorporates low-lying $1/2^+$ and $3/2^+$ baryons together with pseudoscalar and vector mesons [8], two $J^P = 1/2^-$ Ω_c states and one $J^P = 3/2^-$ Ω_c^* were observed, the first two in good agreement with the outcome of [7].

With the aim of incorporating explicitly heavy-quark spin symmetry (HQSS), a proper QCD symmetry that appears when the quark masses become larger than the typical confinement scale, a scheme has been developed in Refs. [5, 9, 6, 10, 11, 12, 13] that implements a consistent $\mathrm{SU}(6)_{\mathrm{lsf}} \times$ HQSS extension of the Weinberg-Tomozawa (WT) interaction, where "lsf" refers to light quark-spin-flavor symmetry. In fact, Refs. [5, 6] are the first baryon-meson molecular analyses, fully consistent with HQSS, of the odd-parity $\Lambda_c(2595)$ $[J = 1/2]$ and $\Lambda_c(2625)$ $[J = 3/2]$ resonances. The same model also generates dynamically the $\Lambda_b(5912)$ and $\Lambda_b(5920)$ narrow resonances, found by LHCb [14], which are HQSS partners [10, 13]. It turns out that the $\Lambda_b(5920)$ resonance is the bottom counterpart of the $\Lambda_c(2625)$, whereas the $\Lambda_b(5912)$ is the bottomed version of the second charmed state that is part of the two-pole structure of the $\Lambda_c(2595)$ [10, 13].

Within this scheme consistent with both chiral symmetry and HQSS, five Ω_c states were found in Ref. [6], three $J = 1/2$ and the two $J = 3/2$ bound states. These five odd-parity Ω_c, Ω_c^* states, stemming from the most attractive $\mathrm{SU}(6)_{\mathrm{lsf}} \times$ HQSS representations, have masses below 3 GeV, and cannot be easily identified with the LHCb resonances. Predicted masses, however, depend not only on the baryon-meson interactions, but also on the adopted renormalization scheme (RS).

In this paper we review the RS used in Ref. [6] and its impact in the generation of the five $\Omega_c^{(*)}$ states. We show how the pole positions can be moved up in energy by implementing a different RS, so as to make then the identification of at least three states with the experimental $\Omega_c^{(*)}$ states feasible [15].

2. Unitarized baryon-meson model with $\mathrm{SU}(6) \times$ HQSS Weinberg-Tomozawa kernel

We consider the charm $C = 1$, strangeness $S = -2$ and isospin $I = 0$ sector, where the $\Omega_c^{(*)}$ excited states are located, as we revisited Ref. [6]. Starting from the baryon-meson pair of pseudoscalar and vector mesons as well as the low-lying $1/2^+$ and $3/2^+$ baryons in the $C = 1, S = -2, I = 0$ sector, the s-wave $\mathrm{SU}(6)_{\mathrm{lsf}} \times$ HQSS WT kernel is used as potential V^J (for a given total angular momentum J) for the Bethe-Salpeter equation (BSE). This leads to a T-matrix

$$T^J(s) = \frac{1}{1 - V^J(s)G^J(s)} V^J(s), \tag{2.1}$$

where $G^J(s)$ is a diagonal matrix that contains the loop functions for the different baryon-meson pairs. The loop function is logarithmically ultraviolet (UV) divergent and needs to be renormalized. This can be done by one-subtraction

$$G_i(s) = \overline{G}_i(s) + G_i(s_{i+}), \tag{2.2}$$

with the finite part of the loop function, $\overline{G}_i(s)$ [16]. The divergent contribution, $G_i(s_{i+})$, needs to be renormalized. Two different renormalization schemes are widely used. On the one hand, the one subtraction at certain scale, $\sqrt{s} = \mu$, such that

$$G_i(\sqrt{s} = \mu) = 0, \tag{2.3}$$

$$G_i^\mu(s_{i+}) = -\overline{G}_i(\mu^2), \tag{2.4}$$

$$G_i^\mu(s) = \overline{G}_i(s) - \overline{G}_i(\mu^2). \tag{2.5}$$

On the other hand, we make finite the UV divergent part of the loop function using a sharp-cutoff regulator Λ in momentum space, so we have

$$G_i^\Lambda(s) = \overline{G}_i(s) + G_i^\Lambda(s_{i+}). \tag{2.6}$$

Note that if one uses channel-dependent cutoffs, the one-subtraction RS, $\mu - RS$, is recovered by choosing in each channel Λ_i such that

$$G_i^{\Lambda_i}(s_{i+}) = -\overline{G}_i(\mu^2). \tag{2.7}$$

However, if one uses a common UV cutoff in a given *CSI* sector, both RSs are independent and will lead to different results.

