\[ \Omega_b^- \to (\Xi^+_c K^-) \pi^- \] and the \( \Omega_c \) states

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Abstract

We study the weak decay \( \Omega_b^- \to (\Xi^+_c K^-) \pi^- \), in view of the narrow \( \Omega_c \) states recently measured by the LHCb collaboration and later confirmed by the Belle collaboration. The \( \Omega_c(3050) \) and \( \Omega_c(3090) \) are described as meson-baryon molecular states, using an extension of the local hidden gauge approach in coupled channels. We investigate the \( \Xi D, \Xi_c \bar{K} \) and \( \Xi'_c \bar{K} \) invariant mass distributions making predictions that could be confronted with future experiments, providing useful information that could help determine the quantum numbers and nature of these states.

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I. INTRODUCTION

The recent discovery of five narrow $\Omega_c$ states by the LHCb collaboration \[1\] in $pp$ collisions, also recently confirmed by the Belle collaboration \[2\] in $e^+e^-$ collisions, motivated an increasing amount of theoretical work with different proposals for their structure. In particular, the correct assignment of quantum numbers $J^P$ remains an open question and it could be the key to understand the nature of these states.

Predictions using quark models for such states and related ones were done in Refs. \[3–15\]. Quark models mostly propose a diquark-quark structure $(ss)c$, where it is assumed that the diquark $(ss)$ is in the ground state $1S$ bound by the attractive antitriplet color structure $3$; and in this case the identical flavor content of the diquark implies that it has spin $S_{ss} = 1$, which then can combine with the spin of the $c$ quark $S_c$ and with orbital momentum $L$ of the diquark $(ss)$ relative to the $c$ quark. One of the typical assignments in the literature consists in the five $1P$ possibilities of combining angular momentum when $L = 1$, which yields two states with $J^P = 1/2^-$, two with $3/2^-$ and one with $5/2^-$, that split due to spin-dependent forces and could correspond to the five narrow $\Omega_c$ states: $\Omega_c(3000)$, $\Omega_c(3050)$, $\Omega_c(3066)$, $\Omega_c(3090)$ and $\Omega_c(3119)$, in the same order of increasing $J$. Alternative assignments also include the possibility of a radial excitation of the diquark $(ss)$ relative to the $c$ quark, where the last two $\Omega_c$ states would be $2S$ excitations with quantum numbers $J^P = 1/2^+$ and $3/2^+$, respectively, whereas the first three $\Omega_c$ states would be the last $1P$ states with quantum numbers $J^P = 3/2^-$ and $5/2^-$, in that order. Interesting discussions on this quark model picture can be found in Ref. \[16\] and similar approaches in Refs. \[17–20\].

Other methods have also been employed to study these states, as QCD Sum Rules in Refs. \[21–27\] and Lattice QCD in Ref. \[28\]. Pentaquark options have been suggested in Refs. \[29–34\]. Some works have emphasized the value of decay properties to obtain information on the nature of these states \[35–37\] and a discussion on the possible quantum numbers was done in Ref. \[38\].

On the other hand, some of these states could actually be pentaquark-like molecules, dynamically generated from meson-baryon interactions in coupled channels with charm $C = 1$, strangeness $S = -2$ and isospin $I = 0$. Predictions in the molecular picture using coupled channels of meson-baryon interactions were done in Refs. \[39–41\]. In this picture the interaction in $S$-wave of baryons with spin-parity $J^P = 1/2^+$ or $J^P = 3/2^+$ with pseudoscalar
mesons leads to meson-baryon systems with $J^P = 1/2^-$ and $J^P = 3/2^-$, respectively. Channels with vector mesons instead of pseudoscalars can also be included resulting in $J^P = 1/2^-$, $3/2^-$ and $5/2^-$. However, most of the recent works adopting this picture manage to relate two or three of the new $\Omega_c$ states to meson-baryon systems with $J^P = 1/2^-$ and $J^P = 3/2^-$, dominated by the pseudoscalar-baryon channels.

In Ref. [41] an $SU(6)_{isf} \times HQSS$ model (HQSS stands for heavy quark spin symmetry) extending the Weinberg-Tomozawa $\pi N$ interaction was employed to make a systematic study of many possible meson-baryon systems. In Ref. [42] the renormalization scheme of Ref. [41] was reviewed, performing an update of the results of the $C = 1, S = -2$ and $I = 0$ sector in view of the new experimental data. The updated results indicate that one can relate the $\Omega_c(3000)$ to a state with $J^P = 1/2^-$ and the $\Omega_c(3050)$ to another state with $J^P = 3/2^-$, with hints that the $\Omega_c(3090)$ or $\Omega_c(3119)$ could also have $J^P = 1/2^-$. In Ref. [40] the molecular picture was developed using $SU(4)$ symmetry to extend the interaction described by vector meson exchange in the local hidden gauge approach. This work was also reviewed under the light of the new experimental data and an updated study was made in Ref. [43], where it was shown that the $\Omega_c(3050)$ and $\Omega_c(3090)$ can be both related to meson-baryon resonances with $J^P = 1/2^-$, stemming from pseudoscalar-baryon$(1/2^+)$ interaction.

