Traces of the hidden-charm S=-1 pentaquark in the $\Lambda_b \to J/\Psi\eta\Lambda$ decay

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Abstract. The hidden charm pentaquark state $P_c(4450)$, observed recently by the LHCb collaboration in the $\Lambda_b \to J/\psi K^- p$ decay, may be of molecular nature, as advocated by some unitary approaches that also predict pentaquark partners in the strangeness S=-1 sector. In this work we argue that a hidden-charm strange pentaquark could also be seen in the decay of the Λ_b, but through the $J/\psi\eta\Lambda$ decay mode, by studying the invariant mass spectrum of $J/\psi\Lambda$ pairs.
In our model we assume a standard weak decay topology, then incorporate the hadronization process and final state interaction effects, and we find that the $J/\psi\eta\Lambda$ final state is populated with the strength similar to that of the $J/\psi K^- p$. We have studied the dependence of our results on reasonable changes in the parameters of the models as well as on the unknown properties of the speculated hidden charm strange pentaquark. We have observed that, while there appear changes in the position of the peak and in the shapes of the distributions, a resonance signal in the $J/\psi\Lambda$ invariant mass spectrum is clearly seen in all the cases. This gives us confidence that such an experimental study could result into a successful proof of the existence of this new state.

1 Introduction

The LHCb collaboration reported recently two exotic structures in the invariant $J/\psi p$ mass spectrum of the $\Lambda_b^0 \to J/\psi K^- p$ process. These pentaquark states were named $P_c(4380)$, with a mass of $4380 \pm 8 \pm 29$ MeV and a width of $205 \pm 18 \pm 86$ MeV, and $P_c(4450)$, with a mass of $4449.8 \pm 1.7 \pm 2.5$ MeV and a width of $39 \pm 5 \pm 19$ MeV [1, 2]. Hidden charm baryon states with similar characteristics of the states reported had already been predicted, employing a molecular picture [3–7] or a quark model approach [8, 9]. A list of early references on pentaquark states can be seen in Ref. [10].

We would like to remind the reader that a theoretical study of the $\Lambda_b^0 \to J/\psi K^- p$ reaction was done in [11], prior to the experimental study of Ref. [1], predicting the contribution of the tail of the $\Lambda(1405)$ in the $K^- p$ invariant mass distribution. The analysis of [1] contains such a contribution in agreement in shape with the predictions, where the absolute normalization was unknown. Moreover, it was shown in [12] that the distributions in the pentaquark channel, i.e. in the invariant $J/\psi p$ mass

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spectrum of [1], could be explained via the incorporation of the hidden charm N^* states predicted in [3–6], which are molecular states mostly made from $\bar{D}^*\Sigma_c$ or $\bar{D}^*\Sigma_c^*$ components and having a small coupling to $J/\psi p$, one of their open decay channels. We also recall that when the hidden charm N^* resonances were theoretically predicted as molecular states in several unitary approaches, some partner hidden charm strange Λ^* states were also predicted.

It is unlikely that there are no partners of the states found in [1], and indeed, in [3, 4] states of spin-parity $3/2^-$ with hidden charm but strangeness $S = -1$ were found, mostly made of $\bar{D}^*\Xi_c$ or $\bar{D}^*\Xi_c'$, decaying into $J/\psi\Lambda$. In view of this, the decay of Ξ_b^- into $J/\psi K^-\Lambda$ was suggested in [13] as a suitable reaction to find a hidden charm strange state. Predictions for the $\bar{K}\Lambda$ and $J/\psi\Lambda$ mass distributions were done, and, playing with uncertainties, it was shown that a clear peak in the $J/\psi\Lambda$ mass distribution should show up. This reaction is presently being considered by the LHCb collaboration. However, since there is a much smaller statistics in the production of Ξ_b^- than that of Λ_b, it is interesting to explore alternative reactions to observe this strangeness $S = -1$ hidden-charm pentaquark. In [14] the authors suggest to employ the $J/\psi\eta\Lambda$ decay mode of the Λ_b. Since the $\eta\Lambda$ pair is populated with a weight $\sqrt{2}/3$ relative to the K^-p pair in the primary $\Lambda_b^0 \to J/\psi MB$ reaction [11], the $\Lambda_b \to J/\psi\,\eta\Lambda$ decay rate should be similar as that found for $J/\psi K^-p$ final states in the study of the non-strange pentaquark, and the new strange state should be looked for in the $J/\psi\Lambda$ mass distribution instead of the $J/\psi p$ one. The possible existence of an strange $S = -1$ pentaquark partner was also studied in [15] from the non-strange decay mode $\Lambda_b \to J/\psi K^0\Lambda$.

In these proceedings we will concentrate on the case of a $J^P = 1/2^-$ strange pentaquark, although in the original paper [14] (as well as in Ref. [15]) the possibilities to produce both $J^P = 1/2^-$ and $J^P = 3/2^-$ pentaquarks have beed explored.