The different dynamically-generated excited $\Omega_c^{(*)}$ are obtained as poles of the scattering amplitudes in each J sector for $(C = 1, S = -2, I = 0)$. We look at both the first and second Riemann sheets (FRS and SRS) of the variable $\sqrt{s}$, with the following prescription: the poles of the FRS that appear on the real axis below threshold are interpreted as bound states, whereas the poles on the SRS below the real axis and above threshold are identified with resonances. The mass and the width can be found from the position of the pole on the complex energy plane:

$$T_{ij}(s) \simeq \frac{g_i g_j}{\sqrt{s} - \sqrt{s_R}}, \tag{2.8}$$

where the quantity $\sqrt{s_R} = M_R - \mathrm{i}\Gamma_\mathrm{R}/2$ provides the mass (M_R) and the width (Γ_R) of the state, and g_i is the complex coupling of the resonance to the channel i.

3. Excited $\Omega_c^{(*)}$ states

The LHCb Collaboration studied the $\Xi_c^+ K^-$ spectrum using pp collisions and five new narrow excited Ω_c^0 states were identified: the $\Omega_c^0(3000)$, $\Omega_c^0(3050)$, $\Omega_c^0(3066)$, $\Omega_c^0(3090)$ and the $\Omega_c^0(3119)$, the last three also observed in the $\Xi_c^{\prime +} K^-$ decay. Moreover, a sixth broad structure with 3188 MeV was found in the $\Xi_c^+ K^-$ spectrum.

Table 1: Ω_c an Ω_c^* excited states as reported in Ref. [6]. We label them from **a** to **e**, according to their energy position (taken from [15]).

Name	M_R (MeV)	Γ_R (MeV)	J
a	2810.9	0	1/2
b	2814.3	0	3/2
c	2884.5	0	1/2
d	2941.6	0	1/2
e	2980.0	0	3/2

Table 2: Ω_c and Ω_c^* excited states obtained using $\alpha = 1.16$ (taken from [15]).

Name	M_R (MeV)	Γ_R (MeV)	J	M_R^{exp}	Γ_R^{exp}
a	2922.2	0	1/2	—	—
b	2928.1	0	3/2	—	—
c	2941.3	0	1/2	—	—
d	2999.9	0.06	1/2	3000.4	4.5
e	3036.3	0	3/2	3050.2	0.8

3.1 One-subtraction renormalization

As previously mentioned, in Ref. [6] five excited Ω_c states with $J = 1/2^-$ and $J = 3/2^-$ were predicted, with masses below 3 GeV (Table 1). The position in mass makes it difficult to identify any of them with the LHCb resonances.

These states were dynamically generated by solving a coupled-channel BSE using a $\mathrm{SU}(6)_{\mathrm{lsf}} \times$ HQSS-extended WT interaction as a kernel, whereas the baryon-meson loops were renormalized with one-substraction at the scale $\mu = \sqrt{\alpha\left(m_{th}^2 + M_{th}^2\right)}$, with $\alpha = 1$, and m_{th} and M_{th} the masses of the meson and baryon of the channel with the lowest threshold in the given *CSI* sector [4]. However, it is possible to permit some freedom and slightly modify the choice of the subtraction point by changing α. Allowing for just moderately changes, we find that for $\alpha = 1.16$ the two last states (**d** and **e** in Table 1) are now located near the experimental $\Omega_c(3000)$ and $\Omega_c(3050)$ (see Table 2). The state with mass 2999.9 MeV is mainly generated by $\Xi_c^{\prime +}\bar{K}$, whereas the state at 3036.3 MeV has a dominant $\Xi_c^*\bar{K}$ component that can be reconciled with the experimental decay $\Xi_c^+ K^-$ if we allow for $\Xi_c^*\bar{K} \to \Xi_c K$ $d-$wave transition [15].

Given our results, we need to explore a different RS to evaluate the impact of the renormalization procedure in the predictions of $\Omega_c^{(*)}$ in a controlled manner. Thus, we use the relation between the subtraction constants and the cutoff scheme, and employ a common UV cutoff for all baryon-meson loops within reasonable limits. In this manner, we avoid any fictitious reduction of any baryon-meson channel with the use of a small cutoff value and we prevent an arbitrary variation of the subtraction constants.

3.2 Common cutoff regularization

In order to identify our five dynamically generated $\Omega_c^{(*)}$ of Table 1 using the new subtraction constants, we first have to determine how their masses and widths change as we adiabatically vary

Table 3: Ω_c and Ω_c^* excited states calculated using the subtraction constants from a cutoff of $\Lambda = 1090$ MeV (taken from [15]).