A similar approach which also describes the meson-baryon interaction through vector meson exchange was recently developed in Ref. [44], using an extension of the local hidden gauge approach [45–49] and taking into account the spin-flavor wave functions of the baryons. In the present work we will follow the description of the $\Omega_c$ states of Ref. [44]; extensive discussions on the methods and results can be found there, as well as predictions of higher energy states of meson-baryon nature. In this framework, which will be discussed in more detail in the next section, the heavy quarks are treated as spectators, what implies that the dominant terms of the interaction come from light vector exchange, therefore the interaction is consistent with $SU(3)$ chiral Lagrangians and heavy quark spin symmetry is preserved for the dominant terms in the $(1/m_Q)$ counting [50–52]. A remarkable agreement of both masses and widths of the $\Omega_c(3050)$ and $\Omega_c(3090)$ was obtained from the pseudoscalar-baryon$(1/2^+)$ interaction, in accordance with results of Ref. [43]; also an extra sector of pseudoscalar-baryon$(3/2^+)$ could be related to the $\Omega_c(3119)$, therefore assigning to it the quantum numbers $J^P = 3/2^-$. 

3
Other works on the molecular picture followed [53–55]. In Ref. [53] the authors propose that the broad structure found around 3188 MeV [1, 2] could be related with a molecular $\Xi D$ state due to the proximity of its threshold around 3185 MeV. As we shall see in the next section, in our approach [44] the molecular state dominated by the $\Xi D$ channel corresponds to the $\Omega_c(3090)$.

It is clear that any constraint on the quantum numbers of these new states is essential to discriminate between the different theoretical models. It is possible — or if we dare to say, likely — that some of the new $\Omega_c$ states are $1P$ and/or $2S$ excitations as discussed in the quark model picture, and some are dynamically generated meson-baryon molecules. However, answering these questions would automatically open new ones, since in both pictures, choosing one assignment over the other implies that there should be other partner states, which could be below or above the energy range searched experimentally; some of them might be detectable only in other decay channels. Understanding the properties of the observed states, like the narrowness, the couplings to one channel or another, the presence or absence of certain predicted states, and so on, require more information, not just on the quantum numbers, but also the measurement of such states in different reactions. Radiative transitions could play an important role on this search, such as the confirmation of the feeddown mechanism $\Omega_c(3066) \rightarrow \Xi_c' K \rightarrow \Xi_c \gamma K$, which might be the cause of the enhancement near threshold on the LHCb data [1], even tough additional states have not been completely ruled out as an alternative explanation. The search for isospin-violating decays to $\Omega_c \pi^0$ or electric dipole transitions to $\Omega_c \gamma$ could also bring valuable new information. In view of this exciting spectroscopy challenge, where any correct explanation brings with itself an equal or greater number of new questions, we will try to shed light on the discussion looking for different reactions where the new $\Omega_c$ states could be spotted, and discuss peculiarities of each case that could help distinguish among the various different interpretations on the recent literature.

In the present work we propose the experimental study of these new states through the decay of $\Omega_c^-$ baryons, as suggested in Ref. [56]. The mass and lifetime of the $\Omega_c^-$ were recently measured by the LHCb collaboration [57], obtaining results compatible with the previous measurements of the same collaboration [58] and also with the ones of the CDF collaboration [59], but not with the results of the DØ collaboration [60]. We shall adopt the mass value listed by the Particle Data Group [61], which is quite close to the most recent
measurement of LHCb.

We will discuss the decay $\Omega^{-}_b \to (\Xi^{+}_c K^-)\pi^-$, which could be performed by the LHCb collaboration \cite{56} in the future and could be very useful to distinguish states with different structures and quantum numbers. First we present a brief summary of the main results of Ref. \cite{44} and comment on the formalism employed there to obtain the $\Omega_c(3050)$ and $\Omega_c(3090)$ as meson-baryon molecules (further details can be found in the Appendix section). Next we will discuss the $\Omega^{-}_b \to (\Xi^{+}_c K^-)\pi^-$ decay and how the coupled channels approach naturally accounts for the dynamical generation of the $\Omega_c$ states from the hadronization that takes place after the conversion of the $b$ quark into a $c$ quark. Then we show the results of how these two states would be seen in the $\Omega^{-}_b$ decay if the molecular picture of Ref. \cite{44} is correct, providing solid predictions that could easily be put to proof in the near future.

\section{The $\Omega_c(3050)$ and $\Omega_c(3090)$ in the Molecular Picture}

In Ref. \cite{44} a thorough discussion was made about the meson-baryon interaction due to the exchange of vector mesons, and an extension of the local hidden gauge approach \cite{45–49} was used together with a method that takes into account the information of the spin-flavor wave functions of the baryons (see Appendix and Ref. \cite{44} for details). It was shown that considering the heavy quarks as spectators, the interactions can be obtained using only $SU(3)$ symmetry (without the need of $SU(4)$ in the dominant terms), respecting heavy quark spin symmetry in the leading terms where only light vectors are exchanged \cite{50–52}.

The procedure begins with the choice of meson-baryon channels in the $C = 1, S = -2$ and $I = 0$ sector, interacting in $S$-wave, with a determined total spin $J$. In the case of $J = 1/2$, we can have the pseudoscalar-baryon interactions, where the baryons have $J = 1/2$, and also vector-baryon contributions that result in $J = 1/2$. However, in Ref. \cite{44} it was shown that the coupling of vector-baryon channels with the pseudoscalar-baryon ones gives only a small contribution that can be safely neglected. Therefore, the states generated from the vector-baryon interaction decouple from the ones with pseudoscalars and can be treated separately. Since they do not play any role in the generation of the $\Omega_c(3050)$ and $\Omega_c(3090)$ in our approach, we will leave them aside in the present work. As presented in Ref. \cite{44}, a third state, the $\Omega_c(3119)$, can also be related with another pseudoscalar-baryon system,
where the baryons have $J = 3/2$, assigning the quantum numbers $J^P = 3/2^-$ to the $\Omega_c(3119)$. However, since in our approach this state does not mix with the channels of $J = 1/2$, its decay into $\Xi_c\bar{K}$ calls for additional mechanisms with the exchange of pseudoscalars, which go beyond the scope of our approach here.

