Possibility of the existence of charmed exotica

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We employ the chiral quark-soliton model to describe excited baryons with one heavy quark. Identifying known charmed baryons with multiplets allowed by the model, we argue that apart from regular excitations of the ground-state multiplets, two out of five narrow Ω_c^0 states, recently reported by the LHCb Collaboration, may correspond to the exotic pentaquarks. This interpretation can be easily verified experimentally, since exotic Ω_c^0 states—contrary to the regular excitations—form isospin triplets rather than singlets.

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I. INTRODUCTION

In a very recent paper, the LHCb Collaboration announced five or even six Ω_c^0 states with masses in the range of 3–3.2 GeV [1]. Naturally, they correspond to the excitations of the ground-state multiplets of charmed baryons that in this case form two SU(3) sextets: $1/2^+$ and $3/2^+$. In a recent paper [2], we have shown that these two sextets together with the ground state $\bar{\mathbf{3}}$ that comprises $\Lambda_c(2280)$ and $\Xi_c(2470)$ can be successfully described in terms of the chiral quark-soliton model (χQSM) supplemented by an interaction with a heavy quark in such a way that heavy-quark symmetry [3] is respected. A great advantage of the χQSM consists in a rather restrictive mass formula linking the spectra of light baryons with the heavy ones in question.

In the present paper, we consider excitations of these ground-state multiplets that are predicted within the χQSM. They fall into two distinct categories: the regular excitations that correspond to one-particle excitation of the initial quark configuration, and the exotic ones, which in the present work are identified with collective rotations of the soliton [4–6]. Since different assignments of the Ω_c^0 states are possible, we propose criteria that have to be fulfilled by these excitations. In conclusion, we argue that the most probable assignment is that $\Omega_c^0(3050)$ and $\Omega_c^0(3119)$, which are very narrow with the decay widths around 1 MeV, correspond to the isospin triplet of pentaquarks in the SU(3) $\overline{\mathbf{15}}$, while the remaining states, including rather wide bumps above 3.2 GeV, correspond to the quark excitations of the ground-state sextets and are, therefore, isospin singlets.

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The LHCb discovery triggered several attempts to get insight into the nature of the excited Ω_c^0's in different approaches. This includes the QCD sum rules [7–9], the constituent quark models [10], and lattice QCD [11]. In Refs. [12–14], the new states are treated as bound states of a charm quark and a light diquark, and the authors of Ref. [15] viewed the new states as $\Xi_c K$ and $\Xi_c' K$ molecular states and in some approaches [16] as pentaquarks.

The paper is organized as follows. First, we briefly describe the model and provide formulas for masses and discuss the decay widths (where possible). Next, we compare the χQSM predictions with spectra of excited Λ_c and Ξ_c, and then we discuss possible assignments of newly discovered Ω_c^0 states within the pattern of mass splittings predicted by the model. Finally, we conclude and give estimates of masses of other members of $\overline{\mathbf{15}}$ and excited **6**.

II. CHIRAL QUARK-SOLITON MODEL FOR EXCITED HEAVY BARYONS

The χQSM is based on an argument of Witten [17] that in the limit of large number of colors, N_c relativistic valence quarks generate chiral mean fields represented by a distortion of a Dirac sea that, in turn, influence the valence quarks themselves (for a review, see Ref. [18]) forming a self-organized configuration called a *soliton*. The schematic pattern of light-quark energy levels corresponding to this scenario is depicted in Fig. 1(a). Next, rotations of the soliton, both in flavor and configuration spaces, are quantized semiclassically, and the collective Hamiltonian is computed. The model predicts rotational baryon spectra that satisfy the following selection rules:

(i) allowed SU(3) representations must contain states with hypercharge $Y' = N_c/3$,

(ii) the isospin $\boldsymbol{T}'$ of the states with $Y' = N_c/3$ couples with the soliton spin $\boldsymbol{J}$ to a singlet $\boldsymbol{T}' + \boldsymbol{J} = 0$.

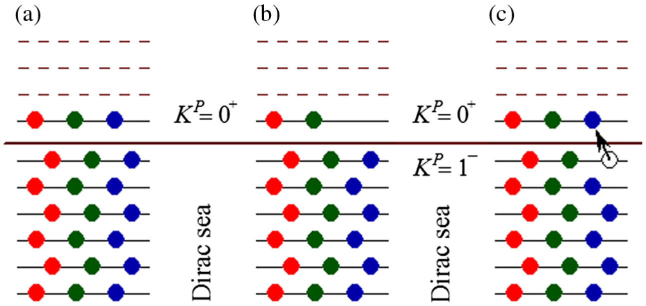

FIG. 1. Schematic pattern of light- (u and d) quark levels in a self-consistent soliton configuration. In (a), all sea levels are filled, and N_c ($=3$ in the figure) valence quarks occupy the $K^P = 0^+$ lowest positive energy level. Unoccupied positive energy levels are depicted by dashed lines. In (b), one valence quark has been stripped off, and the soliton has to be supplemented by a heavy quark not shown in the figure. In (c), a possible excitation of a sea level quark conjectured to be $K^P = 1^-$ to the valence level is shown, and again the soliton has to couple to a heavy quark. Levels for strange quarks that exhibit different filling pattern are not shown.

In the case of light parity (+) baryons, the lowest allowed representations are 8 of spin 1/2, 10 of spin 3/2, and also exotic $\overline{\mathbf{10}}$ of spin 1/2 with the lightest state corresponding to the putative $\Theta^+(1540)$.

