Compact pentaquark structures

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We study the possibility that at least one of the two pentaquark structures recently reported by LHCb [1–3] could be described as a compact pentaquark state, and we give predictions for new channels that can be studied by the experimentalists if this hypothesis is correct. We use very general arguments dictated by symmetry considerations, in order to describe the pentaquark states within a group theory approach. A complete classification of all possible states and quantum numbers, which can be useful both to the experimentalists in their search for new findings and to theoretical model builders, is given, without the introduction of any particular dynamical model. Some predictions are finally given by means of a Gürsey-Radicati (GR) inspired mass formula. We reproduce the mass and the quantum numbers of the lightest pentaquark state reported by LHCb \((J^P = 3^-)\) with a parameter-free mass formula, fixed on the well-established baryons. We predict other pentaquark resonances (giving their masses, and suggesting possible decay channels) which belong to the same multiplet as the lightest one. Finally, we compute the partial decay widths for all the predicted pentaquark resonances.

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I. INTRODUCTION

The LHCb collaboration has recently reported the observation of exotic structures in \(\Lambda_b\) decay [1], further supported by another two articles of the LHCb collaboration [2, 3]. The decay can proceed according to the diagram in fig. 1, which involves conventional hadrons:

\[ \Lambda_b^0 \rightarrow J/\psi + \Lambda^* \]  

or it can be characterised by exotic contributions, which are referred to as charmonium-pentaquark states (fig. 2):

\[ \Lambda_b^0 \rightarrow P_c^+ + K^- \]  

The LHCb collaboration found two resonant structures:

FIG. 1: Feynman diagrams for \((\Lambda_b^0 \rightarrow J/\psi + \Lambda^*)\) (Fig. taken from Ref. [1]; APS copyright)

FIG. 2: Feynman diagrams for \((\Lambda_b^0 \rightarrow P_c^+ + K^-)\) (Fig. taken from Ref. [1]; APS copyright)

a lower mass state at \(4380 \pm 8 \pm 28\) MeV, with a width of \(205 \pm 18 \pm 86\) MeV, and a higher mass one at \(4449.8 \pm 1.7 \pm 2.5\) MeV, with a width of \(39 \pm 5 \pm 19\) MeV, in the \(J/\Psi p\) invariant mass spectrum. Moreover, according to the LHCb collaboration [1], the preferred \(J^P\) assignments are \(3/2^-\) and \(5/2^+\), respectively.

Since the observation of the two resonant structures, many explanations have been proposed for the LHCb pentaquark states. Meson-baryon molecules were suggested in [4–9]. Pentaquark states of diquark-diquark-antiquark nature were suggested in [11, 12], and \(\bar{D}\) soliton states in [15]. The molecular interpretation works well for the heaviest resonant state (see, for example [4]). Therefore, in this study, we focus on the lightest pentaquark structure (with \(J^P = 3^-\)), by means of a multiquark approach. From the LHCb quantum numbers of the lightest resonant state, we show that it can be described as a pentaquark state with spin \(S = \frac{3}{2}\). We show that the ground state multiplet of the charmonium pentaquark states is a \(SU_f(3)\) octet, and we studied all the charmonium pentaquark states which belong to the octet, predicted their masses, and suggested possible decay channels in which the experimentalists can observe them. By using an effective Lagrangian [26] for the \(P_c J/\Psi\) coupling, in combination with the branching ratio \(B(P_c^+ \rightarrow J/\Psi p)\) upper limit extracted by Wang [27], and with our predicted masses, we compute the partial decay widths for the predicted pentaquark resonances.
II. CLASSIFICATION OF THE \( qqcc \) MULTIPLETS AS BASED ON SYMMETRY PROPERTIES

In order to classify the pentaquark multiplets, we made use, as much as possible, of symmetry principles, without introducing any explicit dynamical models. We made use of the Young tableau technique, adopting for each representation the notation \( [f]_d = [f_1, \ldots, f_n]_d \), where \( f_i \) denotes the number of boxes in the \( i \)-th row of the Young tableau, and \( d \) is the dimension of the representation.