2 The $\Lambda_b \to J/\psi\,\eta\Lambda$ decay process

The study of the $\Lambda_b \to J/\psi\,\eta\Lambda$ decay follows the same approach as that presented in [11, 16] for $\Lambda_b \to J/\psi\; K^-\; p$ [14, 17]. At quark level, both processes are identical and proceed through the transition diagram depicted in figure 1, where we can see the W-exchange weak process transforming the b quark into $c\bar{c}s$, followed by the hadronization of a pair of quarks which eventually produces a meson and a baryon, in addition to the J/ψ. The process depicted in figure 1 assumes that the elementary weak decay involves only the b quark of the Λ_b, while the u and d quarks remain as spectators, the reason being that one expects one-body operators in a microscopical evaluation to have larger strength than two- or multi-body operators. According to this assumption, since the Λ_b has isospin $I = 0$, so does the spectator ud pair, which, combined with the s quark after the weak decay, can only form $I = 0$ Λ states. The findings of the experimental analysis of Ref. [1] clearly support this hypothesis.

For the hadronization process we introduce a $\bar{q}q$ pair between two quarks with the quantum numbers of the vacuum, $\bar{u}u + \bar{d}d + \bar{s}s$. The dominant contribution of the hadronization preserves the spectator role of the ud pair, which ends up into the final baryon, and requires the involvement of the s quark, which ends up into the final meson. Any other topology that would bring the u or d quark into the final meson requires a large momentum transfer that supresses the mechanism. If, in addition, we wish to have the meson-baryon pair in s-wave, it will have negative parity, forcing the s quark prior to hadronization to have also this parity and thus be in an excited state. Since in the final K^- or η mesons the s quark is in its ground state, this also implies that the s quark produced immediately after the weak process must participate actively in the process of hadronization, which proceeds as shown in figure 1. A further discussion on different hadronization mechanisms can be found in [18].

The technical way to implement the hadronization and produce meson-baryon pairs in the final state follows the same steps as in [19–21] for meson decays and in [11, 22] for the Λ_b decay. The

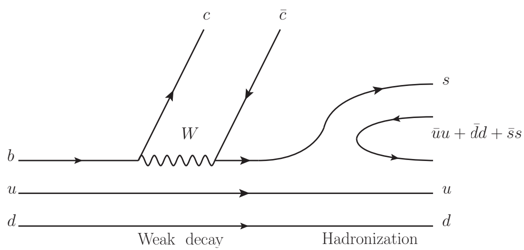

Figure 1: Diagram describing the weak decay of the Λ_b into the J/ψ and a meson-baryon pair formed through a hadronization mechanism.

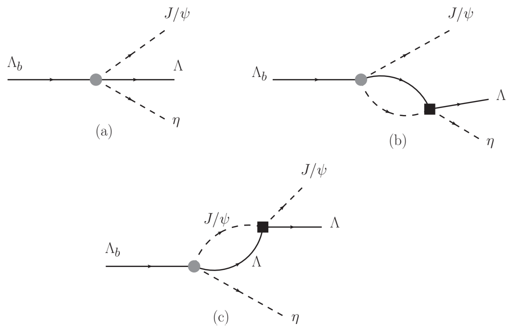

Figure 2: Diagrammatic representation of the decay amplitude for $\Lambda_b \to J/\psi\,\eta\Lambda$: a) tree level, b) the $\eta\Lambda$ production through the coupled channel interaction of the initially produced $\eta\Lambda$ and $\bar{K}N$ meson-baryon pairs, c) $J/\psi\Lambda \to J/\psi\Lambda$ interaction.

flavor decomposition of the Λ_b state is:

$$|\Lambda_b\rangle = \frac{1}{\sqrt{2}}|b(ud-du)\rangle\,, \tag{1}$$

which becomes, after the weak process

$$|H\rangle = \frac{1}{\sqrt{2}}|s(ud-du)\rangle\,, \tag{2}$$

or, upon hadronization,

$$|H\rangle = \frac{1}{\sqrt{2}}|s(\bar{u}u+\bar{d}d+\bar{s}s)(ud-du)\rangle\,, \tag{3}$$

which can be written in terms of the $q\bar{q}$ matrix P, as

$$|H\rangle = \frac{1}{\sqrt{2}}\sum_{i=1}^{3}|P_{3i}q_i(ud-du)\rangle\,, \quad \text{where} \quad P = \begin{pmatrix} u\bar{u} & u\bar{d} & u\bar{s} \\ d\bar{u} & d\bar{d} & d\bar{s} \\ s\bar{u} & s\bar{d} & s\bar{s} \end{pmatrix} \text{ and } q = \begin{pmatrix} u \\ d \\ s \end{pmatrix}. \tag{4}$$