In the recent paper [2] following [4], we have extended this model to baryons involving one heavy quark. In this case, the valence level is occupied by $N_c - 1$ light quarks [see Fig. 1(b)] that couple with a heavy quark Q to form a color singlet. The first selection rule in this case reads $Y' = (N_c - 1)/3$. Therefore, the lowest allowed SU(3) representations correspond to the soliton of spin 0 in $\bar{\mathbf{3}}$ and spin 1 in 6. Soliton spin couples with heavy-quark spin to form a spin-1/2 SU(3) triplet and two sextets of spin 1/2 and 3/2 that are subject to a hyperfine splitting. This pattern is confirmed by the data not only qualitatively but also quantitatively as shown in Ref. [2].

The next allowed representation of the rotational excitations corresponds to the exotic $\overline{\mathbf{15}}$ of spin 0 or spin 1. As we will show below, the spin-1 soliton has a lower mass, and when it couples with a heavy quark, it forms spin-1/2 or -3/2 exotic multiplets that should be hyperfine split similarly to the ground-state sextets.

The rotational states described above do not change the parity of the ground-state soliton, and, therefore, they correspond to positive parity. In the present approach, negative parity states are generated by soliton configurations with one light valence quark excited from the valence level or from the Dirac sea. In this way, one can successfully describe the light baryon spectrum up to 2 GeV [6]. In this case, the second selection rule above is modified: $\boldsymbol{T}' + \boldsymbol{J} = \boldsymbol{K}$, where $\boldsymbol{K}$ denotes the so-called *grand spin* of the excited valence quark. Let us remind that the energy levels of the Dirac operator in the presence of the chiral field with hedgehog symmetry are classified by an integer K^P where $\boldsymbol{K} = \boldsymbol{l} + \boldsymbol{s} + \boldsymbol{t}$ with $\boldsymbol{l}$ standing for quark angular momentum, $\boldsymbol{s}$ for its spin, and $\boldsymbol{t}$ for isospin [18]. P denotes parity. The soliton configuration with an excited quark develops its own rotational band. The selection rules for excited quark solitons can be, therefore, summarized as follows:

(i) allowed SU(3) representations must contain states with hypercharge $Y' = (N_c - 1)/3$,
(ii) the isospin $\boldsymbol{T}'$ of the states with $Y' = (N_c - 1)/3$ couples with the soliton spin $\boldsymbol{J}$ as $\boldsymbol{T}' + \boldsymbol{J} = \boldsymbol{K}$, where $\boldsymbol{K}$ is the grand spin of the excited level.

The formula for the soliton mass in the chiral limit for the states in the SU(3) representation $\mathcal{R}$ has been derived in Ref. [5] and reads

$$\mathcal{M}^{(K)} = M_{\rm sol}^{(K)} + \frac{1}{2I_2}\left[C_2(\mathcal{R}) - T'(T'+1) - \frac{3}{4}Y'^2\right] + \frac{1}{2I_1}[(1-a_K)T'(T'+1) + a_K J(J+1) - a_K(1-a_K)K(K+1)], \tag{1}$$

where $C_2(\mathcal{R})$ stands for the SU(3) Casimir operator. $M_{\rm sol}^{(K)} \sim N_c$ denotes classical soliton mass, $I_{1,2}$ represent moments of inertia, and a_K is a parameter that takes into account one-quark excitation. Although all these parameters can be, in principle, calculated in a specific model, we shall follow here a so-called *model-independent* approach introduced in the context of the Skyrme model in Ref. [19], where all parameters are extracted from the experimental data.

Note that $a_K = 0$ if all valence quarks occupy the ground-state level and the soliton spin $J = T'$. For solitons constructed from an excited valence quark configuration $a_K \neq 0$, and the soliton spin J takes the following values:

$$J = |T' - K|, \ldots, |T' + K|. \tag{2}$$

In the case when the strange quark mass $m_s > m_{u,d} \simeq 0$, the soliton mass (1) has to be supplemented by the chiral symmetry-breaking Hamiltonian [5]

$$H_{\rm br} = \alpha D_{88}^{(8)} + \beta\hat{Y} + \frac{\gamma}{\sqrt{3}}\sum_{i=1}^{3} D_{8i}^{(8)}\hat{T}'_i + \frac{\delta}{\sqrt{3}}\sum_{i=1}^{3} D_{8i}^{(8)}\hat{K}_i, \tag{3}$$

which has to be evaluated between the collective wave functions [5,6] that depend on the flavor rotation matrix A,

$$\Psi^{(\mathcal{R};YTT_3)}_{(\mathcal{R}^*;-Y'T'T'_3)}(A) = \sqrt{\dim(\mathcal{R})}(-)^{T'_3-Y'/2} \times D^{(\mathcal{R})*}_{(Y,T,T_3)(Y',T',-T'_3)}(A) \tag{4}$$

coupled to the spin rotational wave function that depends on the rotational matrix S and to the excited quark function χ_{K_3},

$$\Phi^{(\mathcal{R})}_{B,J,J_3,(T',K)} = \sqrt{\frac{2J+1}{2K+1}} \sum_{T'_3,J'_3,K'_3} \left(\begin{array}{cc|c} T' & J & K \\ -T'_3 & J'_3 & K'_3 \end{array} \right) \times (-)^{-(T'+T'_3)} \Psi^{(\mathcal{R};B)}_{(\mathcal{R}^*;-Y'T'T'_3)}(A) D^{(J)*}_{J'_3 J_3}(S) \chi_{K'_3}, \tag{5}$$

where index $(\mathcal{R};YTT_3)$ corresponds to the SU(3) quantum numbers of a given baryon in representation $\mathcal{R}$, and spin index $(\mathcal{R}^*;-Y'T'T'_3)$ is confined to a fixed value of Y' and formally transforms as a member of a representation conjugated to $\mathcal{R}$. The functions $D^{(\mathcal{R})}$ and $D^{(J)}$ are the SU(3) and SU(2) Wigner matrices, respectively, and $\chi_{K'_3} = |K', K'_3\rangle$. $\mathcal{O}(m_s)$ parameters α, β, γ, and δ are computable in terms of single quark wave functions of valence and sea quarks. Their explicit form can be found in, e.g., Ref. [5].