In agreement with the \( LHC_b \) hypothesis [1], we think the charmonium pentaquark wave function as \( qqcc \) where \( q = u, d, s \) is a light quark and \( c \) is the heavy charm quark. Let us first discuss the possible configurations of \( qqq \) quarks in the \( qqcc \) system.

The \( c\bar{c} \) can be in a colour octet or singlet with spin 0 or 1. The colour wave function of the \( qqcc \) system must be an \( SU_3(3) \) singlet, so the remaining three light quarks are also in a colour-singlet or in a colour-octet.

The orbital symmetry of the quark wave function depends on the quantum numbers of the resonant state \( J^P = \frac{3}{2}^+ \). Indeed, the parity \( P \) of the pentaquark system is:

\[
P | qqcc \geqslant (-1)^{l+1}
\]

and so \( l \) must be even. The total angular momentum is \( J = \frac{3}{2} \), and so \( l = 0 \) or \( l = 2 \). In this paper, we hypothesise that the lightest charmonium pentaquark state reported by the \( LHC_b \) collaboration is a ground state pentaquark with \( l = 0 \), and so each quark is in an S-wave.

The three light quarks must satisfy the Pauli principle. As a consequence, since the \( q^2 \) orbital part is completely symmetric, the spin-flavour and the colour part are conjugate: spin-flavour symmetric state if they are in a colour singlet; or spin-flavour mixed symmetry state if they are in a colour octet. Therefore, the allowed \( SU_{sf}(6) \) spin-flavour pentaquark configurations are a 56-plet ([3]_{56}) and a 70-plet ([21]_{70}) which correspond respectively to the three light quarks in a colour singlet and in a colour octet. In Tab. 1 the analysis of the flavour and spin content of the spin-flavour 56-plet and of the 70-plet, i.e. their decomposition into the representations of \( SU_f(3) \otimes SU_s(2) \), is reported.

The \( SU_{sf}(6) \) 70-plet contains an \( SU_f(3) \) flavour octet \([21]_8\) and a decuplet \([3]_{10}\), while the 56-plet contains an \( SU_f(3) \) flavour singlet [111]_1, two octets \([21]_8\) and a decuplet \([3]_{10}\). The allowed \( SU_f(3) \) flavour representations to which the charmonium pentaquark states can belong are therefore:

\[
[111]_1, [21]_8, [3]_{10}
\]

Since the charmonium pentaquark state, as reported by \( LHC_b \), has a quark content \( uudc\bar{c} \), it does not have strange quarks, and so the strangeness \( S = 0 \). In the case of 3 flavours \((u, d, s)\), the hypercharge \( Y \) is defined as:

\[
Y = B + S
\]

where \( B \) is the barionic number and \( S \) is the strangeness. Since the charmonium pentaquark state, as reported by \( LHC_b \), has a quark content \( uudc\bar{c} \), it does not have strange quarks, and so the strangeness \( S = 0 \), the charm \( C \) is 0, the barionic number \( B = 1 \), and then \( Y \) must be equal to 1.

Therefore the pentaquark state must be found in a \( Y = 1 \) submultiplet of the allowed flavour states of Eq. 4. Following this reasoning, we must exclude the singlet [111]_1, because it does not have any \( Y = 1 \) submultiplets; therefore, the remaining possible \( SU_f(3) \) multiplets for the charmonium pentaquark states are:

\[
[21]_8, [3]_{10}
\]

III. THE EXTENSION OF THE GÜRSEY-RADICATI MASS FORMULA

In order to determine the mass splitting between the multiplets of Eq. 6 we made use of a Gürsey-Radicati (GR)-inspired formula [22]. As yet, there is experimental evidence of only two charmonium pentaquark states. This is not sufficient to determine all parameters in the GR mass formula, and then to predict the masses of the other pentaquarks. For this reason, we use the values of the parameters determined from the three-quark spectrum (see Tab. 1), assuming that the coefficients in the GR formula are the same for different quark systems. The simplest GR formula extension which permits us to

\[
\begin{array}{ccc}
SU_{sf}(6) & SU_f(3) & SU_s(2) \\
\hline
[3]_{56} & [3]_{10} & [3]_{4} \\
[21]_8 & [21]_2 \\
[21]_{70} & [3]_{10} & [21]_2 \\
[21]_8 & [3]_{4} \\
[21]_8 & [21]_2 \\
[111]_1 & [21]_2 \\
\end{array}
\]
distinguish the different multiplets of $SU_f(3)$ is