The channels chosen for the $J^P = 1/2^-$ pseudoscalar-baryon states and their respective threshold masses are shown in Table I.

| States | $\Xi_c\bar{K}$ | $\Xi_c^*\bar{K}$ | $\Xi D$ | $\Omega_c\eta$ |
|--------|----------------|-----------------|--------|---------------|
| Threshold | 2965 | 3074 | 3185 | 3243 |

Using an extension of the local hidden gauge Lagrangians [45–49] and the baryon spin-flavor wave functions with the heavy quark as spectator, the pseudoscalar-baryon interaction of each channel due to the exchange of vector mesons was calculated in Ref. [44], and then used as the kernel of the Bethe-Salpeter equation in the on-shell approximation where the multiple meson-baryon loops are summed up in order to obtain a unitarized scattering amplitude [62, 63]

$$ T = [1 - VG]^{-1} V, $$

where $V$ is the matrix describing the interaction between every channel (see Appendix for more details) and $G$ is the meson-baryon loop function

$$ G_l = i \int \frac{d^4q}{(2\pi)^4} \frac{M_l}{E_l(q)} \frac{1}{k^0 + p^0 - q^0 - E_l(q) + i\epsilon} \frac{1}{q^2 - m_l^2 + i\epsilon}, $$

where $k^0 + p^0 = \sqrt{s}$ and $\omega_l, E_l$, are the meson and baryon energies in the $l$-channel, respectively, and $m_l, M_l$ the meson and baryon masses. As discussed in Ref. [44] a common sharp cutoff of 650 MeV was used to regularize the propagators. We will adopt the same cutoff value along the present work.

When solving Eq. (1) as a function of the center-of-mass energy $\sqrt{s}$, one can go to the complex energy plane and look for poles in the second Riemann sheet, which correspond to the dynamically generated states such that the real part of the pole is the mass of the state and its width is given by twice the imaginary part of the pole. In Table II we show the
poles found that correspond to the $\Omega_c(3050)$ and $\Omega_c(3090)$ states. In addition, we evaluate the couplings $g_i$ of the states obtained to the different channels and $g_i G_i$, which for $S$-wave gives the strength of the wave function at the origin [44, 64, 65].

TABLE II. The coupling constants to various channels for the poles in the $J^P = 1/2^-$ sector, with $q_{max} = 650$ MeV, and $g_i G_i^H$ in MeV.

|       | $\Xi_c \bar{K}$ | $\Xi'_c \bar{K}$ | $\Xi D$ | $\Omega_c \eta$ |
|-------|-----------------|-----------------|--------|-----------------|
| 3054.05 + i0.44 | -0.06 + i0.14 | 1.94 + i0.01 | -2.14 + i0.26 | 1.98 + i0.01 |
| $g_i$   | -1.40 - i3.85   | -34.41 - i0.30 | 9.33 - i1.10 | -16.81 - i0.11 |
| $g_i G_i^H$ | 5.05 + i10.19  | -9.97 - i3.67  | -29.82 + i0.31 | -3.59 - i2.23 |
| 3091.28 + i5.12 | 0.18 - i0.37  | 0.31 + i0.25  | 5.83 - i0.20  | 0.38 + i0.23  |
| $g_i$   | 0.18 - i0.37    | 0.31 + i0.25   | 5.83 - i0.20  | 0.38 + i0.23  |
| $g_i G_i^H$ | 5.05 + i10.19  | -9.97 - i3.67  | -29.82 + i0.31 | -3.59 - i2.23 |

III. THE $\Omega_b^− \to (\Xi_c \bar{K}) \pi^−$, $\Omega_b^− \to (\Xi'_c \bar{K}) \pi^−$ AND $\Omega_b^− \to (\Xi D) \pi^−$ DECAYS

As discussed in the previous sections and extensively in Ref. [44], the $\Omega_c(3050)$ and $\Omega_c(3090)$ are both strong candidates of meson-baryon molecules, dynamically generated from the coupled channels interaction of baryons with $J^P = 1/2^+$ and pseudoscalar mesons ($J^P = 0^−$) in $S$-wave, therefore in this picture their spin-parity assignment is $J^P = 1/2^−$. Since every quark has intrinsic $J^P = 1/2^+$, in order to have $\Xi_c^+ K^−$ in $S$-wave in the final state, the hadronization must occur between the $c$ quark and one $s$ quark, as shown in Fig. 1.

The reason for this approach is that the $b$ quark must turn into a $c$ quark, then both $s$ quarks act as spectators in this decay, so their quantum numbers are fixed a priori as $J^P = 1/2^+$, what implies that the $c$ quark must be in $L = 1$ to account for the negative parity of the final state, and the hadronization must necessarily involve it so it can be deexcited (see Ref. [66] for details). That being said, we include the $\bar{q}q$ terms in such a way that the hadronization will occur between the $c$ quark and an $s$ quark.

$$css \to c(\bar{u}u + \bar{d}d + \bar{s}s)s \equiv H, \quad (3)$$

$$H = \sum_i c \bar{q}_i q_i ss \equiv \sum_i \Phi_{4i} q_i ss, \quad (4)$$
FIG. 1. $\Omega^{-}_b$ decay at quark level with emission of a $\pi^-$ and subsequent hadronization.