Writing the matrix P in terms of the meson states, $P \to \phi$, where the η,η' mixing [23] has been assumed,

$$\phi = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{3}} + \frac{\eta'}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{3}} + \frac{\eta'}{\sqrt{6}} & K^0 \\ K^- & \bar{K}^0 & -\frac{\eta}{\sqrt{3}} + \frac{2\eta'}{\sqrt{6}} \end{pmatrix}, \tag{5}$$

the hadronized state becomes:

$$|H\rangle = \frac{1}{\sqrt{2}}\left(K^-u(ud-du) + \bar{K}^0d(ud-du) + \frac{1}{\sqrt{3}}\left(-\eta+\sqrt{2}\eta'\right)s(ud-du)\right) \tag{6}$$

By the former equation one obtains the mixed antisymmetric representation of the octet of baryons and taking the results of [24] (see also [11]) one finds the final representation

$$|H\rangle = |K^-p\rangle + |\bar{K}^0n\rangle + \frac{\sqrt{2}}{3}|\eta\Lambda\rangle \tag{7}$$

where we have omitted the $|\eta'\Lambda\rangle$ contribution because of the large mass of the η' meson [11]. Note that, in order to accomodate the wave-function of Ref. [24] to the phase convention of the chiral Lagrangian, a change of sign on the Σ^+, Ξ^0 and Λ states needs to be done. This prescription changes the sign of the $\eta\Lambda$ coefficient in eq. (7) with respect to that employed in [14]. This affects the numerical results, but qualitative features and the conclusions of Ref. [14] remain the same.

The final step consists in taking into account the final state interaction of the meson-baryon pairs. The amplitude for the $\Lambda_b \to J/\psi\,\eta\Lambda$ decay will then be built from the diagrams of figure 2, where we can see the direct tree-level process, depicted by diagram (a), the final-state interaction contribution of the meson-baryon pair into $\eta\Lambda$ production (b), and the final-state $J/\psi\Lambda \to J/\psi\Lambda$ interaction (c). The corresponding amplitude can be written as:

$$\mathcal{M}(M_{\eta\Lambda}, M_{J/\psi\Lambda}) = V_p\Big[h_{\eta\Lambda} + \sum_i h_i G_i(M_{\eta\Lambda}) t_{i,\eta\Lambda}(M_{\eta\Lambda}) + h_{\eta\Lambda} G_{J/\psi\Lambda}(M_{J/\psi\Lambda})\, t_{J/\psi\Lambda,J/\psi\Lambda}(M_{J/\psi\Lambda})\Big], \quad (8)$$

where the weights h_i, obtained from Eq. (7), are:

$$h_{\pi^0\Sigma^0} = h_{\pi^+\Sigma^-} = h_{\pi^-\Sigma^+} = 0\,,\; h_{\eta\Lambda} = \frac{\sqrt{2}}{3}\,,\; h_{K^-p} = h_{\bar{K}^0 n} = 1\,,\; h_{K^+\Xi^-} = h_{K^0\Xi^0} = 0\,, \quad (9)$$

and where G_i, with $i = K^-p, \bar{K}^0 n, \eta\Lambda$, denotes the meson-baryon loop function, chosen in accordance with the model for the scattering matrix $t_{i,\eta\Lambda}$ [25]. Similarly, we take the loop function $G_{J/\psi\Lambda}$ employed in the model of [3, 4] on which, as we will show below, we base our prescription for $t_{J/\psi\Lambda,J/\psi\Lambda}$. The factor V_p, which includes the common dynamics of the production of the different pairs, is unknown and we take it as constant, see Ref. [22] for a more detailed argumentation.

We note that, since our model will rely on the strange pentaquark predicted in Refs. [3, 4] at an energy around 4550 MeV, which couples strongly to $\bar{D}^{*0}\Xi_c'$ states, one should consider the possibility that the influence of this resonance in the final $J/\psi\Lambda$ mass distribution could also be due to the creation of a virtual $\bar{D}^*\eta\Xi_c'$ state in a first step of the Λ_b decay, through the mechanism of figure 3, followed by multiple interactions to generate the resonance, which would eventually decay into a $J/\psi\Lambda$ pair in the final state, represented by the diagrams of figure 4. However, this configuration requires a different topology, as seen in figure 3, in which the *ud* quarks of the Λ_b do not act as a coupled spectator pair. Although it is hard to quantify the size of the amplitude of figure 1 with respect to that of figure 3, the fact that in this later case one of the spectator quarks ends up in the charmed meson and the other one goes to the baryon makes us believe that the corresponding amplitude will be reduced. We will therefore assume the dominance of the mechanism of figure 1 over that of figure 3 by a factor of two or more and will give predictions for different relative signs of the two processes. We note that the lowest order contribution to the $\Lambda_b \to J/\psi\,\eta\Lambda$ decay induced by virtual $\bar{D}^*\eta\Xi_c'$ states is the amplitude of figure 4 (a). We have checked, by explicit numerical evaluation, that the next-order contribution of figure 4 (b), involving the additional final state interaction of the $\eta\Lambda$ pair, gives a negligible correction, hence it will be ignored in the results presented in Sect. 3.