In order to construct a heavy baryon in the present model, we have to strip off one light quark from the valence level and quantize the soliton with a new constraint $Y' = (N_c - 1)/3$. The pertinent light-quark configuration is shown in Fig. 1(b). Such a soliton is coupled with a heavy quark to form a color singlet, and the collective Hamiltonian has to be supplemented by a hyperfine interaction, which we parametrize as follows [2]:

$$H_{\mathrm{hf}} = \frac{2}{3}\frac{\kappa}{m_Q} \boldsymbol{J} \cdot \boldsymbol{J}_Q, \tag{6}$$

where κ is flavor independent. The operators $\boldsymbol{J}$ and $\boldsymbol{J}_Q$ represent the spin operators for the soliton and the heavy quark, respectively.

III. PHENOMENOLOGY OF HEAVY BARYONS IN χQSM

A. Light sector phenomenology

In order to estimate the heavy baryon masses in the χQSM in the model-independent approach, one fixes model parameters from the light sector and uses them for predictions in the heavy-quark sector. This procedure, however, suffers from different systematic uncertainties. For example, there exist corrections to $M_{\mathrm{sol}} \sim N_c$ that are of the order $\mathcal{O}(N_c^0)$ related to the Casimir energy [20,21] and meson loops [22–25], which are beyond control in the present approach. Obviously, in a model-independent approach, these corrections are accommodated in M_{sol} and also in $1/I_{1,2}$. It is, however, unknown how they depend on the soliton quantum numbers and how they change in the presence of a heavy quark due to, for example, nontrivial color interactions between the soliton and an extra quark.

The splittings between multiplets are under much better control than the absolute masses. For example, moment of inertia I_1 can be determined from the mass difference of the mean octet ($\mathcal{M}_{\mathbf{8}} \sim 1150$ MeV) and decuplet ($\mathcal{M}_{\mathbf{10}} \sim 1380$ MeV) masses. Indeed, it follows from (1):

$$\frac{1}{I_1} = \frac{2}{3}(\mathcal{M}_{\mathbf{10}} - \mathcal{M}_{\mathbf{8}}) = 153 \text{ MeV}, \tag{7}$$

which agrees well with the much more complete analysis of Ref. [26] giving $1/I_1 = 160$ MeV.

It is, however, much more difficult to estimate the second moment of inertia I_2, as it contributes only to the masses of exotic pentaquarks. Given the fact that the nonexotic members of $\overline{\mathbf{10}}$ can mix with regular baryons [27], $1/I_2$ estimation suffers from large uncertainty. Also, the mass of Θ^+, whose existence is still upheld by the LEPS Collaboration [28,29], DIANA Collaboration [30], and a part of the CLAS experiment [31] (see, however, the critique in Ref. [32]), suffers from an uncertainty of 20 MeV: 1520–1540 MeV. The best way to extract $1/I_2$ is to use the mass of the exotic Ξ_5, since it does not mix with low mass regular hyperons. Using the values from Refs. [27,33], we obtain

$$\frac{1}{I_2} = 400\text{–}450 \text{ MeV} \tag{8}$$

to be compared with even a larger estimate of Ref. [26]: $1/I_2 = 470$ MeV.

Splittings inside SU(3) multiplets are expressed in terms of $\mathcal{O}(m_s)$ parameters: α, β, and γ. A rather detailed phenomenological analysis, which includes wave function corrections, isospin splittings, and decay rates yields a rather well-constrained result [26], which has been used in Ref. [2] and which we shall be using here as well:

$$\alpha = -255 \text{ MeV}, \quad \beta = -140 \text{ MeV}, \quad \gamma = -101 \text{ MeV}. \tag{9}$$

B. Ground-state multiplets

In order to estimate the masses of the states in $\bar{\mathbf{3}}$ and 6, we have used the general formula (1) with one modification. Since the mean fields are generated by $N_c - 1$ valence quarks [see Fig. 1(b)], we have modified $\mathcal{O}(N_c)$ model parameters by the scaling factor $\rho = (N_c - 1)/N_c$, namely,

$I_{1,2} \to \rho I_{1,2}$ and $\alpha \to \bar{\alpha} = \rho\alpha$. This procedure has been applied in [2] both for average mass splittings between the multiplets and for m_s splittings within multiplets of ground-state baryons. While the rescaling works very well for m_s splittings, it is much less accurate for the moments of inertia $I_{1,2}$. Strictly speaking, rescaling by a factor $(N_c - 1)/N_c$ should work well only for quantities dominated by valence levels, which is probably not the case for I_1. Indeed, the rescaling factor that reproduces well $6 - \bar{\mathbf{3}}$ splitting is equal to $\rho = 0.9$ rather than $2/3$.