$$M_{GR} = M_0 + AS(S + 1) + DY + +E \left[ I(I + 1) - \frac{1}{4}Y^2 \right] + GC_2(SU(3)) + FN_C$$  (7)

where $M_0$ is a scale parameter: this means that, for example, in baryons each quark gives a contribution of roughly $\frac{1}{3}M_0$ to the whole mass. $I$ and $Y$ are the isospin and hypercharge, respectively, while $C_2(SU(3))$ is the eigenvalue of the $SU_f(3)$ Casimir operator. Finally, $N_C$ is a counter of $c$ quarks or $\bar{c}$ antiquarks. This term takes into account the mass difference between a $c$ quark (or a $\bar{c}$ antiquark) in relation to the light quarks ($u, d$). The coefficients $A, D, E, G, F$ and the scale parameter $M_0$ have been fixed by using the well-established baryons spectrum.

Tab. II reports the baryons used to fix the parameters in Eq. 7, the $SU_f(3)$ multiplet which they were assigned to, the corresponding eigenvalues of the Casimir operator $C_2(SU(3))$, their quantum numbers, and the values of $N_c$.

In Tab. III all the parameters, with their corresponding values, are reported.

| baryon | $SU_f(3)$ | $C_2(SU(3))$ spin | $Y$ | $I$ | $N_c$ |
|--------|----------|------------------|----|----|-------|
| $\Lambda(1116)$ | [21]$_8$ | 3 | $\frac{1}{2}$ | 0 | 0 |
| $\Lambda^+_c(2286)$ | [11]$_3$ | $\frac{3}{2}$ | $\frac{1}{2}$ | 0 | 1 |
| $\Sigma^+_c(2455)$ | [2]$_6$ | $\frac{3}{2}$ | $\frac{1}{2}$ | $\frac{1}{2}$ | 1 |
| $\Xi^+_c(2471)$ | [11]$_3$ | $\frac{3}{2}$ | $\frac{1}{2}$ | $-1$ | $\frac{1}{2}$ | 1 |
| $\Xi^+_c(2576)$ | [2]$_6$ | $\frac{3}{2}$ | $\frac{1}{2}$ | $-1$ | $\frac{1}{2}$ | 1 |
| $\Omega^0_c(2695)$ | [2]$_6$ | $\frac{3}{2}$ | $\frac{1}{2}$ | $-\frac{3}{2}$ | 0 | 1 |
| $\Omega^+_c(2766)$ | [2]$_6$ | $\frac{3}{2}$ | $\frac{1}{2}$ | $-\frac{3}{2}$ | 0 | 1 |

TABLE II: On the left side, the list of baryons used to fix the parameters in Eq. 7 is reported. In the following columns we report the multiplet which they were assigned to, with the corresponding eigenvalues of the Casimir operator $C_2(SU(3))$, their quantum numbers, and the values of $N_c$, respectively.

| $M_0$ | $A$ | $D$ | $E$ | $F$ | $G$ |
|-------|-----|-----|-----|-----|-----|
| value (MeV) | 940.8 | 23.6 | -157.3 | 32.0 | 1365.7 | 52.5 |

TABLE III: values of the parameters in the GR extended mass formula.

In order to show the reliability of the values obtained with the GR mass formula extension, we calculated the predicted mass of the two charmed baryons $\Xi_c(2520)$ and $\Xi_c(2645)$ reported by the PDG [13]. The quantum number assignments and the predicted masses are reported in Tabs. IV and V respectively.

| baryon | $SU_f(3)$ | $C_2(SU(3))$ spin | $Y$ | $I$ | $N_c$ |
|--------|----------|------------------|----|----|-------|
| $\Sigma_C(2520)$ | $\Sigma^+_{C}$ | 2518.41$^{+0.21}_{-0.19}$ | 2526 |
| $\Xi_C(2645)$ | $\Xi^+_{C}$ | 2645.9$^{+0.5}_{-0.5}$ | 2646 |