where in the last step we have identified the $(q\bar{q})$ matrix in Eq. (5) with the matrix $\Phi$ of pseudoscalar mesons in Eq. (6),

$$(q\bar{q}) = \begin{pmatrix} u\bar{u} & u\bar{d} & u\bar{s} & u\bar{c} \\ d\bar{u} & d\bar{d} & d\bar{s} & d\bar{c} \\ s\bar{u} & s\bar{d} & s\bar{s} & s\bar{c} \\ c\bar{u} & c\bar{d} & c\bar{s} & c\bar{c} \end{pmatrix} \equiv \Phi, \quad (5)$$

$$\Phi = \begin{pmatrix} \frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{3}} \eta + \frac{1}{\sqrt{6}} \eta' \\ -\frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{3}} \eta + \frac{1}{\sqrt{6}} \eta' \\ \pi^+ \\ K^+ \\ D^0 \end{pmatrix} \begin{pmatrix} \pi^- & K^0 & D^- \\ K^- & D^0 & D_s^0 \end{pmatrix}, \quad (6)$$

where we have included the mixing between $\eta$ and $\eta'$ [67] for proper matching of $\Phi$ with the $q\bar{q}$ matrix.

Then replacing the $c\bar{q}$ meson terms we get

$$|H\rangle = D^0 \, u\bar{s}s + D^+ d\bar{s}s + \ldots \quad (7)$$

where we have already neglected the heavy combination of $D_s^+ sss$ which could only contribute to states with $J^P = 3/2^-$ since $sss$ corresponds to the $\Omega^-$. In terms of spin-flavor wave functions, the ground state $\Omega^-$ is flavor symmetric, hence it can only have symmetric spin $\uparrow\uparrow\uparrow$, i.e. $J^P = 3/2^+$, since the color singlet provides the antisymmetric part of the wave function. Besides, $D_s^+ \Omega^-$ is far above (threshold at 3641 MeV) the energy range of the channel space considered to generate the meson-baryon molecules.
Next, we need to relate the \( \text{uss} \) and \( \text{dss} \) with the corresponding baryons, in this case, the \( \Xi \) baryons with \( J^P = 1/2^+ \). For consistence with our approach in Ref. [44], we use the proper spin-flavor wave functions, inspired in Ref. [68], where

\[ \Xi^0 \equiv \frac{1}{\sqrt{2}} (\phi_{\text{MS}} \chi_{\text{MS}} + \phi_{\text{MA}} \chi_{\text{MA}}), \]  

with the Mixed-Symmetric and Mixed-Antisymmetric flavor and spin wave functions\(^1\):

\[ \phi_{\text{MS}} = \frac{1}{\sqrt{6}} [s(us + su) - 2uss], \quad \phi_{\text{MA}} = -\frac{1}{\sqrt{2}} [s(us - su)], \]

\[ \chi_{\text{MS}} = \frac{1}{\sqrt{6}} (\uparrow\uparrow\downarrow + \uparrow\downarrow\uparrow - 2 \downarrow\uparrow\uparrow), \quad \chi_{\text{MA}} = \frac{1}{\sqrt{2}} \uparrow (\uparrow\downarrow - \downarrow\uparrow). \]

Thus, only the mixed symmetric will contribute and we get the following weight for \( \Xi^0 \)

\[ \langle \text{uss} | \Xi^0 \rangle = \frac{1}{\sqrt{6}} \langle \text{uss} | sus + ssu - 2uss \rangle \]

\[ \equiv -\frac{2}{\sqrt{6}}. \]  

(9)

Equivalently, for the \( \Xi^- \) wave function we only replace \( u \rightarrow d \) and change the sign of \( \phi_{\text{MS}} \) and \( \phi_{\text{MA}} \) [44]. Again, only the mixed symmetric part will contribute and we get

\[ \langle \text{dss} | \Xi^- \rangle = -\frac{1}{\sqrt{6}} \langle \text{dss} | sds + ssd - 2dss \rangle \]

\[ \equiv \frac{2}{\sqrt{6}}. \]  

(10)

Note that we have dropped the common factor \( 1/\sqrt{2} \) from Eq. (8), which can be absorbed in the constant factor \( V_P \) that we will introduce next.

Then after the hadronization we will have the combination

\[ |H\rangle = -\frac{2}{\sqrt{6}} D^0 \Xi^0 + \frac{2}{\sqrt{6}} D^+ \Xi^- = \frac{2}{\sqrt{3}} |\Xi D, I = 0\rangle, \]  

(11)

where we have absorbed the global minus sign when we changed to the isospin 0 combination in the order baryon-meson\(^2\):

\[ |\Xi D, I = 0\rangle = -\frac{1}{\sqrt{2}} \left| \Xi^0 D^0 - \Xi^- D^+ \right\rangle. \]  

(12)

\(^1\) Note that the phase convention of \( \Xi^0 \) is changed in respect to Ref. [68]. This is necessary to be consistent with the chiral Lagrangians, as in Table III of Ref. [69]. See also the footnote of Ref. [70].