Let us briefly summarize the results of Ref. [2]:

(1) The lowest-lying heavy baryons can be indeed grouped in two SU(3) multiplets depicted in Fig. 2: an antitriplet of spin $1/2$ and two sextets of spin $1/2$ and $3/2$.

(2) The sextets are subject to the hyperfine splitting (6) that scales like $1/m_Q$, and the value of the splitting parameter for the charm quark is $\kappa/m_c = 70$ MeV.

(3) Within each multiplet, $\mathcal{R}$ isospin submultiplets split proportionally to the hypercharge $\delta_{\mathcal{R}} Y$ with parameters $\delta_{\bar{\mathbf{3}}} = -180$ MeV and $\delta_{\mathbf{6}} = -120$ MeV. These values extracted from the heavy baryon data are the same for b and c baryons; they are, however, lower by 11% than the values obtained from the splittings of the light baryon octet and decuplet with the help of Eq. (9). This can be explained by an 11% reduction of the strange quark mass in the presence of a heavy quark Q, since the ratio $\delta_{\bar{\mathbf{3}}}/\delta_{\mathbf{6}}$ is the same for both determinations.

(4) Splittings between average $\bar{\mathbf{3}}$ and 6 masses are proportional to $1/I_1$ and are equal in charm and bottom sectors. The value of $1/I_1$ extracted from heavy baryon spectra and from the light baryon spectra require a tiny rescaling factor $\rho = 0.9$.

(5) The model predicts a sum rule that links particles from different multiplets and allows us to calculate Ω_Q^* mass, which is very well satisfied for $Q = c$ and gives a prediction for yet unmeasured Ω_b^*.

C. Exotic $\overline{\mathbf{15}}$ as a rotational excitation

Analogous to the pentaquark $\overline{10}$ representation, also in the present case, the soliton admits exotic representations with the lowest one being $\overline{\mathbf{15}}$ (see Fig. 2). In this section, we study the properties of heavy pentaquarks constructed from a $\overline{\mathbf{15}}$ soliton and a heavy quark. The next possible exotic representation is $\overline{\mathbf{15}'} = (p = 0, q = 4)$ with spin $J = 1$, which, however, is heavier than $\overline{\mathbf{15}}$.

As we can see from Fig. 2, the soliton in $\overline{\mathbf{15}}$ can be quantized both as spin $J = 0$ and 1 (remember that the isospin of the states on $Y' = 2/3$ line corresponds to spin[1]).

[1]From now on, we use numerical values of the quantum numbers corresponding to $N_c = 3$, which does not allow for proper N_c counting.

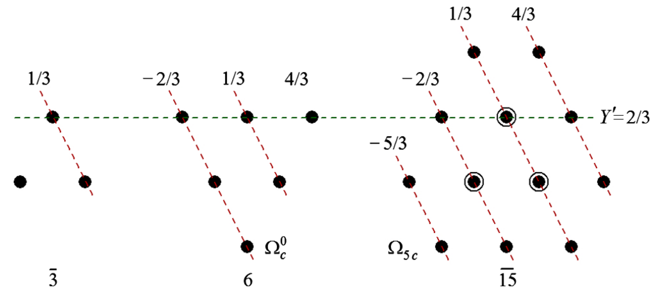

FIG. 2. Rotational band of a soliton with one valence quark stripped off. The soliton spin corresponds to the isospin T' of states on the quantization line $Y' = 2/3$. We show three lowest allowed representations: the antitriplet of spin 0, the sextet of spin 1, and the lowest exotic representation $\overline{\mathbf{15}}$ of spin 1 or 0. Diagonal lines indicate the states of equal charges (shown above the lines). Heavy-quark charge has to be added.

In order to estimate the masses of the states in $\overline{\mathbf{15}}$, we shall use the general formula (1) with the rescaled moments of inertia $I_{1,2} \to \rho I_{1,2}$:

$$\mathcal{M}_{\overline{\mathbf{15}},J=0} = M_{\rm sol} + \frac{5}{2}\frac{1}{\rho I_2},$$
$$\mathcal{M}_{\overline{\mathbf{15}},J=1} = M_{\rm sol} + \frac{3}{2}\frac{1}{\rho I_2} + \frac{1}{\rho I_1}. \qquad (10)$$

Interestingly, the mass difference

$$\Delta_{\overline{\mathbf{15}}} = \mathcal{M}_{\overline{\mathbf{15}},J=0} - \mathcal{M}_{\overline{\mathbf{15}},J=1} = \frac{1}{\rho}\left(\frac{1}{I_2} - \frac{1}{I_1}\right) \qquad (11)$$

is *positive*, since both in the model calculations and model-independent analyses $I_1 \sim 3I_2$, which means—counterintuitively—that the spin-1 soliton is lighter than the spin-0 one.