Adding the process initiated by an intermediate $\bar{D}^*\eta\Xi_c'$ state followed by final state interactions leading to a $J/\psi\Lambda$ pair and an η meson represented by the diagram of figure 4 (a), the final amplitude for $\Lambda_b \to J/\psi\,\eta\Lambda$ decay, producing a strange pentaquark with $J^P = 1/2^-$ becomes:

$$\begin{aligned}\mathcal{M}(M_{\eta\Lambda}, M_{J/\psi\Lambda}) = V_p \Bigg[& h_{\eta\Lambda} + \sum_i h_i G_i(M_{\eta\Lambda}) t_{i,\eta\Lambda}(M_{\eta\Lambda}) \\ & + h_{\eta\Lambda} G_{J/\psi\Lambda}(M_{J/\psi\Lambda})\, t_{J/\psi\Lambda,J/\psi\Lambda}(M_{J/\psi\Lambda}) + \beta\, G_{\bar{D}^*\Xi_c'}(M_{\bar{D}^*\Xi_c'})\, t_{\bar{D}^*\Xi_c',J/\psi\Lambda}(M_{J/\psi\Lambda})\Bigg]\,, \end{aligned} \quad (10)$$

where β is a dimensionless parameter which controls the strength of Λ_b decaying virtually into $\bar{D}^*\eta\Xi_c'$, relative to its decay into $J/\psi\,\eta\Lambda$.

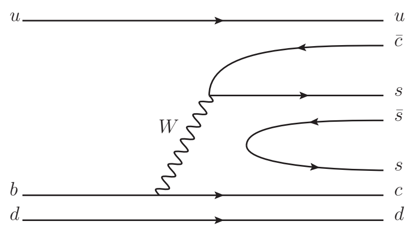

Figure 3: Diagram describing the weak decay Λ_b into the $\bar{D}^*$ and a $\eta\Xi'_c$ pair.

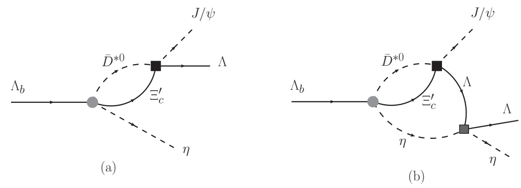

Figure 4: Diagrammatic representation of the decay of the Λ_b into a virtual $\eta\bar{D}^{*0}\Xi'_c$ intermediate state, followed by the $\bar{D}^{*0}\Xi'_c \to J/\psi\Lambda$ conversion process, (a), and implementing also the final state interaction of the final $\eta\Lambda$ pair, (b).

Finally, the double differential cross-section for the $\Lambda_b \to J/\psi\, \eta\Lambda$ decay process reads [11]:

$$\frac{d^2\Gamma}{dM_{\eta\Lambda}dM_{J/\psi\Lambda}} = \frac{3}{16\pi^3}\frac{M_\Lambda M_{\eta\Lambda}M_{J/\psi\Lambda}}{M^2_{\Lambda_b}}|\mathcal{M}(M_{\eta\Lambda}, M_{J/\psi\Lambda})|^2 \,, \tag{11}$$

with $\mathcal{M}$ being that of Eqs. (8) or (10), corresponding to an s-wave weak vertex and, hence, producing a pentaquark with $J = 1/2$. Then fixing, for example, the invariant mass $M_{J/\psi\Lambda}$, one can integrate over $M_{\eta\Lambda}$ in order to obtain $d\Gamma/dM_{J/\psi\Lambda}$.

Now we shall briefly describe the theoretical models employed to obtain the amplitudes $t_{i,\eta\Lambda}$, $t_{J/\psi\Lambda,J/\psi\Lambda}$ and $t_{\bar{D}^*\Xi'_c,J/\psi\Lambda}$, which account for the final state interaction effects.

The $S = -1$ meson-baryon amplitude with $\eta\Lambda$ in the final state, see figure 2 (b) , is determined from the coupled-channel unitary model of Ref. [25–27], developed with the aim of improving upon the knowledge of the chiral interation at next-to-leading order (NLO). The parameters of the model were fitted to a large set of experimental scattering data in different two-body channels, threshold branching ratios, and the precise SIDDHARTA value of the energy shift and width of kaonic hidrogen into consideration, see [25] for details. Differently to other works also involving chiral Lagrangian NLO terms, as e.g. [28–30], the model was also constrained to reproduce the $K^-p \to K^+\Xi^-, K^0\Xi^0$ reactions, since they are especially sensitive to the NLO parameters [25–27].

