On the Origin of the XYZ Mesons

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Abstract. In this talk we present a mechanism giving rise to exotic XYZ four-quark states in the meson spectra within a constituent quark model approach. We discuss its generalization to five-quark states in the heavy baryon sector. Finally, we revise some other works in the literature and experimental data where this mechanism may be working.

INTRODUCTION

Before the discovery of the $X(3872)$ [1], the hadronic experimental data were classified either as $qqq$ or $qar{q}$ states according to $SU(3)$ irreducible representations based on Gell-Mann conjecture [2]. However, since 2003 more than twenty newly observed meson resonances reported by different experimental collaborations appeared, close to a two-meson threshold, presenting properties that make a simple quark-antiquark structure unlikely [3], the so-called XYZ states. Although this observation could be coincidental due to the large number of thresholds in the energy region where the XYZ mesons have been reported, it could also point to a close relation between some particular thresholds and resonances contributing to the standard quark-antiquark heavy meson spectroscopy.

More than a decade has elapsed since the discovery of the $X(3872)$, and no compelling explanation for the pattern of XYZ mesons has emerged. Several alternatives have been proposed in the literature to address these XYZ states [4], being the most common ones conventional quarkonium [5], which consists of a color-singlet heavy quark-antiquark pair: $(Q\bar{Q})_1$; meson-meson molecules [6], which consists of color-singlet $Q\bar{q}$ and $Q\bar{q}$ mesons bound by hadronic interactions: $(Q\bar{q})_1 + (\bar{Q}q)_1$; quarkonium hybrids [7], which consists of a color-octet $Q\bar{Q}$ pair to which a gluonic excitation is bound: $(Q\bar{Q})_8 + g$; four-quark states [8], which consists of a $Q\bar{Q}$ pair and a light quark $q$ and antiquark $\bar{q}$ bound by interquark potentials into a color singlet: $(Q\bar{Q}q\bar{q})_1$; and diquarkonium [9], which consists of a color-antitriplet $Qq$ diquark and a color-triplet $\bar{Q}\bar{q}$ diquark bound by the QCD color force: $(Qq)_{\bar{3}} + (\bar{Q}q)_{3}$. However, none of the models that has been proposed provide a plausible pattern for all the $XYZ$ mesons that have been observed.

In this talk we analyze heavy hadron spectroscopy beyond open flavor thresholds considering higher order Fock space components looking for a general pattern of the $XYZ$ states. We highlight the pivotal role played by hadron-hadron thresholds in heavy hadron spectroscopy and under which conditions, if any, they could cause a resonance to appear. We will analyze the thresholds open for four-quark states contributing to heavy meson spectroscopy and we will show how these thresholds may entangle so a bound four-quark state or a resonance may emerge and the conditions required for that [10]. We will extend our arguments to the heavy baryon spectra, where one could also find contributions with a involved structure, such as compact five–quark states beyond simple $ND$ resonances [11, 12]. The study of these contributions requires from a full coupled-channel approach including all possible physical states contributing to a given set of quantum numbers $(T, J)$, as has been demonstrated in Ref. [13] for the charmonium spectrum.
FORMALISM

Standard mesons ($q\bar{q}$) and baryons ($qqq$) are the only clusters of quarks where it is not possible to construct a color singlet using a subset of their constituents. This, however, is not the case for multiquark combinations, and in particular for four–quark and five–quark states addressing the meson and baryon spectra, respectively. Thus, when dealing with higher order Fock space contributions to hadron spectroscopy, one has to discriminate between possible multiquark bound states or resonances and simple pieces of the hadron–hadron continuum. For this purpose, one has to analyze the two–hadron states that constitute the threshold for each set of quantum numbers. These thresholds have to be determined assuming quantum number conservation within exactly the same scheme (parameters and interactions) used for the multiquark calculation. If other models, parametrizations or experimental masses are used, then multiquark states might be misidentified as members of the hadron spectra while being simple pieces of the continuum.