In order to estimate the masses of exotic heavy baryons, it is useful to relate the mean $\overline{\mathbf{15}}$ mass to the mean 6 mass:

$$\mathcal{M}_{\overline{\mathbf{15}},J=1} = \mathcal{M}_{\mathbf{6}} + \frac{1}{\rho I_2}, \qquad (12)$$

where we have from [2] $\mathcal{M}_{\mathbf{6}} = 2580$ MeV. Given rather large uncertainty of I_2 (8) and of the factor $\rho = 1$–0.66 we get

$$\mathcal{M}_{\overline{\mathbf{15}},J=1} = 2980\text{–}3260 \text{ MeV}. \qquad (13)$$

Finally, we have to calculate matrix elements of the symmetry-breaking Hamiltonian (3). The result reads

$$\Delta_s M_{\overline{\mathbf{15}}} = Y\left(\beta + \frac{17}{144}(\alpha - 2\gamma)\right) + \left(-\frac{2}{27} + \frac{1}{24}\left(T(T+1) - \frac{1}{4}Y^2\right)\right)(\alpha - 2\gamma). \qquad (14)$$

Note that in this case, the δ term does not contribute. Using values from Eq. (9), we obtain $\delta_{\Omega_c} = 180$ MeV, which should be further reduced by 11% giving $\delta_{\Omega_c} = 160$ MeV. We, therefore, predict that Ω_c from $\overline{\mathbf{15}}$ has mass in the range of 3140–3370 MeV before the hyperfine splitting, which we estimate using $\kappa/m_c = 70$ MeV to be -50 and $+20$ MeV for spin 1/2 and 3/2, respectively. Therefore, we see that these rough estimates indicate that some of the states seen by the LHCb might actually be exotic Ω_{5c} pentaquarks. At this point, one should remember that these estimates are subject to the uncertainties due to the $\mathcal{O}(N_c^0)$ corrections discussed above.

The χQSM allows us to estimate the decay widths that proceed through the transition of the light sector associated with the emission of the pseudoscalar meson φ $(=\pi, K, \eta)$. The heavy quark remains in the first approximation intact [3] and acts merely as a spectator. The transition operator can be expressed in terms of three coupling constants and the collective operators:

$$\mathcal{O}_\varphi = 3\left[G_0 D^{(8)}_{\varphi i} - G_1 d_{3bc} D^{(8)}_{\varphi b}\hat{T}'_c - G_2 \frac{1}{\sqrt{3}} D^{(8)}_{\varphi 8}\hat{T}'_i\right] \times \frac{p_i}{M_1 + M_2}. \tag{15}$$

In the present case, we will have transitions $\overline{\mathbf{15}}_1 \to \bar{\mathbf{3}}_0$ (where the lower index refers to $T' = J$) that includes decays of exotic Ω_c measured by the LHCb or $\overline{\mathbf{15}}_1 \to \mathbf{6}_1$ that includes, e.g., decays to $\Omega_c(2535) + \pi$ that have much larger phase space. Sandwiching operator (15) between rotational wave functions (4), one can calculate the effective decay constants

$$\begin{aligned} \overline{\mathbf{15}}_1 \to \bar{\mathbf{3}}_0 \qquad & G_{\bar{\mathbf{3}}} = G_0 - \frac{1}{2}G_1, \\ \overline{\mathbf{15}}_1 \to \mathbf{6}_1 \qquad & G_{\mathbf{6}} = G_0 - \frac{1}{2}G_1 - G_2. \end{aligned} \tag{16}$$

In this normalization, the pion-nucleon decay constant $(g_{\pi NN} \sim 13)$ reads

$$g_{\pi NN} = \frac{7}{10}\left(G_0 + \frac{1}{2}G_1 + \frac{1}{14}G_2\right).$$

Interestingly, in the constituent quark limit [34,35] of the χQSM

$$G_0 = (N_c + 2)G, \qquad G_1 = 4G, \qquad G_2 = 2G. \tag{17}$$

In the present case, however, due to the fact that we have only $N_c - 1$ occupied valence levels, constant G_0 should be replaced

$$G_0 \to \bar{G}_0 = (N_c + 1)G. \tag{18}$$

With this replacement, $G_{\mathbf{6}} = 0$, an effect similar to the nullification of the Θ^+ width in the same limit [35]. We, therefore, expect exotic $\overline{\mathbf{15}}_1$ pentaquarks to have small widths, even if $G_{\bar{\mathbf{3}}} \neq 0$ in the constituent quark limit. We have checked that for a reasonable set of parameters $G_{0,1,2}$, one can indeed get the total decay width being of the order of 1 MeV.

In the present approach, we cannot that easily calculate decay widths that include D mesons. Fortunately, the states that we discuss in this paper are lying below the threshold for such decays.

D. Excited $\bar{\mathbf{3}}$ and 6 multiplets as one-quark excitations

Possible one-quark excitations of the soliton depicted in Fig. 1(b) have been discussed by Diakonov in Ref. [4]. By comparing possible excitation energies with the ones in the light sector, he has come to the conclusion that the most favorable transition that would lead to excited parity $(-)$ heavy baryons corresponds to the transition from a $K^P = 1^-$ sea level to an unoccupied $K^P = 0^+$ state [see Fig. 1(c)]. Such a transition is not allowed in the light baryon sector. The very existence of a $K^P = 1^-$ level as a sea level of the highest energy is, of course, a plausible conjecture that has to be confirmed by model calculations.

The first allowed SU(3) representation for the one-quark excited soliton is again $\bar{\mathbf{3}}$ with $T' = 0$, which—according to (2) for $K = 1$—is quantized as spin 1. From (1), we have

$$\mathcal{M}'_{\bar{\mathbf{3}}} = M'_{\rm sol} + \frac{1}{2I_2} + \frac{1}{I_1}(a_1^2). \tag{19}$$

We will treat $\mathcal{M}'_{\bar{\mathbf{3}}}$ as a phenomenological parameter. The next possibility is flavor 6 with $T' = 1$, which may couple with $K = 1$ to $J = 0$, 1, and 2. From (1), we have

$$\mathcal{M}'_{\mathbf{6}J} = \mathcal{M}'_{\bar{\mathbf{3}}} + \frac{1 - a_1}{I_1} + \frac{a_1}{I_1} \times \begin{cases} -1 & \text{for } J = 0 \\ 0 & \text{for } J = 1 \\ 2 & \text{for } J = 2 \end{cases}. \tag{20}$$