More especifically, the meson-baryon amplitudes of Ref. [25] are built from a kernel obtained from the SU(3) chiral Lagrangian up to NLO:

$$v_{ij} = v^{WT}_{ij} + v^{NLO}_{ij} + v^{Born,D}_{ij} + V^{Born,C}_{ij} \tag{12}$$

where

$$v^{WT}_{ij} = -\frac{C_{ij}(2\sqrt{s} - M_i - M_j)}{4f^2}N_iN_j\,, \quad v^{NLO}_{ij} = \frac{D_{ij} - 2(k_{i,\mu}k^\mu_j)L_{ij}}{f^2}N_iN_j \quad \text{and} \quad N_i = \sqrt{\frac{M_i + E_i}{2M_i}}\,. \tag{13}$$

The indices i, j stand for any of the ten meson-baryon channels in the neutral $S = -1$ sector, namely K^-p, $\bar{K}^0n$, $\pi^0\Lambda$, $\pi^0\Sigma^0$, $\pi^-\Sigma^+$, $\pi^+\Sigma^-$, $\eta\Lambda$, $\eta\Sigma^0$, $K^+\Xi^-$ and $K^0\Xi^0$, while M_i and E_i are the mass and energy, respectively, of the baryon of the channel "i", and $k_{i,\mu}$ is the corresponding meson four-momenta. The Weinberg-Tomozawa (WT) term is written in terms of SU(3) coefficients C_{ij}, the pion decay constant f_π; while the other 7 low energy constants embedded in the matrices D_{ij} and L_{ij} of the NLO term, which can be found, for example, in the appendices of Ref. [25]. The direct $v^{Bor,D}_{ij}$ and crossed $v^{Born,C}_{ij}$

Born terms are given by

$$v_{ij}^{Born,D} = -\sum_{k=1}^{8} \frac{C^{(\mathrm{Born})}_{\bar{i}i,k} C^{(\mathrm{Born})}_{\bar{j}j,k}}{12f^2} \mathcal{N}_i \mathcal{N}_j \frac{(\sqrt{s}-M_i)(\sqrt{s}-M_k)(\sqrt{s}-M_j)}{s-M_k^2} \ , \tag{14}$$

$$\begin{aligned} v_{ij}^{Born,C} = \sum_{k=1}^{8} \frac{C^{(\mathrm{Born})}_{\bar{j}k,i} C^{(\mathrm{Born})}_{\bar{i}k,j}}{12f^2} \mathcal{N}_i \mathcal{N}_j \times \Bigg[\sqrt{s} + M_k - \frac{(M_i+M_k)(M_j+M_k)}{2(M_i+E_i)\left(M_j+E_j\right)}(\sqrt{s}-M_k+M_i+M_j) \\ + \frac{(M_i+M_k)(M_j+M_k)}{4q_iq_j} \left\{ \sqrt{s} + M_k - M_i - M_j - \frac{s+M_k^2-m_i^2-m_j^2-2E_iE_j}{2(M_i+E_i)\left(M_j+E_j\right)}(\sqrt{s}-M_k+M_i+M_j) \right\} \\ \times \ln \frac{s+M_k^2-m_i^2-m_j^2-2E_iE_j-2q_iq_j}{s+M_k^2-m_i^2-m_j^2-2E_iE_j+2q_iq_j} \Bigg] \ , \end{aligned} \tag{15}$$

where the k label refers to the intermediate baryon involved in the process. The coefficients $C^{(\mathrm{Born})}_{\bar{i}i,k}$ can be found in Ref. [31] and they depend on the axial vector constants D and F; q_i is the center-of-mass (CM) three-momenta in the corresponding channel, and m_i denote the corresponding meson mass.

To obtain chiral unitary amplitudes we solve the on-shell factorized Bethe-Salpeter equation:

$$t_{ij} = v_{ij} + v_{il} G_l t_{lj} \ , \tag{16}$$

where the meson-baryon loop function G_l is obtained employing dimensional regularization

$$G_l = \frac{2M_l}{(4\pi)^2} \left\{ a_l + \ln \frac{M_l^2}{\mu^2} + \frac{m_l^2 - M_l^2 + s}{2s} \ln \frac{m_l^2}{M_l^2} + \frac{q_{\mathrm{cm}}}{\sqrt{s}} \ln \left[\frac{(s+2\sqrt{s}q_{\mathrm{cm}})^2 - (M_l^2-m_l^2)^2}{(s-2\sqrt{s}q_{\mathrm{cm}})^2 - (M_l^2-m_l^2)^2} \right] \right\}, \tag{17}$$

where the regularization scale μ is taken to be 1 GeV , and a_l are subtraction constants, which, together with the low energy parameters of the Lagrangian, were determined from fits to data in Ref. [25–27].

With respect to the final state interaction in the $J/\psi\Lambda$ sector, represented by the diagrams of figure 2(c) and figure 4(a), we recall that two states with strangeness and hidden charm with $J^P = 3/2^-$ and $I = 0$ were found in Refs. [3, 4] as meson-baryon molecules, having pole positions $\sqrt{s} = 4368 - 2.8\mathrm{i}$ and $\sqrt{s} = 4547 - 6.4\mathrm{i}$ and coupling to $J/\psi\Lambda$ states with strength $\mid g_{J/\psi\Lambda} \mid = 0.47$ and 0.61, respectively. The magnitude of each of these couplings is relatively small compared to the coupling of the pole to the main meson-baryon component, which for the lower energy pole is $\bar{D}^*\Xi_c$, with $\mid g_{\bar{D}^*\Xi_c} \mid = 3.6$, while for the higher energy one is $\bar{D}^*\Xi_c'$, with $\mid g_{\bar{D}^*\Xi_c'} \mid = 2.6$. In either case, $\mid g_{J/\psi\Lambda} \mid$ is large enough to create a peak in the mass distribution, as we shall see. As candidate for the strangeness -1 pentaquark, we will consider the state at higher energy since its mass is close to the non-strange pentaquark found in [1]. One must however accept that the mass obtained for these states has uncertainties since, unlike in other sectors,there are no experimental data to constrain the parameters of the theory. We therefore take the nominal value of about $M_R = 4550$ MeV for the mass of the strange pentaquark and will explore the stability of our results to variations of this mass. We shall take $\Gamma_R = 10$ MeV in agreement with the findings of [3, 4]. Our explorations are implemented employing the following Breit-Wigner representation for the $t_{J/\psi\Lambda,J/\psi\Lambda}$ and $t_{\bar{D}^*\Xi_c',J/\psi\Lambda}$ amplitudes