\(^2\) Recall the isospin doublets:

\[ D = \begin{pmatrix} D^+ \\ -D^0 \end{pmatrix}, \quad \Xi = \begin{pmatrix} \Xi^0 \\ -\Xi^- \end{pmatrix}. \]
Now we can proceed to construct the amplitude of the process $\Omega_b^- \to (\Xi_c K) \pi^-$. It is instructive to first look at the process $\Omega_b^- \to (\Xi D) \pi^-$, as depicted in Fig. 2. From Eq. (11) we see that after the emission of a pion, the hadronization involving the $css$ quarks generates a $\Xi D$ pair in isospin 0. Thus we can write this process as the sum of a tree-level contribution and the final state interaction of $\Xi D$ going through the molecular states of Table II. This information is contained in the diagonal $t$ matrix element $t^{I=0}_{\Xi D \to \Xi D}$, with $t$ the same as in Eq. (1), where all the coupled channels dynamics is included.

\begin{equation}
\Omega_b^- \to \pi^- \Xi D \quad \text{process. Tree-level (left) plus } \Xi D \text{ rescattering (right).}
\end{equation}

Then for the tree-level contribution (left diagram of Fig. 2) we simply write

\begin{equation}
t_{\text{tree}} = V_P \frac{2}{\sqrt{3}},
\end{equation}

where the weight of $\Xi D$ production comes from Eq. (11) and $V_P$ contains all information related to the $\Omega_b^-$ weak decay, a common unknown factor in all process we will investigate. On the other hand, for the $\Xi D$ rescattering (right diagram of Fig. 2) we will have

\begin{equation}
t_{\text{loop}} = V_P \frac{2}{\sqrt{3}} G_{\Xi D} [M_{\text{inv}}(\Xi D)] t^{I=0}_{\Xi D \to \Xi D} [M_{\text{inv}}(\Xi D)],
\end{equation}

where $G_{\Xi D}$ is the propagator of the baryon-meson loop, as in Eq. (2). Then the amplitude of the process $\Omega_b^- \to \pi^- \Xi D$ is given by

\begin{equation}
t_{\Omega_b^- \to \pi^- \Xi D} = V_P \frac{2}{\sqrt{3}} (1 + G_{\Xi D} [M_{\text{inv}}(\Xi D)] t^{I=0}_{\Xi D \to \Xi D} [M_{\text{inv}}(\Xi D)]) .
\end{equation}

With this amplitude we can write the $\Xi D$ invariant mass distribution

\begin{equation}
\frac{d\Gamma}{dM_{\text{inv}}(\Xi D)} = \frac{1}{(2\pi)^3} \frac{1}{4M_{\Omega_b^-}^2} 2M_{\Omega_b^-} 2M_{\Xi} p_{\pi^-} p_D \left| t_{\Omega_b^- \to \pi^- \Xi D} \right|^2 ,
\end{equation}

10
where \( p_{\pi^-} \) is the pion momentum in the \( \Omega_b^- \) rest frame

\[
p_{\pi^-} = \frac{\lambda^{1/2}(\lambda, m_\pi^2, M_{\Omega b}^2)}{2M_{\Omega b}^-},
\]

(17)

and \( \tilde{p}_D \) is the \( D \) momentum in the \( \Xi D \) rest frame

\[
\tilde{p}_D = \frac{\lambda^{1/2}(\lambda, m_D^2, m_\Xi^2)}{2M_{\Xi D}^{-}}.
\]

(18)

For the \( \Omega_b^- \rightarrow (\Xi_c \bar{K}) \pi^- \) process there is no tree-level contribution, since the hadronization only produces a \( \Xi D \) pair, as in Eq. (11), then the only contribution comes from the diagram in Fig. 3.

Due to our coupled channels approach, the transition \( \Xi D \rightarrow \Xi_c \bar{K} \) is already contained in the \( t \) matrix and the production of \( \Xi_c \bar{K} \) (also in isospin 0) appears naturally. The corresponding amplitude will be

\[
t_{\Omega b^- \rightarrow \pi^- \Xi_c \bar{K}} = V_P \frac{2}{\sqrt{3}} G_{\Xi D}[M_{\Xi D}(\Xi_c \bar{K})] t_{\Xi D \rightarrow \Xi_c \bar{K}}[M_{\Xi D}(\Xi_c \bar{K})],
\]

(19)

where \( t_{\Xi D \rightarrow \Xi_c \bar{K}} \) is the transition amplitude of \( \Xi D \rightarrow \Xi_c \bar{K} \), from the same \( t \) matrix. Then the \( \Xi_c \bar{K} \) invariant mass distribution is also analogous,

\[
\frac{d\Gamma}{dM_{\Xi D}(\Xi_c \bar{K})} = \frac{1}{(2\pi)^3} \frac{4M_{\Omega b}^-}{2M_{\Omega b}^-} \frac{2M_{\Xi_c}}{2M_{\Xi_c} \rho_{\pi^-} \tilde{p}_{\Xi_c}} \left| t_{\Omega b^- \rightarrow \pi^- \Xi_c \bar{K}} \right|^2,
\]

(20)

where

\[
p_{\pi^-} = \frac{\lambda^{1/2}(\lambda, m_\pi^2, M_{\Xi c \bar{K}}^2)}{2M_{\Omega b}^-},
\]

(21)
and
\[ \bar{p}_K = \frac{\lambda^{1/2}(M_{\text{inv}}(\Xi_c \bar{K}), m_{\bar{K}}^2, M_{\Xi_c}^2)}{2M_{\text{inv}}(\Xi_c \bar{K})}. \] (22)

Analogously, we can also calculate the invariant mass distribution for the final state $\Xi'_c \bar{K}$, replacing $\Xi_c$ by $\Xi'_c$ in the previous equations and taking the matrix element of the $t$ matrix corresponding to the transition $\Xi D \rightarrow \Xi'_c \bar{K}$.

It is also interesting to look at the case of coalescence, where the $\Xi D$ pair merge into the resonance regardless of the final decay channel, as depicted in Fig. 4.

FIG. 4. Resonance coalescence in the $\Omega_b^- \rightarrow \pi^- R_i$ process through $\Xi D$ rescattering, where $R_i$ is the $\Omega_c(3050)$ or $\Omega_c(3090)$.