Both the $\bar{\mathbf{3}}$ and the $\mathbf{6}$ are subject to the m_s splittings proportional to the hypercharge Y. For $\bar{\mathbf{3}}$, the splitting parameter is given by the same formula as for the ground-state antitriplet, and, therefore, we know its numerical value [2]:

$$\delta'_{\bar{\mathbf{3}}} = \frac{3}{8}\bar{\alpha} + \beta = \delta_{\bar{\mathbf{3}}} = -180 \text{ MeV}. \tag{21}$$

In the case of the sextet, the splittings depend on the soliton spin and read

$$\delta'_{6J} = \delta_6 - \frac{3}{20}\delta \times \begin{cases} 2 & \text{for } J=0 \\ 1 & \text{for } J=1, \\ -1 & \text{for } J=2 \end{cases} \tag{22}$$

where $\delta_6 = -120$ MeV [2] corresponds to the ground-state sextet splitting. Unfortunately, since we do not know the value of a new parameter δ, we have no handle on the values of different δ'_{6J}.

Furthermore, according to Eq. (6), we have to include the hyperfine splittings with, however, different chromomagnetic constant κ'. The model predicts two SU(3) triplets of spin 1/2 and 3/2, two sextets of spin 1/2 and 3/2, and two sextets of spin 3/2 and 5/2 split by

$$\Delta^{\rm hf}_{\bar{3}} = \Delta^{\rm hf}_{6J=1} = \frac{\kappa'}{m_c}, \qquad \Delta^{\rm hf}_{6J=2} = \frac{5}{3}\frac{\kappa'}{m_c} \tag{23}$$

and one sextet, presumably the lightest one, corresponding to $J = 0$ with no hyperfine splitting.

It is relatively easy to check the χQSM predictions for excited $\bar{3}$, since there are rather well-measured candidates. Indeed, for $(1/2)^-$ we have $\Lambda_c(2592)$ and $\Xi_c(2790)$, and for $(3/2)^-$ there exist $\Lambda_c(2628)$ and $\Xi_c(2818)$. From this assignment, we get $\delta'_{\bar{3}} = -198$ and -190 MeV, respectively, in relative good agreement with Eq. (21). Furthermore, we can extract two other parameters:

$$\frac{\kappa'}{m_c} = \frac{1}{3}(M_{\Lambda_c(2628)} + 2M_{\Xi_c(2818)}) - \frac{1}{3}(M_{\Lambda_c(2252)} + 2M_{\Xi_c(2790)}) = 30 \text{ MeV}, \tag{24}$$

$$\mathcal{M}'_{\bar{3}} = \frac{2}{9}(M_{\Lambda_c(2628)} + 2M_{\Xi_c(2818)}) + \frac{1}{9}(M_{\Lambda_c(2252)} + 2M_{\Xi_c(2790)}) = 2744 \text{ MeV}. \tag{25}$$

In the next section, we shall discuss the phenomenological application of the χQSM to the charmed sextet.

IV. POSSIBLE INTERPRETATIONS OF THE LHCB Ω_c STATES

The natural scenario, which we will follow in this analysis, is that higher spin states (or more precisely, higher J states) become heavier as the spin increases. This assumption leads then to a 6 spectrum depicted schematically in Fig. 3 with the $(J = 0, 1/2^-)$ state corresponding to $\Omega_c(3000)$. This spectrum has to be supplemented by two possible states $1/2^+$ and $3/2^+$ belonging to the exotic $\overline{\mathbf{15}}$.

The splittings $\Delta_{1,2}$ in Fig. 3 correspond to the Ω_c masses for given J before the hyperfine splitting and read

$$\Delta_1 = \frac{a_1}{I_1} + \frac{3}{20}\delta, \qquad \Delta_2 = 2\Delta_1. \tag{26}$$

The χQSM predicts five Ω_c states belonging to the excited sextet. Therefore, we may try to identify all five LCHb resonances with these states. The corresponding scenario is summarized in Table I. We see that in this scenario, the relation between the mass splittings (26) is badly broken. This can be further illustrated by observing that the Ω_c masses satisfy in the χQSM two orthogonal sum rules $\sigma_1 = \sigma_2 = 0$ with

$$\sigma_1 = 6\Omega_c(J=0,1/2^-) - \Omega_c(J=1,1/2^-) - 8\Omega_c(J=1,3/2^-) + 3\Omega_c(J=2,5/2^-),$$
$$\sigma_2 = -4\Omega_c(J=0,1/2^-) + 9\Omega_c(J=1,1/2^-) - 3\Omega_c(J=1,3/2^-) - 5\Omega_c(J=2,3/2^-) + 3\Omega_c(J=2,5/2^-), \tag{27}$$

which are numerically badly violated by the assignment of the minimal scenario. Let us mention that the authors of Ref. [12], who try to interpret the LHCb states within the quark-diquark model, came to the similar conclusion.

As can be seen from Table I, the parameter for the hyperfine splitting deviates considerably from that determined from the experimental data for the excited $\bar{3}'$ given in Eq. (24). Also, in this scenario the widths of the would-be hyperfine split partners are very different. All these arguments suggest that such minimal scenario is not realistic in the mean-field picture of baryons.