Name	M_R (MeV)	Γ_R (MeV)	J	M_R^{exp}	Γ_R^{exp}
a	2963.95	0.0	1/2	—	—
c	2994.26	1.85	1/2	3000.4	4.5
b	3048.7	0.0	3/2	3050.2	0.8
d	3116.81	3.72	1/2	3119.1/ 3090.2	1.1/ 8.7
e	3155.37	0.17	3/2	—	—

the values of the subtraction constants. This is done by

$$G_i(s) = \overline{G}_i(s) - (1-x)\overline{G}_i(\mu^2) + xG_i^{\Lambda}(s_{i+}), \tag{3.1}$$

where x is a parameter that changes adiabatically from 0 to 1, and $\mu^2 = (m_{th}^2 + M_{th}^2)$. In this way, we can follow the original $\Omega_c^{(*)}$ in the complex energy plane as we modified our prescription to use a common cutoff for the computation of the subtraction constants.

In Table 3 we show our results for $\Omega_c^{(*)}$ for a cutoff of $\Lambda = 1090$ MeV. We find that three poles (those named **c**, **b** and **d**) can be identified with the three experimental states at 3000 MeV, 3050 MeV, and 3119 or 3090 MeV. This identification is due to the closeness in energy to the experimental states and because of the dominant contribution to their dynamical generation by the experimental $\Xi_c\bar{K}$ and $\Xi_c'\bar{K}$ channels.

Next, we aim at assessing the dependence of our results on the cutoff. Therefore, we examine higher and lower values of the cutoff. In Fig. 1, we show the pole positions for $\Lambda = 1090$ MeV (Table 3) and two additional cutoffs, approximately 100 MeV apart. We find that for cutoffs below 800 MeV, all states generated dynamically become heavier and wider than the experimental ones, whereas for bigger values than 1300–1350 MeV, our states appear well below 3 GeV. Coming back to Fig. 1, we find that some (probably at least three) of the states observed by LHCb [1] can be identified with three of our $\Omega_c^{(*)}$. In order to make the experimental identification possible, a significant coupling to the $\Xi_c\bar{K}$ channel has to be obtained, often via $\Xi_c^*\bar{K}$ and $\Xi_c\bar{K}^*$ allowing for the $d-$wave transitions.

The molecular nature of the five $\Omega_c^{(*)}$ narrow states has been recently analyzed in Refs. [7, 8], as previously mentioned, as well as the observed broad 3188 MeV in Ref. [17]. In Ref. [7] two $J = 1/2$ baryon-meson molecular states could be identified with the experimental $\Omega_c(3050)$ and $\Omega_c(3090)$. These two states have been reproduced in the $J = 1/2$ sector in Ref. [8], whereas a $J = 3/2$ molecular state has been also identified with the experimental $\Omega_c(3119)$, as the authors incorporate baryon $3/2^+$-pseudoscalar meson states into the hidden gauge approach. Compared to these works, our model for $\Lambda = 1090$ identifies $J = 1/2^-$ $\Omega_c(3000)$, $\Omega_c(3119/3090)$ and $J = 3/2^-$ $\Omega_c(3050)$. This due to the fact that we use a different RS as well as different interaction matrices, in particular for the channels involving D, D^* and light vector mesons.

With regards to the broad structure observed by the LHCb Collaboration around 3188 MeV, the authors of Ref. [17] have indicated that it could be the superposition of two $D\Xi$ bound states. In our case, it is difficult to reach any identification with the experimental result, since most likely our model would also have to consider states from less attractive $\mathrm{SU}(6)_{\mathrm{lsf}}\times$HQSS multiplets [6]. Also,