The value of the amplitude in the process $\Omega_b^- \rightarrow \pi^- R_i$, where $R_i$ is one of the molecular states of Table II, is proportional to the coupling of that resonance to the $\Xi D$ channel,
\[ t_{\Omega_b^- \rightarrow \pi^- R_i} = V_P \frac{2}{\sqrt{3}} G_{\Xi D}(M_{R_i}) g_{R_i, \Xi D}, \] (23)
where the propagator is calculated at the resonance mass $M_{R_i}$. With this quantity we can calculate the equivalent of the integrated mass distribution around the $R_i$ resonance, which does not depend on its decay mode,
\[ \Gamma_{\Omega_b^- \rightarrow \pi^- R_i} = \frac{1}{8\pi} \frac{1}{M_{\Omega_b^-}^2} 2M_{\Omega_b^-} 2M_{R_i} \bar{p}'_{\pi^-} \left| t_{\Omega_b^- \rightarrow \pi^- \Xi_c \bar{K}}(M_{R_i}) \right|^2, \] (24)
where
\[ \bar{p}'_{\pi^-} = \frac{\lambda^{1/2}(M_{\Omega_b^-}^2, m_{\pi^-}^2, M_{R_i}^2)}{2M_{\Omega_b^-}}. \] (25)
IV. RESULTS

It is interesting to see how these processes show the importance of the coupled channels. The decay of the $\Omega_c$ states into $\Xi D$ is kinematically forbidden below the corresponding threshold at 3185 MeV, then we cannot see the corresponding peaks in the $\Xi D$ invariant mass distribution, but we can see their indirect effect, both from the meson-baryon loop in Eq. (15) and from the amplitude $t^{I=0}_{\Xi D \rightarrow \Xi D}$. The corresponding invariant mass distribution, Eq. (16), is shown in Fig. 5 by the solid line. To compare with the case where only tree-level contributes, we remove $t_{\text{loop}}$ and keep only $t_{\text{tree}}$, normalizing the curve such that it has the same area as the solid curve in the energy range shown, which is plotted as the dashed line in Fig. 5.

![Graph showing the invariant mass distribution of $\Xi D$](image)

**FIG. 5.** $\Xi D$ invariant mass distribution from Eq. (16). Solid line: using the complete amplitude of Eq. (15). Dashed line: removing the $G_{\Xi D} t^{I=0}_{\Xi D \rightarrow \Xi D}$ term (only tree-level contribution) and normalizing such that both curves have the same area.

The $\Xi_c K$ threshold is at 2965 MeV, then we can see clearly, of course, the peaks of the $\Omega_c$ states in the $\Xi_c K$ invariant mass distribution. According to Eq. (11), we expect only $\Xi D$ production from the hadronization that occurs right after the $\Omega_b^-$ decay, which means we have no tree-level contribution for $\Xi_c K$ production. However, the transition to $\Xi_c K$ through off-shell $\Xi D$ loops arises naturally from the coupled channels approach. In fact, both the $\Omega_c(3050)$ and $\Omega_c(3090)$ couple strongly to $\Xi D$ (see Table II), and their formation from the
ΞD state formed in the first step of the Ω_b^− decay with subsequent transition to Ξ_cK (going through the ΞD virtual state) is not only possible, but expected.

In Fig. 6 we show the Ξ_cK invariant mass distribution. The only unknown quantity is the global factor V_P, common to all amplitudes we investigate here. The ratio between the intensity of each peak does not depend on V_P, so all ratios are predictions that could be confronted with future experiments. We can see that the intensity of the Ω_c(3050) peak is about 65% higher than the Ω_c(3090) peak. Note that we are using the same normalization for all the reactions, hence the ratio of the strength at the peaks in Fig. 6 to the strength of the ΞD mass distribution (solid line) in Fig. 5 is also a prediction. In arbitrary units, the ΞD distribution has a maximum of about 125 around 3240 MeV, whereas in the Ξ_cK distribution the Ω_c(3050) and Ω_c(3090) peaks have an intensity of about $15.50 \times 10^3$ and $9.45 \times 10^3$, respectively, therefore we predict that the Ξ_cK distribution in the vicinity of the resonances peaks is roughly two orders of magnitude higher than the ΞD distribution.

![Graph of Ξ_cK invariant mass distribution from Eq. (20).](image)

Cusp effects in the Ξ_cK distribution also appear at the Ξ'_cK and ΞD thresholds at 3074 MeV and 3185 MeV, respectively, but their intensity is very small compared to the peaks of the resonances and cannot be seen clearly in Fig. 6.

We also notice that apparently no significant interference pattern is seen between both states, even tough they have the same quantum numbers, a feature that also agrees with
the fit performed by the LHCb collaboration [1, 16].

As for the $\Xi_c'K$ invariant mass distribution, only the $\Omega_c(3090)$ can be seen, as shown in Fig. 7, since this channel is also open for the decay of this state, whereas the $\Omega_c(3050)$ is below the threshold of $\Xi_c'K$.

![Fig. 7. $\Xi_c'K$ invariant mass distribution from Eq. (20), replacing $\Xi_c$ by $\Xi_c'$ and taking $t_{\Xi D \to \Xi_c'K}$.](image)

Again we can compare the intensity of the peaks. In the $\Xi_c'K$ distribution the $\Omega_c(3090)$ peak has an intensity of $3.86 \times 10^3$, which is about 40% of the intensity it has in the $\Xi_cK$.

Finally, it is interesting to compare the results of $\Gamma_{R_i \to \pi^- R_i}$, given by Eq. (24), with the integrated invariant mass distribution of Eq. (20) around the peak of each resonance,

$$\int_{R_i} \frac{d\Gamma}{dM_{\text{inv}}(\Xi_cK)} dM_{\text{inv}}(\Xi_cK).$$

In Table III we show the results of Eqs. (24) and (26) for the $\Omega_c(3050)$ and $\Omega_c(3090)$, and for the latter we also show the integrated $\Xi_c'K$ distribution.