TABLE IV: On the left side, the $\Sigma_C$ and $\Xi_C$ baryons. On the right side, the $SU_f(3)$ multiplet which they were assigned to, with eigenvalues of the Casimir operator $C_2(SU(3))$, the assigned quantum numbers, and the values of $N_c$.

| baryon | isospin | exp. masses (MeV) | predicted masses (MeV) |
|--------|---------|------------------|-----------------------|
| $\Sigma_C(2520)$ | $\Sigma^+_{C}$ | 2518.41$^{+0.21}_{-0.19}$ | 2526 |
| $\Sigma_C(2520)$ | $\Sigma^0_{C}$ | 2518.48$^{+0.20}_{-0.20}$ | 2526 |
| $\Xi_C(2645)$ | $\Xi^+_{C}$ | 2645.9$^{+0.5}_{-0.5}$ | 2646 |
| $\Xi_C(2645)$ | $\Xi^0_{C}$ | 2645.9$^{+0.5}_{-0.5}$ | 2646 |

TABLE V: The first column shows $\Sigma_C$ and $\Xi_C$ baryons; the possible isospin configurations are reported in the second column, while in the third column the corresponding experimental masses as from PDG [13]. In the last column the predicted masses, calculated by means of the mass formula of Eq. 7 using the coefficients of Tab. III, are reported.

IV. APPLICATION OF THE GR FORMULA TO THE PENTAQUARK STATES

In Eq. 6 we reported the possible $SU_f(3)$ multiplets for the charmonium pentaquark states. We hypothesise that the charmonium pentaquark state $J^P = \frac{3}{2}^-$, reported by the LHCb collaboration, belongs to the lowest mass $SU_f(3)$ multiplet. According to the GR formula 7, the mass splitting between the different $SU_f(3)$ multiplets of Eq. 6 is due to the different eigenvalues of the Casimir operator $C_2(SU(3))$, and so it is proportional to the coefficient $G$ (reported in Tab. III). Since $G$ is positive ($G = 52.5$ MeV), the lowest mass multiplet is the one with the minimum Casimir operator eigenvalue, and so it is the octet (see Tab. VI). In Tab. VI, each multiplet, with the corresponding eigenvalues of the Casimir operator $C_2(SU(3))$, is reported.
The result is in agreement with the mass reported by the LHCb collaboration. Despite the simplicity of the approach that we used, this state is labelled with \(P_{21}\), and its theoretical mass, predicted by means of our GR formula extension, is \(M = 4404\) MeV. Its mass is in agreement with the mass reported by the LHCb collaboration: \(M = 4380 \pm 8 \pm 29\) MeV.

### V. Decay Channels

We will now explore the possible decay channels in which the other predicted states of the octet can be observed. These channels will be described in detail. The state \(P^{0+}(4404)\) is a part of an isospin doublet. In order to observe its isospin partner \((P^{00}(4404))\), a possible decay channel could be:

\[
\Lambda_b^0 \rightarrow P^{00} + \bar{K}^{00}, P^{00} \rightarrow J/\Psi + n.
\]

The corresponding Feynman diagram is reported in Fig. 4.

![FIG. 4: \(\Lambda_b\) baryon decay in \(P^{00}(4404)\) and \(\bar{K}^{00}\), where \(P^{00}(4404)\) is the neutral pentaquark state, a member of the isospin doublet with \(Y = 1\).](image)

With respect to the other charmonium pentaquark states of the octet, with strangeness, we have to focus on the decays of bottom baryons with strange quarks. Let us consider the following \(\Xi_b^-\) decay:

\[
\Xi_b^- \rightarrow J/\psi + \Xi^-.
\]