Given a general four–quark state, $(q_1q_2\bar{q}_3\bar{q}_4)$, two different thresholds are allowed, $(q_1q_2\bar{q}_3)q_4$ and $(q_1\bar{q}_2)(q_3\bar{q}_4)$. If the four–quark system contains identical quarks, like for instance $(\bar{Q}Q\bar{n}\bar{n})$ (in the following $n$ stands for a light quark and $\bar{Q}$ for a heavy $c$ or $b$ quark), the two thresholds are identical, i.e., $(\bar{Q}Q)(\bar{n}\bar{n})$. The importance of this particular feature lies on the fact that a modification of the four–quark interaction would not necessarily translate into the mass of the two free-meson state. Therefore, the unique necessary condition required to have a four–quark bound state would be the existence of a sufficiently attractive interaction between quarks that do not coexist in the two free-meson states. This hypothesis was demonstrated by means of the Lippmann–Schwinger formalism in Ref. [13], concluding the existence of a single stable isoscalar doubly charmed meson with quantum numbers $J^P = 1^+$. For those cases containing a heavy quark and its corresponding heavy antiquark ($Q\bar{Q}\bar{n}\bar{n}$) the situation is remarkably different. Two different thresholds are allowed, namely $(\bar{Q}Q)(\bar{n}\bar{n})$ and $(\bar{n}\bar{n})(Q\bar{Q})$. It has been proved [14] that ground state solutions of the Schrödinger ($q_1q_2\bar{q}_3\bar{q}_4$) two–body problem are concave in $(m_{q_1}^{-1} + m_{\bar{q}_1}^{-1})$ and hence $\bar{M}_{Q\bar{Q}} + \bar{M}_{Q\bar{n}} > M_{\bar{Q}\bar{Q}} + M_{n\bar{n}}$. This property is enforced both by nature and by all models in the literature unless forced to do otherwise. Although this relation among ground-state masses makes the assumption of a strictly flavor independent potential, one should bear in mind that ground states of heavy mesons are perfectly reproduced by a Cornell-like potential [15], that it is flavor independent. The color-spin dependence of the potential would go in favor of this relation for ground states (spin zero) because the color-spin interaction is attractive for spin zero and comes suppressed as $1/(m_\ell m_{\ell'})$, making even lighter the mesons on the right hand side. Regarding the spin independent part, the binding of a coulombic system is proportional to the reduced mass of the interacting particles. Thus, for a two-meson threshold with a heavy-light light-heavy quark structure, the binding of any of the two mesons is proportional to the reduced mass of each meson, being close to the mass of the light quark. However, if the two-meson state presents a heavy-heavy light-light quark structure, the binding of the heavy-heavy meson increases proportionally to the mass of the heavy particle while that of the light-light meson remains constant, becoming this threshold lighter than the heavy-light light-heavy two-meson structure. Thus, it implies that in all relevant cases the lowest two-meson threshold for any $(Q\bar{Q}\bar{n}\bar{n})$ state will be the one made of quarkonium-light mesons, i.e., $(\bar{Q}Q)(\bar{n}\bar{n})$ (see Fig. 1 of Ref. [16]). The interaction between the heavy, $(Q\bar{Q})$, and light, $(\bar{n}\bar{n})$, mesons forming the lowest threshold is almost negligible, due to the absence of a light pseudoscalar exchange mechanism between them [6]. Hence, any attractive effect in the four–quark system must have its origin in the interaction of the higher channel $(\bar{Q}Q)(n\bar{Q})$ or due to the coupled channel effect of the two thresholds, $(\bar{Q}Q)(\bar{n}\bar{n}) \leftrightarrow (\bar{Q}Q)(n\bar{Q})$ [17].