From these results we can draw some interesting conclusions. Let us look first at the $\Omega_c(3050)$. In Table II we can see that this state is dominated by the $\Xi_c'K$ channel, and also has sizable contributions from the higher channels $\Xi D$ and $\Omega_c\eta$. However, the pole is at 3054 MeV, which is 20 MeV below the $\Xi_c'K$ threshold, so the only open channel is $\Xi_cK$, to which the resonance couples very weakly and the phase space available is only about 90 MeV. This feature explains two points: 1) The narrowness of the state, whose upper limit of the width
TABLE III. Comparison between integrated invariant mass distributions around each resonance and the corresponding coalescence.

| State | $\Omega_c(3050)$ | $\Omega_c(3090)$ |
|-------|------------------|------------------|
| Pole  | $3054.05 + i0.44$ | $3091.28 + i5.12$ |
| Coalescence Eq. (24) [a. u.] | 21289 | 215237 |
| Channel | $\Xi_c\bar{K}$ | $\Xi_c\bar{K}$ | $\Xi'_c\bar{K}$ |
| Interval [MeV] | [3049, 3057] | [3057, 3120] | [3074, 3120] |
| Integral Eq. (26) [a. u.] | 21344 | 133482 | 51074 |

reported by the LHCb collaboration is $0.8 \pm 0.2 \pm 0.1$ MeV [1], in excellent agreement with our result: $\Gamma = 2 \times \text{Im}(R_i) = 0.88$ MeV; 2) The good agreement of $\Gamma_{\Omega_c \rightarrow \pi^- R_i}$, 21289, given by Eq. (24), with the integrated invariant mass distribution around the resonance peak, 21344, given by Eq. (26) (the small difference is irrelevant and comes essentially from the choice of the interval of integration). This happens because the state is very narrow and the only channel open for decay is the $\Xi_c\bar{K}$, so the value we obtain from the coalescence (which is independent of the decay channel), matches the value obtained from the integration over the only channel available ($\Xi_c\bar{K}$) as it should be.

On the other hand, the $\Omega_c(3090)$ is dominated by the $\Xi D$ channel, with some contribution from the other channels. Even tough this state couples very strongly to $\Xi D$, it is almost 100 MeV below the respective threshold. It can only decay to $\Xi_c\bar{K}$ and $\Xi'_c\bar{K}$. In both cases the coupling is small, and in the latter channel the phase space available is less than 20 MeV (see Table II). This again, explains two main features: 1) The narrowness, with a width not so small as in the case of the $\Omega_c(3050)$, but still very narrow, since the decay into $\Xi_c\bar{K}$ is reasonable, with more than 120 MeV of phase space available, although the coupling to that channel is weak. The LHCb reports a width of $8.7 \pm 1.0 \pm 0.8$ MeV [1], in fair agreement with our result of 10.24 MeV; 2) The fact that the integrated invariant mass distribution around the peak, 133482 in the $\Xi_c\bar{K}$ distribution, is about 2/3 of the total given by the coalescence, 215237, which is also expected, since in Eq. (26) we are integrating only in the $\Xi_c\bar{K}$ channel, whereas this state can also decay into $\Xi'_c\bar{K}$. The sum of both integrals, in $\Xi_c\bar{K}$ and $\Xi'_c\bar{K}$, is 184556, close to the total given by the coalescence, but still below. This is also expected since Eq. (24) is actually an approximation that is valid in the limit of zero width, which
works pretty well for the $\Omega_c(3050)$ but is not so good for the $\Omega_c(3090)$, with already 10 MeV of width. This can also be verified calculating the contribution of each channel to the total width, in terms of the coupling of the resonance to that channel. For $\Xi_cK$

$$\Gamma_{R,\Xi_cK} = \frac{1}{8\pi} \frac{1}{M_{R_i}^2} 2M_{R_i} 2M_{\Xi_c} |g_{R_i,\Xi_cK}|^2 \tilde{p}_2,$$

(27)

where $M_{R_i}$ is the mass of the resonance and $\tilde{p}_2$ is given by

$$\tilde{p}_2 = \frac{\lambda^{1/2}(M_{R_i}^2, m_K^2, M_{\Xi_c}^2)}{2M_{R_i}}.$$  

(28)

Using this equation we get $\Gamma_{\Xi_cK} = 0.83$ MeV for the $\Omega_c(3050)$, in accordance with the value obtained from the pole position. For the $\Omega_c(3090)$ we get $\Gamma_{\Xi_cK} = 7.24$ MeV for the $\Xi_cK$, and replacing $\Xi_c$ by $\Xi_c'$ and the corresponding coupling, we get $\Gamma_{\Xi'_cK} = 2.56$ MeV, which add up to 9.80 MeV, a bit below 10.24 MeV that we get from the pole position.