$$t_{J/\psi\Lambda,J/\psi\Lambda} = \frac{g^2_{J/\psi\Lambda}}{M_{J/\psi\Lambda} - M_R + \mathrm{i}\,\frac{\Gamma_R}{2}} \quad \text{and} \quad t_{\bar{D}^*\Xi_c',J/\psi\Lambda} = \frac{g_{\bar{D}^*\Xi_c'}\, g_{J/\psi\Lambda}}{M_{J/\psi\Lambda} - M_R + \mathrm{i}\,\frac{\Gamma_R}{2}} \ . \tag{18}$$

Then the production of the resonance is done through the $J/\psi\Lambda \to J/\psi\Lambda$ and $D^*\Xi_c' \to J/\psi\Lambda$ amplitudes, parametrized through the expressions given in Eqs. (18), as seen in diagrams of figure 2(c) and figure 4(a), respectively, as well as in Eqs. (8), (10). The values of the couplings are $g_{J/\psi\Lambda} = -0.61 - 0.06i$ and $g_{\bar{D}^*\Xi_c'} = 2.61 - 0.13i$. The loop functions $G_{J/\psi\Lambda}$ and $G_{\bar{D}^*\Xi_c'}$ appearing in these equations are taken from [3, 4], where a dimensional regularization method with a scale $\mu = 1000$ MeV was employed, using subtraction constants $a_{J/\psi\Lambda} = a_{\bar{D}^*\Xi_c'} = -2.3$.

3 Results and discussions

All the results will be presented for three different parameterizations, the same as in [27], of the $S = -1$ meson-baryon amplitude with $\eta\Lambda$ in the final state. The first fit, labeled as "WT+NLO", is taken from Ref. [25] (where it is was called "NLO"). In this first work taking into account the $K^-p \to K^+\Xi^-, K^0\Xi^0$ experimental data the authors neglected the direct and crossed Born terms and considered only the lowest-order WT term of the Lagrangian as well as the NLO terms. Such an assumption, which was in agreement with the findings of the other groups working with coupled-channel unitary models up to NLO [28–30], appeared to be not completely correct for the $K^-p \to K\Xi$ reactions [26, 27], which are crucial for determination of the NLO parameters since the WT term has no direct contribution for these transitions. The Born terms, eqs. (14,15), were added to the model, see Refs. [26, 27] for more details, and new fits were performed. The second parameterization, labelled as "WT+NLO+Born", is one of such fits taken from Ref. [26]. The last model, denoted as "WT+NLO+Born η chan" is the latest, and probably most reliable, fit from Ref. [27], which includes the data on the K^-p scattering into $\eta\Lambda$ and $\eta\Sigma^0$ final states. The corresponding sets of parameters and resulting $\chi^2/d.o.f.$ for all three proposed parameterizations can be found in table 1 of Ref. [27].

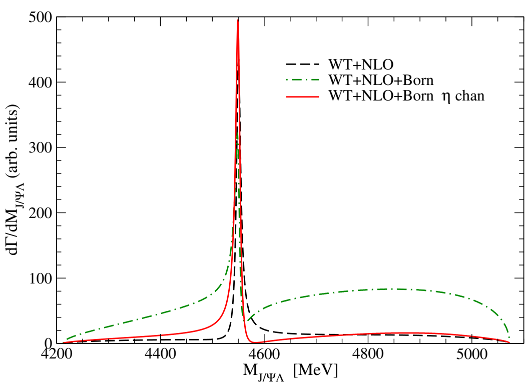

Figure 5: (Color online) Invariant mass distributions of $J/\psi\Lambda$ states produced in the decay $\Lambda_b \to J/\psi\,\eta\Lambda$, obtained for three models discussed in the text.