Similar arguments could be used in the case of the heavy baryon spectra. Given a general five–quark state contributing to the heavy (charm or bottom) baryon spectrum, $(nnn\bar{Q}\bar{n})$, two different thresholds are allowed, $(nnn\bar{Q})\bar{n}$ and $(nn\bar{Q})(\bar{n}\bar{n})$. A straightforward generalization of the concave behavior in $(m_{q_1}^{-1} + m_{\bar{q}_1}^{-1})$ of the ground state solutions of the Schrödinger ($q_1q_2\bar{q}_3\bar{q}_4$) two–body problem to the five–quark system could be obtained within a quark-diquark model if $m_{q_1} \leq m_{q_2} \leq m_{\bar{q}_1}$. Then $M_{q_1\bar{q}_1} + M_{q_2\bar{q}_1} \leq M_{q_1\bar{q}_1} + M_{q_2\bar{q}_1}$, because the intervals in $1/\mu$ of the left hand side and right hand side have the same middle, but the left hand side one is wider that the right hand side one. Now, in a crude quark-diquark model, one can translate this as $M_{q_1q_2\bar{q}_1} + M_{q_3\bar{q}_1} \leq M_{q_1q_2\bar{q}_1} + M_{q_3q_4\bar{q}_1}$, as it is observed in Fig. 5 of Ref. [12], except for the higher spin states where the angular momentum coupling rules impose further restrictions.

In those energy regions where bound and unbound solutions of the multiquark hamiltonian coexist, methods based on infinite expansions become inefficient to hunt a bound state close to an unbound solution, because too many basis states would be required to disentangle them. The most interesting case where this may happen is in the vicinity of a two–hadron threshold, because both the two free–hadron state and a feasible slightly bound multiquark state are

\begin{align*}
M_{D^*} + M_{D_s^*} & = 4014 \text{ MeV} > M_{J/\psi} + M_c = 3879 \text{ MeV} \\
\text{We thank to J. M. Richard for this simple and nice argument that does not make any assumption on the shape of the interaction, linear or not, although it assumes a quark-diquark ansatz.}
\end{align*}
solutions of the same hamiltonian. Such cases have been studied by means of the Lippmann-Schwinger equation in Ref. [18], looking at the Fredholm determinant $D_T(E)$ at zero energy [19]. If there are no interactions then $D_T(0) = 1$, if the system is attractive then $D_T(0) < 1$, and if a bound state exists then $D_T(0) < 0$. All states made of $S$ wave $(Q\bar{n})-(n\bar{Q})$ mesons up to $J = 2$ were scrutinized. A few channels were found to be slightly attractive, $D\bar{D}$ with $(I)J^{PC} = (0)^{0}$$^{+}$, $D\bar{D}′$ with $(0)^{1}$$^{+}$ and $D′\bar{D}′$ with $(0)^{0}$$^{+}$, $(0)^{2}$$^{+}$, and $(1)^{2}$$^{+}$, close to the results of Ref. [6]. However, the only bound state appeared in the $(I)J^{PC} = (0)^{1}$$^{+}$ channel as a consequence of the coupling between $D\bar{D}′$ and $J/\Psi\omega$ two–meson channels.

The conclusions of Refs. [8] and [18] point to a convoluted four–quark molecular structure with a dominant $D\bar{D}′$ component for the $X(3872)$. However, this is not the only supernumerary state that has appeared in the charm and bottom sectors during the last years. A comprehensive list of such $XYZ$ mesons and their properties can be found in Ref. [3]. 20 states have been reported by different experimental collaborations in the charmonium sector above the $D\bar{D}$ threshold, 15 of them neutral and 5 charged. In the bottom sector 2 charged and 1 neutral state have been reported. Of those 23 states only 8 have been observed independently by two different collaborations and with significance greater than $5\sigma$: $X(3872), X(3915), X(2200), G(3900), Y(4140), Y(4260), Y(4360)$, and $Z'(3900)$ (see Table I of Ref. [3]). Therefore, although some of them might not resist cross-check examination by independent experimental collaborations, others are clearly established as real resonances and therefore they have to be accounted for in any description of the meson spectra. While some of them, like the $X(2200)$, seem to fit nicely within a naive quark-antiquark scheme, others do not.