This decay is present in nature and was discovered by the D0 collaboration ([19]). In analogy with the exotic \(\Lambda_b^0\) decay of Fig. 2, we can expect that also in the case of \(\Xi_b^-\) baryon there is another possible exotic decay channel:

\[
\Xi_b^- \rightarrow P^{10}/P^{10} + K^-, P^{10} / P^{10} \rightarrow J/\Psi + \Sigma/\Lambda,\]

where \(P^{10}(4609)\) and \(P^{10}(4535)\) have the same quark content (\(uudc\)), and belong to the isospin triplet, and to the isosinglet, respectively (see Fig. 3). Since they have the same quark content and both are neutral, they can both come from the \(\Xi_b^-\) decay. The charmonium pentaquark state \(P^{1-}(4609)\) can be observed in the following decay process:

\[
\Xi_b^- \rightarrow P^{1-} + \bar{K}^0, P^{1-} \rightarrow J/\Psi + \Sigma^-.
\]
The difference between the two suggested decays for the \( \Xi_b \) baryon (Eq. 10 and Eq. 11) is in the final state: in the case of the final state of Eq. 10, a couple of quarks \( \bar{u}u \) comes from the vacuum, while, in the decay of Eq. 11, the couple of quarks \( \bar{u}u \) is replaced with a couple \( dd \). The baryon \( \Xi_b \) is a member of an isodoublet. The decay of its isospin partner \( \Xi_b^0 \):

\[
\Xi_b^0 \rightarrow P^{1+} + K^-, P^{1+} \rightarrow J/\Psi + \Sigma^+ \quad (12)
\]
is probably the most important one from the experimental point of view, since all the final state particles are charged and, therefore, easier to detect.

In order to have a final pentaquark state with two strange quarks \( s \), we need a double strange baryon in the initial state. The known decay channel of the \( \Omega_b \) baryon is:

\[
\Omega_b^- \rightarrow J/\psi + \Omega^- \quad (13)
\]

This decay was discovered by the D0 detector at the Fermilab Tevatron collider [20]. Another possible \( \Omega_b^- \) decay channel may be, in analogy with the exotic \( \Lambda_b \) decay channel of Fig. 2:

\[
\Omega_b^- \rightarrow P^{20} + K^-, P^{20} \rightarrow J/\Psi + \Xi^0 \quad (14)
\]
The state \( P^{20}(4719) \) of Eq. 14 is a part of an isospin doublet (see Fig. 3). In order to observe its isospin partner \( (P^{2-}(4719)) \), a possible decay channel could be:

\[
\Omega_b^- \rightarrow P^{2-} + K^0, P^{2-} \rightarrow J/\Psi + \Xi^- \quad (15)
\]
The difference between the \( \Omega_b^- \) decays of Eq. 14 and that of Eq. 13 is, respectively, the creation of a couple \( \bar{u}u \) and \( dd \) from the vacuum.

### VI. PARTIAL DECAY WIDTHS

We adopt the effective Lagrangian for the \( P_c N J/\Psi \) couplings from Ref. [26] as follows:

\[
L_{P_c N J/\Psi}^{3/2-} = i\overline{P}_c\mu \left[ \frac{g_1}{2M_N} \Gamma_\nu N \right] \psi^{\mu\nu} + \text{H.c.} \quad (16)
\]

\[
- i\overline{P}_c\mu \left[ \frac{ig_2}{(2M_N)^2} \Gamma^\nu \partial_\nu N + \frac{ig_3}{(2M_N)^2} \Gamma^\nu N \partial_\nu \right] \psi^{\mu\nu} + \text{H.c.}
\]

where \( P_c \) is the pentaquark field with spin-parity \( J^P = \frac{3}{2}^- \), \( N \) and \( \psi \) are the nucleon and the \( J/\Psi \) fields, respectively. The \( \Gamma \) matrices are defined as follows:

\[
\Gamma_\nu = \begin{pmatrix} \gamma_\nu \gamma_5 \\ \gamma_\nu \end{pmatrix}, \quad \Gamma^- = \begin{pmatrix} \gamma_\nu \\ 0 \end{pmatrix} \quad (17)
\]

As noticed by Wang [27] in the pentaquark state decays into \( J/\psi \), the momentum of the final states are fairly small compared with the nucleon mass. Thus, the higher partial wave terms proportional to \((p/M_N)^2\) and \((p/M_N)^3\) can be neglected, so we only consider the first term in Eq. (10). This approximation leads to the following expression for the \( P^0_c(4380) \) partial decay width in the \( N J/\psi \) channel [26]:

\[
\Gamma(P_c^0 \rightarrow N J/\psi) = \frac{\bar{g}^2 N J/\psi}{12\pi} \frac{p_N}{M_{P_c^0}} (E_N + M_N)
\]

\[\times [2E_N(E_N - M_N) + (M_{P_c^0} - M_N)^2 + 2M_{J/\psi}^2] \quad (18)\]

with

\[
\bar{g}_{N J/\psi} = \frac{g_1}{2M_N} \quad (19)
\]