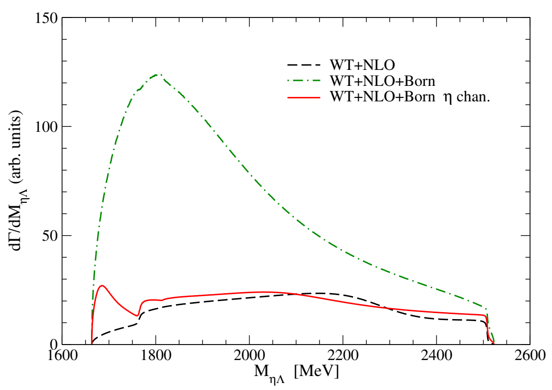

Figure 6: (Color online) The same as figure 5 but for the invariant mass distributions of $\eta\Lambda$ states produced in the decay $\Lambda_b \to J/\psi\,\eta\Lambda$.

We start this section by presenting, in figure 5, the invariant mass distributions of $J/\psi\Lambda$ states produced in the decay $\Lambda_b \to J/\psi\,\eta\Lambda$, obtained from the simpler s-wave weak decay approach of Eq. (8). For all the three models of the $\eta\Lambda$ interaction, the peak of the pentaquark is clearly observed at 4550 MeV, the value of the mass M_R employed in the parametrization of Eq. (18). We can also conclude that the $J/\psi\Lambda$ Breit-Wigner term, eq. (18), has a dominant contribution to the scattering amplitude in the resonance region for all the discussed parameterization.

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The invariant mass distribution of $\eta\Lambda$ pairs is shown in figure 6, where the $J/\psi\Lambda$ resonant structure has disappeared since the invariant $M_{J/\psi\Lambda}$ masses have been integrated out. For "WT+NLO+Born" case we see a broad peaked shape as about 1750 MeV. This structure is not associated to any resonant state since it appears at different energies in different channels and no pole in the complex plane was found either. On the other hand, the new "WT+NLO+Born η chan" model clearly reproduces the resonant structure from $\Lambda(1670)$, see [27] for details.

In the following, results will be presented for only one model of the strong $\eta\Lambda$ interaction, namely the "WT+NLO+Born η chan" one, which is the most reliable fit as argued in Ref. [27]. As we see in figure 5 this fit also gives the smallest signal, and correspondingly it is the most sensitive to other background processes.

The $J/\psi\Lambda$ invariant mass distributions displayed in Figs. 7 and 8, for different values of the pentaquark coupling to $J/\psi\Lambda$ and for different values of the pentaquark mass, respectively, show obvious trends. From figure 7 we can conclude that the pentaquark could be seen over the background even if its coupling to $J/\psi\Lambda$ states were as low as $|\, g_{J/\psi\Lambda} \,| = 0.48$. The unitary approaches of Refs. [3–6] predict values for this coupling in between $0.5 - 1.0$, which make us believe that the strange pentaquark could leave a clear signature in the $J/\psi\Lambda$ mass spectrum.

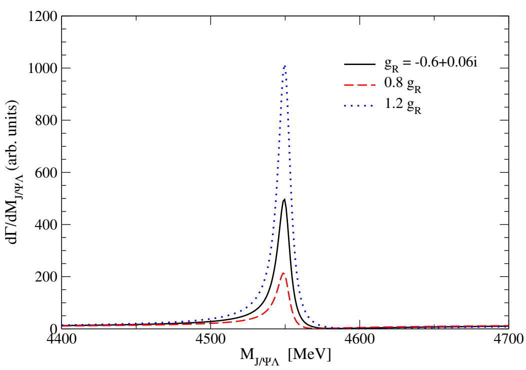

Figure 7: (Color online) Invariant mass distributions of $J/\psi\Lambda$ states produced in the decay $\Lambda_b \to J/\psi\ \eta\Lambda$, obtained for "WT+NLO+Born η chan" model and for different values of the coupling of the pentaquark to $J/\psi\Lambda$.

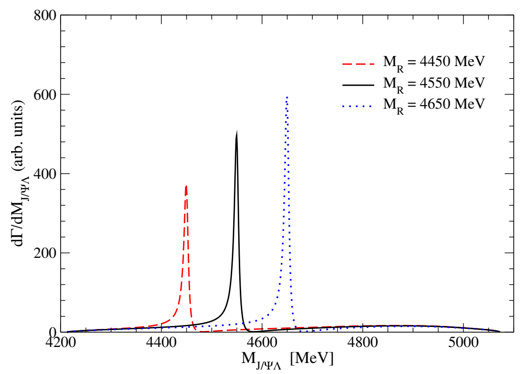

Figure 8: (Color online) Invariant mass distributions of $J/\psi\Lambda$ states produced in the decay $\Lambda_b \to J/\psi\ \eta\Lambda$, obtained for "WT+NLO+Born η chan" model and for different values of the pentaquark mass states.