**RESULTS AND DISCUSSION**

When four-quark components are considered in the wave function of charmonium, there are 72 $cc\bar{n}\bar{n}$ combinations of quantum numbers for total orbital angular momentum $L < 3$. Therefore, the question is not whether it is possible to design a model, or a formalism, able to match one of the newly observed $XYZ$ states with a particular set of these quantum numbers but to understand where the attraction comes from and to explain the systematic that predicts where, if anywhere, experimentalists and theoreticians alike should look into. Since the lowest threshold interaction is rather weak, the possibility to obtain bound states may only stem from the vicinity of an attractive $(c\bar{n})(n\bar{c})$ threshold coupled sufficiently as to bind the system as it occurs with the $X(3872)$ [18]. Although for heavier mesons interactions are more attractive [16], the effect of channel coupling may not be enough to favor binding. Thus, to check the efficiency of this mechanism, we solved with the HH formalism the bottom counterpart of the $X(3872)$, $b\bar{b}n\bar{n}$ with quantum numbers $L = 0, S = 1, I = 0, C = +1,$ and $P = +1$. Within the CQCM of Ref. [20] the corresponding lowest thresholds, $b\bar{b}n\bar{b}$ (10611 MeV) and $\Upsilon\omega$ (10155 MeV), are 456 MeV apart. We show in Fig. 1 (upper panel) the convergence pattern of the energy of the four-quark system as a function of the hyperangular momenta $K$. It can be clearly seen how the energy of the four–quark system (red line) is converging to the lowest threshold $\Upsilon\omega$ (horizontal blue line), what is a sharp signal of an unbound state. One could however play around with the model parameters to almost degenerate the highest $b\bar{b}n\bar{n}$ and $n\bar{b}b\bar{n}$ both thresholds by adding attraction in the heavy-light two-meson state that would also give rise to a strong $(Q\bar{Q})(n\bar{n}) \leftrightarrow (\bar{Q}\bar{Q})(n\bar{n})$ coupling. To neatly illustrate this conclusion we return to the four–quark state $b\bar{b}n\bar{n}$ with quantum numbers $L = 0,$ $I = 0,$ $(0)1^{+}$, the effective wave function of the $b\bar{b}n\bar{n}$ state is weak and therefore the coupled channel effects are simultaneously increased, bound states may appear for a subset of quantum numbers. Hence, threshold vicinity is a required but not sufficient condition to bind a four–quark state. An additional condition is required to allow the emergence of such bound states. Such condition is the existence of an attractive interaction in the higher $(Q\bar{Q})(n\bar{n})$ two–meson system that would also give rise to a strong $(Q\bar{Q})(n\bar{n}) \leftrightarrow (\bar{Q}\bar{Q})(n\bar{n})$ coupling. 3We have slightly increased the $a_{c}(bn)$ strong coupling constant from 0.55 to 0.85, what would move the gap between thresholds without changing the lowest threshold.
\( S = 1, I = 0, C = +1, \) and \( P = +1. \) In Refs. [16, 18] it was proved that the interaction provided by the CQCM model is attractive for these quantum numbers in the charm sector, although does not present a \( D\bar{D}^* \) bound state. In the bottom sector the attraction is enhanced [16]. Thus, we have solved the four–body problem as a function of the threshold energy difference, \( \Delta = E((b\bar{n})(n\bar{b})) - E((b\bar{b})(n\bar{n})) \), ranging from 500 MeV to \(-200 \) MeV. We show in Fig. 2 the energy normalized to the mass of the lowest two-meson threshold for each particular case, i.e., values smaller than 1 will point to a bound state and those larger to an unbound state. In this case, the results can be separated into two distinct categories: (i) \( \Delta \gtrsim 50 \) MeV and (ii) \( \Delta \lesssim 50 \) MeV. When the thresholds are separated by more than 50 MeV, the attractive interaction in the \( B\bar{B}^* \) system and the coupled channel effect is not sufficient to overcome the threshold energy gap, and therefore the four–quark system evolves to an unbound two–meson state. However, when the thresholds get closer, even reversed for \( \Delta < 0, \) the system becomes a compact four–quark state. Of particular interest are those cases where \( \Delta = 50 \) MeV in Fig. 2. In this case the attraction in the higher channel together with the coupled channel effect barely overcomes the threshold energy difference, and hence its wave function becomes strongly entangled. This would generate a molecular state just close to threshold.