The kinematic variables \( E_N \) and \( p_N \) in Eq. (18) are defined as \( E_N = (M_{P_c^0} + M_N^2 - M_{J/\psi}^2)/(2M_{P_c^0}) \) and \( p_N = \sqrt{E_N^2 - M_N^2} \).

Unfortunately, the branching ratio \( B(P_c^0 \rightarrow J/\Psi \bar{p}) \) is not known at present, so the coupling constant \( g_1 \) of Eq. (19) is unknown. However, by using our pentaquark mass predictions, we can provide an expression of the partial decay widths for the pentaquark states with open strangeness. For example, the \( P_{c^+} \) partial decay width in the \( \Sigma^+ J/\Psi \) channel is given by:

\[
\Gamma(P_{c^+} \rightarrow \Sigma^+ J/\Psi) = \frac{\bar{g}^2 \Sigma^+ J/\Psi}{12\pi} \frac{p_{\Sigma^+}}{M_{P_{c^+}}} (E_{\Sigma^+} + M_{\Sigma^+})
\]

\[\times [2E_{\Sigma^+}(E_{\Sigma^+} - M_{\Sigma^+}) + (M_{P_{c^+}} - M_{\Sigma^+})^2 + 2M_{J/\psi}^2] \quad (20)\]

and the coupling constant \( \bar{g}_{\Sigma^+ J/\Psi} \) is:

\[
\bar{g}_{\Sigma^+ J/\Psi} = \frac{g_1}{2M_{\Sigma^+}} \quad (21)
\]

The expressions for the partial decay widths of the \( \Lambda J/\Psi \), \( \Sigma J/\Psi \), and \( \Xi J/\Psi \) channels are listed in Table VII.

| initial state | channel partial width (MeV) |
|--------------|-----------------------------|
| \( P_{c^0} \) | \( \Lambda J/\Psi \) \( (0.81\Gamma_{N J/\Psi}) \) |
| \( P_{c^0}^{-}, P_{c^0}^{10}, P_{c^+}^{10}, \ldots \) | \( \Sigma J/\Psi \) \( (0.73\Gamma_{N J/\Psi}) \) |
| \( P_{c^0}^{-}, P_{c^0}^{20}, \ldots \) | \( \Xi J/\Psi \) \( (0.65\Gamma_{N J/\Psi}) \) |

TABLE VII: Partial decay widths expressions for \( \Lambda J/\Psi \), \( \Sigma J/\Psi \), and \( \Xi J/\Psi \) channels.

Since the pentaquark states were observed in \( J/\Psi p \) channel, it is natural to expect that they can be produced in \( J/\Psi p \) photoproduction via the s and u-channel process. Wang et al. [27] calculated the pentaquark states cross section in \( J/\Psi \) photoproduction and compared it with the present experimental data [29] [30] [31]. The coupling between \( J/\Psi p \) and the two pentaquark states are extracted by assuming it accounts for their total width and 5%, respectively. As a result, they found that if one assumes that the \( J/\Psi p \) channel saturates the total width
of the two pentaquark states (that is $\mathcal{B}(P_c^+ \to J/\Psi p) = 1$ ) one significantly overestimates the experimental data. In conclusion they found that to be consistent also with the present photoproduction data, it is necessary that the branching ratio for both the pentaquark states is $\mathcal{B}(P_c^+ \to J/\Psi p) \leq 0.05$. Thus, if we use the upper branching ratio limit extracted by Wang [27], that is $\mathcal{B}(P_c^+ \to J/\Psi p) = 0.05$, we obtain that the $P_c(4380)$ partial decay width for the $J/\Psi p$ channel is

$$\Gamma_{NJ/\Psi} = \mathcal{B}(P_c^+ \to J/\Psi p)\Gamma_{\text{tot}} = 10.25\text{ MeV} \quad (22)$$

where $\Gamma_{\text{tot}}$ as reported by the LHCb collaboration, is 205 MeV. The numerical results for the other channels are listed in Table VII.