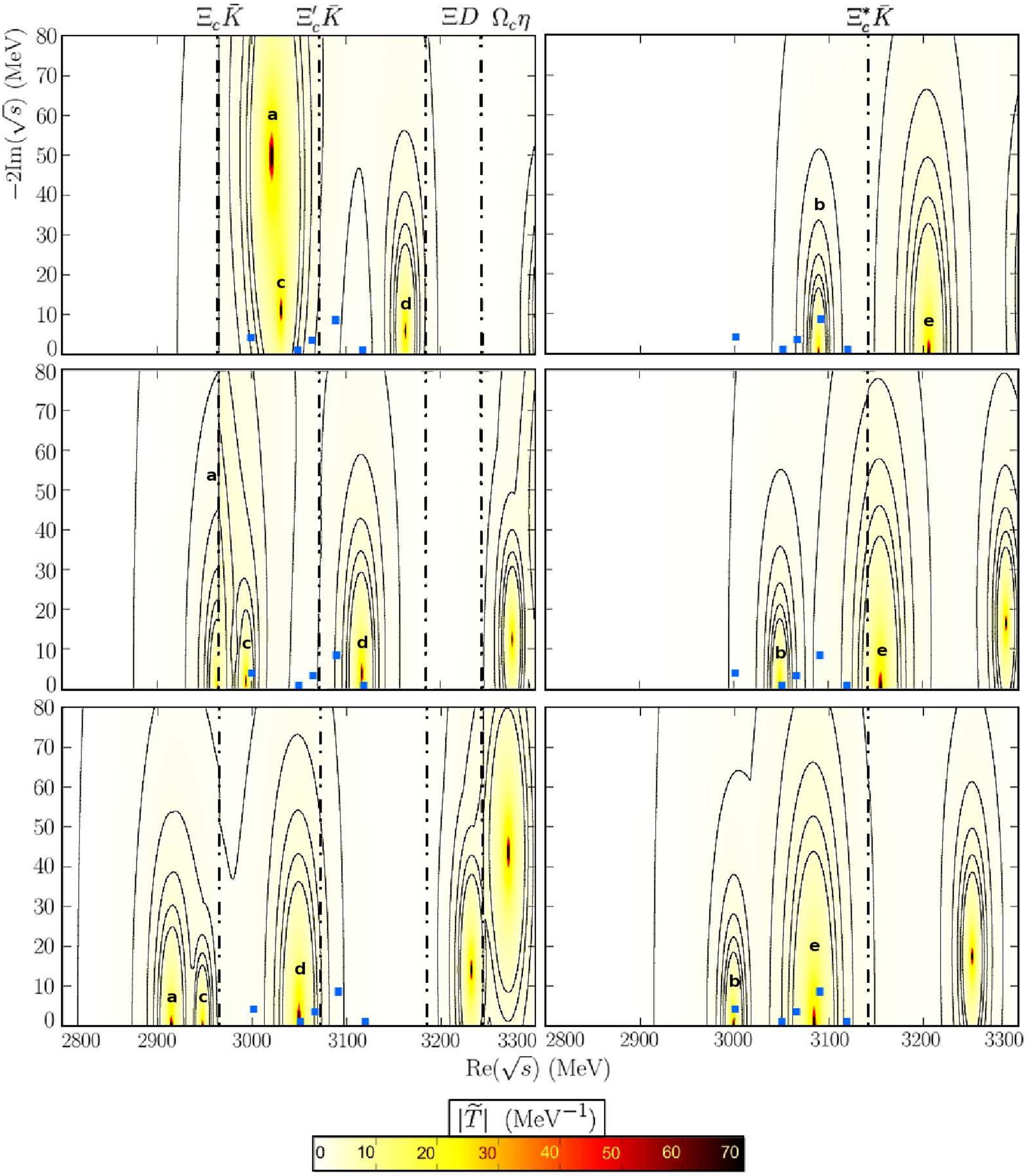

Figure 1: Ω_c and Ω_c^* excited states for different UV cutoffs. The blue squares indicate the experimental results, whereas dashed-dotted lines show the closest baryon-meson thresholds. The left plots are for $J = \frac{1}{2}$ and the right ones for $J = \frac{3}{2}$. For the two largest values of Λ, some resonant states from less attractive $\mathrm{SU(6)}_{\mathrm{lsf}}\times$HQSS multiplets are also visible for higher masses. This figure is taken from [15].

a candidate of a loosely bound molecular state of mass around 3140 MeV is predicted in Ref. [18], whereas in our case we cannot identify it with any of ours. We actually cannot associate our states to those of Ref. [18], as the authors do not consider $\Xi^{(*)}D^{(*)}$ channels.

4. Conclusions

In view of the recent LHCb experimental results [1], we have revisited the $\Omega_c^{(*)}$ states by reviewing the RS used in the unitarized coupled-channel model of Ref. [6] . In this previous work, five odd-parity Ω_c, Ω_c^* states were dynamically generated with masses below 3 GeV, thus making the identification with any of the LHCb resonances difficult.

The predicted masses can be moved to higher energies by implementing a different RS. On the one hand, the common energy-scale used in [6] to perform the subtractions is modified allowing for moderate variations. On the other hand, a common UV cutoff is used to render finite the UV divergent loop functions in all channels.

We conclude that probably at least three of the states observed by LHCb [1] have $J = 1/2^-$ and $J = 3/2^-$. Indeed, those associated to the poles **b** with $J = 3/2$ and **c** with $J = 1/2$ in Table 3 for $\Lambda = 1090$ MeV would belong to the same $\mathrm{SU}(6)_{\mathrm{lsf}} \times$ HQSS multiplets as the $\Lambda_c(2595)$ and $\Lambda_c(2625)$ states, as well as $\Lambda_b(5912)$ and $\Lambda_b(5920)$ resonances.

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