At this point we would like to address another issue concerning the $J/\psi\,\eta$ interaction. We have explicitly considered the $J/\psi\Lambda$ and $\eta\Lambda$ interactions, and we should, in principle, consider the $J/\psi\,\eta$ one too. Interestingly enough the $J/\psi\,\eta$ system in $L = 0$, with quantum numbers $0^-(1^{+-})$, is one of the decay channels of an axial vector state mostly formed from $D^*\bar{D} + c.c.$ around 3840 MeV, as predicted in [32]. Such a state would be a partner state of the $X(3872)$ with negative C parity. Although it is predicted in some tetra-quark models [33], it has not been found experimentally so far. However, even if a resonance which coupled to $J/\psi\,\eta$ existed, there would be no reflection of it in the $J/\psi\Lambda$ mass distribution in our case because the Dalitz plot is not narrow enough. Indeed, we recall that the $J/\psi\Lambda$ mass distribution is obtained by integrating the double differential cross section of Eq. (11) over $M_{\eta\Lambda}$, or equivalently over $M_{J/\psi\,\eta}$. And thus the $J/\psi\,\eta$ state would enhance the $M_{J/\psi\Lambda}$ spectrum over a relatively wide range of energies, from 4.7 GeV to 5 GeV, not leaving any clear peaked structure [14].

We next explore the influence of the strange pentaquark being initiated by the excitation of a virtual $\bar{D}^{*0}\eta\Xi_c'$ state, followed by the multiple scattering of $\bar{D}^{*0}\Xi_c'$ leading to a final $J/\psi\Lambda$ pair and an η meson. As discussed in Sect. 2, the topology for this decay should lead to a reduced amplitude with respect to that of the $J/\psi\,\eta\Lambda$ case. We implement this phenomenologically through the parameter β, as seen in Eq. (10), which is given the values $-0.5, -0.25, 0.0, 0.25, 0.5$ accounting also for different relative sign cases. On the other hand, the smallness of β can be compensated by the fact that the coupling strength of the pentaquark to $\bar{D}^{*0}\Xi_c'$ states is a factor 4 larger than that to $J/\psi\Lambda$.

The obtained results are displayed in figures 9 and 10 for the negative and positive β values correspondingly. In all the cases the pentaquark signal clearly dominates over the background; for the negative values of β the signal is enhanced strongly due to a constructive interference between both mechanisms of the pentaquark production, since $g_{J/\psi\Lambda}$ and $g_{\bar{D}^*\Xi_c'}$ have opposite signs.

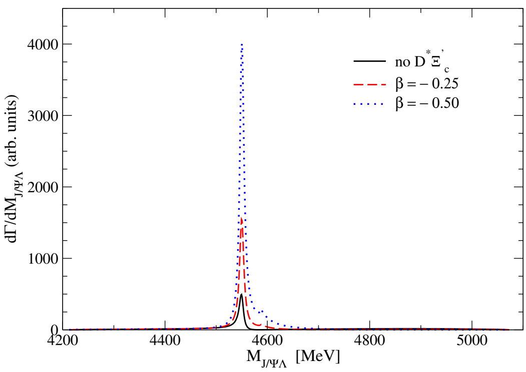

Figure 9: (Color online) Invariant mass distributions of $J/\psi\Lambda$ states produced in the decay $\Lambda_b \rightarrow J/\psi\ \eta\Lambda$, obtained for "WT+NLO+Born η chan" model and for different strengths of the $\bar{D}^{*0}\Xi_c'$ intermediate state contribution.

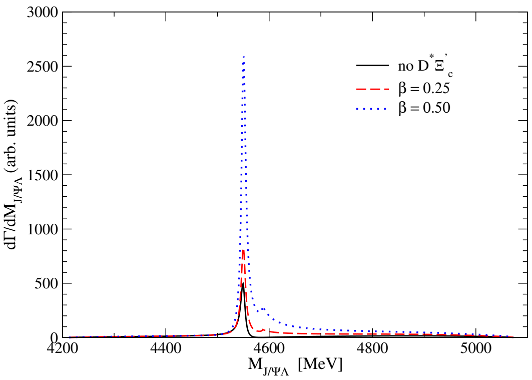

Figure 10: (Color online)The same as figure 9, but for positive β parameter, se eq. (10).

To conclude, the recent finding of two structures in the $J/\psi p$ invariant mass distribution of the $\Lambda_b \rightarrow J/\psi K^- p$ decay, associated to two pentaquark states, together with its plausible explanation in terms of a previously predicted hidden charm baryon molecular state, prompted us to study the decay of the Λ_b into $J/\psi\,\eta\Lambda$ final states [14]. The $\Lambda_b \rightarrow J/\psi\,\eta\Lambda$ decay, being a coupled channel of the $\Lambda_b \rightarrow J/\psi K^- p$ one, will occur with similar strength and one could observe, in the $J/\psi\Lambda$ invariant mass spectrum, possible strange partners of the two non-strange pentaquark states reported by the LHCb collaboration.

With all the options and uncertainties considered, we see that the existence of the molecular state predicted in Refs. [3, 4] leads to a clear signal in the $J/\psi\Lambda$ invariant mass spectrum in the range of about 4450-4650 MeV. Thus we reconfirm the finding of Ref. [14] where different models for the $S = -1$ meson-baryon amplitude with $\eta\Lambda$ in the final state have been considered. This gives us confidence that such an experiment should result into a successful proof of the existence of this new state and we encourage the experimental analysis of this decay channel.

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