| initial state | channel partial width (MeV) |
|--------------|-----------------------------|
| $P_c^{1+}$, $P_c^{10}$, $P_c^{1-}$, $\Sigma J/\Psi$ | 7.59 |
| $P_c^{2+}$, $P_c^{20}$, $\Xi J/\Psi$ | 6.69 |

TABLE VIII: Partial decay widths for $\Lambda J/\Psi$, $\Sigma J/\Psi$ and $\Xi J/\Psi$ channels. The partial decay widths are calculated from the constraint that $\Lambda J/\Psi$ channel accounts for the 5% of the total pentaquark width, as calculated by Wang in (27).

VII. CONCLUSIONS

The LHCb collaboration has recently reported the observation of two exotic structures in $J/\Psi p$ channel [1], which they referred to as charmonium pentaquark states ( with a quark content $uudc\bar{c}$ ) further supported by another two articles by the LHCb collaboration [2] [3]. The significance of each of these states is more than 9 standard deviations. The lightest one has a mass of $4380 \pm 8 \pm 29$ MeV and a width of $205 \pm 18 \pm 86$ MeV, while the heaviest has a mass of $4449.8 \pm 1.7 \pm 2.5$ MeV and a width of $39 \pm 5 \pm 19$ MeV. The preferred $J^P$ assignments, according to the LHCb collaboration [1], are $3/2^-$ and $5/2^+$, respectively.

The earliest prediction for the charmonium pentaquark with $J^P = \frac{3}{2}^-$ was given by J. J. Wu et al. [13]. The heaviest pentaquark state has been apparently well explained by means of a molecular approach [1] [6], and it was also predicted in a molecular approach before the LHCb discovery by [14], in a coupled-channel unitary approach.

As regards the lightest one, molecular models have also been proposed, but the predictions are not so good as for the heaviest state [1] [6]. Some predictions of its mass and quantum numbers were given by [21] in 2012, by means of a potential quark model approach, but these predictions depend strongly on the particular interaction used: colour-magnetic interaction (CM) based on one-gluon exchange, chiral interaction (FS ) based on meson exchange, and instanton-induced interaction (Inst.) based on the non-perturbative QCD vacuum structure.

In this present study, we focused on describing the lightest resonant state ($J^P = \frac{3}{2}^-$), by means of a multiquark approach. An extension of the original GR mass formula [22] which correctly describes the charmed baryon sector was performed (Tab. V), and also proved able to give an unexpected prediction for the mass of the lightest pentaquark state $\frac{3}{2}^-$, which is in agreement with the experimental value within one standard deviation.

We found that the lightest pentaquark state $\frac{3}{2}^-$ belonged to the $SU_f(3)$ octet [21]a. The theoretical mass of the lightest pentaquark state $\frac{3}{2}^-$ predicted by means of the GR formula extension (Eq. 7) is $M = 4404$ MeV, in agreement with the experimental mass $M = 4380 \pm 8 \pm 29$ MeV. We also predicted other pentaquark states, which belong to the same $SU_f(3)$ multiplet as the lightest resonance $J^P = \frac{3}{2}^-$, giving their mass, and suggesting possible decay channels in which they can be observed. We have finally computed the partial decay widths for all the suggested octet-pentaquark decay channels.

As the $\Lambda_b \to J/\Psi K^- p$ decay is expected to be dominated by $\Lambda^* \to K^- p$ resonances [1], we observe that the poor knowledge about the $\Lambda^*$ excited states can affect the estimation of the parameters of the two pentaquark resonances. Moreover, as was noticed by Wang [27], if the two pentaquark candidates are genuine states, their production in photoproduction should be a natural expectation. For these reasons, on the one hand it is important to increase our knowledge about the missing excited states $\Lambda^*$ with new experiments [32], in order to improve the analysis and to extract with more precision the two pentaquark masses and widths. On the other hand, a refined measurement of the $J/\Psi$ photoproduction cross section would provide more information about the nature of the pentaquark states.
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