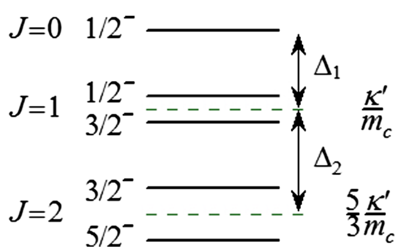

FIG. 3. Schematic spectrum of excited sextets.

TABLE I. Scenario 1. All LHCb Ω_c states are assigned to the excited sextets. This assignment requires hyperfine splitting, which is almost 2 times smaller than in the $\bar{3}$ case and relation (26) is badly broken.

J	S^P	M (MeV)	κ'/m_c (MeV)	Δ_J (MeV)
0	$\frac{1}{2}^-$	3000	not applicable	not applicable
1	$\frac{1}{2}^-$	3050	16	61
	$\frac{3}{2}^-$	3066		
2	$\frac{3}{2}^-$	3090	17	47
	$\frac{5}{2}^-$	3119		

TABLE II. Scenario 2. Only three LHCb states are assigned to the sextets. Using relations (26) and (23), we calculate the masses of $J = 2$ states (marked in italics) that fall into a large bump seen by the LHCb above 3.2 GeV. In this scenario, two narrow states $\Omega_c(3050)$ and $\Omega_c(3119)$ are interpreted as exotic $\overline{\mathbf{15}}$ pentaquarks.

J	S^P	M (MeV)	κ'/m_c (MeV)	Δ_J (MeV)
0	$\frac{1}{2}^-$	3000	not applicable	not applicable
1	$\frac{1}{2}^-$	3066	24	82
	$\frac{3}{2}^-$	3090		
2	$\frac{3}{2}^-$	*3222*	Input	Input
	$\frac{5}{2}^-$	*3262*	24	164

TABLE III. Predicted masses (in MeV) of $1/2^+$ and $3/2^+$ $\overline{\mathbf{15}}$-plet under the assumption that Ω_c members are identified with the observed $\Omega_c(3050)$ and $\Omega_c(3119)$.

	Y	T	$S^P = \frac{1}{2}^+$	$S^P = \frac{3}{2}^+$
B_c	$\frac{5}{3}$	$\frac{1}{2}$	2685	2754
Σ_c	$\frac{2}{3}$	1	2808	2877
Λ_c	$\frac{2}{3}$	0	2806	2875
Ξ_c	$-\frac{1}{3}$	$\frac{1}{2}$	2928	2997
$\Xi_c^{3/2}$	$-\frac{1}{3}$	$\frac{3}{2}$	2931	3000
Ω_c	$-\frac{4}{3}$	1	3050	3119

Given that the minimal scenario does not work, we may try to attribute some of the five narrow LHCb Ω_c's to the possible exotic $\overline{\mathbf{15}}$ multiplet which naturally emerges in our picture. The states $\Omega_c(3050)$ and $\Omega_c(3119)$ are good candidates for the $1/2^+$ and $3/2^+$ hyperfine split Ω_c members of the $\overline{\mathbf{15}}$. First, the corresponding hyperfine splitting parameter $\kappa/m_c \approx 70$ MeV is in excellent agreement with the same parameter determined from the data on the ground-state sextet [2]. Second, the widths of $\Omega_c(3050)$ and $\Omega_c(3119)$ are of order 1 MeV, in agreement with our expectations discussed in Sec. III C above. The assignment of the LHCb states in this scenario is summarized in Table II. We see that in this scenario, the excited sextet states with $J = 2$ have masses above the $\Xi + D$ threshold at 3185 MeV; i.e., they can have rather large widths and are not clearly seen in the LHCb data.

We have tried several other possibilities to distribute the observed states over the negative parity excited sextet and the positive parity $\overline{\mathbf{15}}$; however, all of them give a less consistent picture.

One can check the suggested identification of new Ω_c states in various ways. The simplest one would be to search for the *isospin* partners of Ω_c^0 from the $\overline{\mathbf{15}}$. For example, they can be searched in mass distribution of $\Xi_c^0 + K^-$ or $\Xi_c^+ + \bar{K}^0$; the Ω_c^0's from the sextet do not decay into these channels. Another possibility is to search for the other exotic members of the $\overline{\mathbf{15}}$, especially the lightest B_c baryons (see the next section).

V. MORE ON EXOTIC $\overline{\mathbf{15}}$

A. Partners of $\Omega_c(3050)$ and $\Omega_c(3119)$

In the previous section, we have demonstrated that the favorable scenario is to identify the observed narrow resonances $\Omega_c(3050)$ and $\Omega_c(3119)$ as the $1/2^+$ and $3/2^+$ members of the exotic $\overline{\mathbf{15}}$ multiplet. Now with the help of the mass formula (16), we can predict the masses of other members of the exotic $\overline{\mathbf{15}}$. The parameters of the mass formula (α, β, γ) are fixed by the spectrum of the ground-state light multiplets (9). Note that the spectrum has to be calculated using rescaled $\alpha \to \bar{\alpha} = 2/3\alpha$. Furthermore, the splittings have to be reduced by 11% to account for the effect discussed in Sec. III B. The predicted masses[2] of $\overline{\mathbf{15}}$ are summarized in Table III.[3] Note that with these numbers, we get $\mathcal{M}_{\overline{\mathbf{15}},J=1} = 2935$ MeV, just a little below the lower limit of Eq. (13).