As a prediction we can also state, based on the coalescence results, that the ratio of the $\Omega_c(3050)$ over the $\Omega_c(3090)$ production is about 10% in the $\Omega_b^-$ decay:

$$\frac{\Gamma_{\Omega_b^+ \to \pi^- \Omega_c(3050)}}{\Gamma_{\Omega_b^- \to \pi^- \Omega_c(3090)}} \approx 10\%.$$ 

(29)

V. CONCLUSIONS

We have studied the weak decay $\Omega_b^- \to (\Xi_c^+ K^-) \pi^-$, in view of the narrow $\Omega_c$ states recently measured by the LHCb collaboration and later confirmed by the Belle collaboration. Based on the previous work where the $\Omega_c(3050)$ and $\Omega_c(3090)$ are described as meson-baryon molecular states, using an extension of the local hidden gauge approach in coupled channels, with results in remarkable agreement with experiment, we have investigated the $\Xi D$, $\Xi_cK$ and $\Xi'_cK$ invariant mass distributions, and discussed the role of coupled channels in the process. Predictions that could be confronted with future experiments are presented, providing useful information that could help to determine the quantum numbers and nature of these states. Since $\Omega_b^-$ baryons have already been observed in several experiments, two of them performed by the LHCb collaboration, the present work should encourage such study in the near future, which would certainly bring novel key information for the understanding of these new states.
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APPENDIX

The ingredients needed to obtain the molecular states presented here are the vector(V)-pseudoscalar(P)-pseudoscalar(P) Lagrangian from the local hidden gauge approach,

$$\mathcal{L}_{\text{VPP}} = -ig \left< \Phi, \partial_\mu \Phi \right| V^\mu \right>, \quad (30)$$

for the VPP vertex, where $g = m_V / 2f_\pi$; $f_\pi = 93$ MeV is the pion decay constant, and $m_V$ is the mass of the light vector mesons. Φ is the pseudoscalar mesons matrix from Eq. (6) and $V^\mu$ is the vector mesons matrix,

$$V^\mu = \begin{pmatrix}
\frac{1}{\sqrt{2}} \rho^0 + \frac{1}{\sqrt{2}} \omega & \rho^+ & K^{*+} & \bar{D}^{*0} \\
\rho^- & -\frac{1}{\sqrt{2}} \rho^0 + \frac{1}{\sqrt{2}} \omega & K^{*0} & \bar{D}^{*-} \\
K^{*-} & \bar{K}^{*0} & \phi & D^{*-} \\
D^{*0} & D^{*+} & D^{*++} & J/\psi
\end{pmatrix}. \quad (31)$$

For the VBB vertex, instead of extending from $SU(3)$ to $SU(4)$ the vector(V)-baryon(B)-baryon(B) Lagrangian from the local hidden gauge approach, we construct spin-flavor wave functions for the baryons, considering the heavy quark as spectator and symmetrizing only in the light quarks. Then we write the vector mesons in terms of their quark content and apply them in the wave functions as number operators times the coupling constant from vector meson exchange (the same of Eq. (30)). As discussed in Ref. [44], this method yields the same results as the VBB Lagrangian of the local hidden gauge in $SU(3)$, being consistent at
the same time with this formalism and with the hypothesis of the heavy quark as spectator, ensuring heavy quark spin symmetry in the dominant terms where only light vectors are exchanged.

With these ingredients and the baryon spin-flavor wave functions, one can obtain the coefficients $D_{ij}$ of the $V$ matrix that we plug in the Bethe-Salpeter equation in the form of Eq. (1), given by

$$V_{ij} = D_{ij} \frac{2 \sqrt{8 - M_{B_i} - M_{B_j}}}{4 f_s^2} \sqrt{\frac{M_{B_i} + E_{B_i}}{2M_{B_i}}} \sqrt{\frac{M_{B_j} + E_{B_j}}{2M_{B_j}}} ,$$

(32)

where $M_{B_i,B_j}$ and $E_{B_i,B_j}$ stand for the mass and the center-of-mass energy of the baryons, respectively, and the matrix $D_{ij}$ is given in Table IV.

TABLE IV. $D_{ij}$ coefficients of Eq. (32) for the $J^P = 1/2^-$ meson-baryon states in $S$-wave.

| $J = 1/2$ | $\Xi_c \bar{K}$ | $\Xi_c' \bar{K}$ | $\Xi D$ | $\Omega_c \eta$ |
|-----------|----------------|----------------|--------|----------------|
| $\Xi_c \bar{K}$ | -1 | 0 | $-\frac{1}{\sqrt{2}} \lambda$ | |
| $\Xi_c' \bar{K}$ | -1 | $\frac{1}{\sqrt{6}} \lambda$ | $-\frac{4}{\sqrt{3}} \lambda$ | |
| $\Xi D$ | -2 | $\frac{\sqrt{2}}{3} \lambda$ | | |
| $\Omega_c \eta$ | | | 0 | |

The details of the calculation can be found in Ref. [44]. In Table IV we have the parameter $\lambda$ in some non diagonal matrix elements, which involve transitions from one meson without charm to one with charm, like $\bar{K} \rightarrow D$. In this case we have to exchange a heavy quark, and the propagator of the exchanged vector goes like

$$\frac{1}{(q^0)^2 - |\mathbf{q}|^2 - m_{D_s^*}^2} \approx \frac{1}{(m_D - m_K)^2 - m_{D_s^*}^2} ,$$

(33)

and the ratio to the propagator of the light vectors is

$$\lambda \equiv \frac{-m_V^2}{(m_D - m_K)^2 - m_{D_s^*}^2} \approx 0.25 .$$

(34)

Then we take $\lambda = 1/4$ in all these matrix elements, as it was done in Ref. [71].

The diagonal matrix elements of Table IV coincide with those of Ref. [43], but not all the non diagonal. SU(4) symmetry is used in Ref. [43], but only SU(3) is effectively used in the diagonal terms. The same happens in our approach, with the difference that the heavy
baryon wave functions we have constructed are not eigenstates of SU(4), since we treat the heavy quark as a spectator and use SU(3) for the light quarks. This induces a spin-flavor dependence different from the one of pure SU(4) symmetry.

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