The exotic $\overline{\mathbf{15}}$-plet contains six explicitly exotic states: B_c^+, B_c^{++} (with the minimal quark content $cudd\bar{s}$ and $cuud\bar{s}$), $\Xi_c^{3/2-}, \Xi_c^{3/2++}$ ($cdds\bar{u}$, $cuus\bar{d}$), and Ω_c^-, Ω_c^+ ($cdss\bar{u}$, $cuss\bar{d}$). The detailed properties of the $\overline{\mathbf{15}}$-plet will be studied elsewhere. Here we note that the predicted mass of the lightest $\overline{\mathbf{15}}$ member, the B_c baryon, lies slightly below the strong decay threshold to $(\Lambda_c, \Sigma_c) + K$; hence, we predict that the B_c baryon decays only weakly.

The $\overline{\mathbf{15}}$-plet was discussed for the first time by Diakonov in Ref. [4]. In this paper, the $\overline{\mathbf{15}}$-plet was obtained due to a specific quark transition between quark levels in the mean field (an analog of the Gamow-Teller transition), so the picture there is different from ours. In Ref. [4], the $\overline{\mathbf{15}}$-plet is considerably lighter than in our picture. For example, the mass of the B_c baryon is 2420 MeV. We shall compare in detail the two pictures elsewhere.

B. On excited Ω_b

The mean-field picture of baryons presented here can be easily generalized to baryons with a bottom quark. The main feature of our approach is that the mean field does not depend on the heavy-quark mass. So, if we replace the charm quark by the bottom one, we have to make an overall shift of the masses and rescale the hyperfine splittings.

[2]Note that the predicted masses can be affected by the mixing of nonexotic members of the $\overline{\mathbf{15}}$ with the ground state and excited $\bar{\mathbf{3}}$ and 6, similar to how it happens in the light baryon sector; see Ref. [27].

[3]We adopt the naming scheme suggested by Diakonov [4].

As for the overall shift of the masses, we take the mass difference of the ground-state antitriplets for charmed and bottom baryons:

$$M^b_{\bar{\mathbf{3}}} - M^c_{\bar{\mathbf{3}}} = 3327\ \text{MeV}, \tag{28}$$

which was determined in Ref. [2], where we have also demonstrated that the ratio of the hyperfine mass splittings in the charm and bottom ground-state sextets is close to ~0.3, being in excellent agreement with the mass ratio m_c/m_b.

Performing the overall mass shift and rescaling the hyperfine splittings, we obtain the following prediction for the excited Ω_b: $\Omega_b(6327, 1/2^-)$, $\Omega_b(6404, 1/2^-)$, $\Omega_b(6411, 3/2^-)$, $\Omega_b(6566, 3/2^-)$, and $\Omega_b(6578, 5/2^-)$ belonging to the excited sextets, and $\Omega_b(6409, 1/2^+)$ and $\Omega_b(6430, 3/2^+)$ belonging to the exotic $\overline{\mathbf{15}}$-plet.

VI. SUMMARY AND CONCLUSIONS

The goal of the present paper was to classify the Ω_c baryons that have been recently reported by the LHCb Collaboration [1], employing the mean-field approach. The mean-field picture of baryons being justified by the large-N_c limit offers a unified description of light and heavy baryons. We have shown in Ref. [2] that the universal mean field gives simultaneously a good description of the ground-state $\bar{\mathbf{3}}$ and 6 multiplets of heavy baryons. Also, the ground-state light baryon multiplets are well described [6]. In the present work, we have demonstrated that the same picture predicts the following excited states for heavy-quark baryons in the mass region of 3000–3200 MeV:

(i) two hyperfine split ($1/2^-$ and $3/2^-$) $\bar{\mathbf{3}}'$ which experimentally have very good candidates,

(ii) five excited sextets (rotationally and hyperfine split) with quantum numbers $(J = 0, 1/2^-)$, $(J = 1, 1/2^-, 3/2^-)$, and $(J = 2, 3/2^-, 5/2^-)$, where J denotes the soliton spin,

(iii) two hyperfine split exotic $\overline{\mathbf{15}}$-plets with quantum numbers $1/2^+$ and $3/2^+$.

Because of the universality of our mean-field picture, the basic properties of these excitations are fixed by light baryons and by ground-state multiplets of heavy-quark baryons. The predictions for the excited $\bar{\mathbf{3}}'$-plets are in excellent agreement with the experimental spectrum of the excited Λ_c and Ξ_c.

The observation of the new excited Ω_c^0's allows us to get insight into the excited sextets and $\overline{\mathbf{15}}$-plets. We identify the observed $\Omega_c(3000)$, $\Omega_c(3066)$, and $\Omega_c(3090)$ with $(J = 0, 1/2^-)$ and $(J = 1, 1/2^-, 3/2^-)$ states from the excited sextet, whereas we identify the most narrow $\Omega_c(3050)$ and $\Omega_c(3119)$ states with the $(J = 1, 1/2^+, 3/2^+)$ states from the exotic $\overline{\mathbf{15}}$ multiplet. The remaining two $(J = 2, 3/2^-, 5/2^-)$ states from the sextet have masses above the $\Xi + D$ threshold (3185 MeV), so they are probably hidden in a large bump observed by the LHCb Collaboration above 3200 MeV. It should be stressed that the simplest scenario in which all five LHCb Ω_c^0 states are classified as members of the excited sextets contradicts general mass formulas derived within the χQSM.

The simplest way to falsify our identification is to search for the *isospin* partners of Ω_c^0 from the $\overline{\mathbf{15}}$. For example, they can be searched in the mass distribution of $\Xi_c^0 + K^-$ or $\Xi_c^+ + \bar{K}^0$; the Ω_c^0's from the sextet do not decay into these channels.

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