It should also be noted that \( \Delta_0, \) the \( \Delta \) value for which \( \text{Energy}(4q)/E(\text{lowest}) \) is equal to one, will change depending on the particular set of quantum numbers considered. However, \( \Delta_0 = 0 \) implies that no binding energy is provided by the upper threshold and the off–diagonal terms and therefore such solution will not correspond to a bound state. In that case our calculation will be simply providing a piece of the meson–meson unbound continuum.

The same argument has been drawn in Ref. [21] where the supercharmonium states have been introduced, experimental resonances that appear to contain a \((c\bar{c})\) pair, but have other properties that preclude a description in terms
of only \((c\bar{c})\) (idealized charmonium) basis states. Four-quark supercharmonium configurations occur very near, or even below, the lowest S-wave threshold for production of meson pairs. These states evade Coleman’s argument in 1979 [22], who used the \(1/N_c\) expansion of QCD to show that four-quark color singlets tend to propagate as pairs of mesons. One of these superchrmonium states could be the \(Z(4475)\), that may appear as a linear superposition of \(D\bar{D}^*\) mesons with one of them in a radially excited state \(2^+\). In a simple one-pion exchange model the \(D \rightarrow D + \pi\) vertex is forbidden while the \(D^* \rightarrow D + \pi\) is allowed. This gives rise to weak diagonal interactions but a strong mixing between two thresholds, \(D(1S)\bar{D}^*(2S)\), \(4482\text{ MeV}/c^2\), and \(D^*(1S)\bar{D}(2S)\), \(4433\text{ MeV}/c^2\). Besides, the small mass difference between the thresholds due to the diminishing of the hyperfine splitting when increasing the radial excitation, leads to an amplification of the binding. The mechanism would be the same working for the \(X(3872)\) but with a complete degeneracy between \(D(1S)\bar{D}^*(1S)\) and \(D^*(1S)\bar{D}(1S)\).

A similar situation could be found in the case of the heavy baryon spectrum. An important source of attraction might be the coupled-channel effect of the two thresholds, \((nnn)(Q\bar{n}) \leftrightarrow (nnQ)(n\bar{n})\) [17]. Thus, to check the efficiency of this mechanism, we have performed a calculation considering all physical channels, \((nnn)(Q\bar{n})\) and \((nnQ)(n\bar{n})\). When the \((nnn)(Q\bar{n})\) and \((nnQ)(n\bar{n})\) thresholds are sufficiently far away, the coupled-channel effect is small, and bound states are not found. However, when the thresholds move closer, the coupled-channel strength is increased, and bound states may appear for a subset of quantum numbers. Hence, threshold vicinity is again a required but not sufficient condition to bind a five–quark state. Under these conditions, there are the channels with high spin \(J^P = 5/2^-\) the only ones that may lodge a compact five-quark state for all isospins [12]. The reason stems on the reverse of the ordering of the thresholds, being the lowest threshold \((nnn)(Q\bar{n})\) the one with the more attractive interaction. Of particular interest is the \((T)J^P = (2)5/2^-\) state, that survives the consideration of the break apart thresholds. It may correspond to the \(\Theta_c(3250)\) pentaquark found by the QCD sum rule analysis of Ref. [23] when studying the unexplained structure with a mass of \(3250\text{ MeV}/c^2\) in the \(\Sigma^*_c\pi^-\pi^-\) invariant mass reported recently by the BABAR Collaboration [24]. Such state could therefore be a consequence of the close-to-degeneracy of the lowest thresholds with \(I = 2\) and \(J^P = 5/2^-\), \(\Delta D^*\) and \(\Sigma_c\rho\) and the attractive interaction of the \(\Delta D^*\) system [12]. It could be detected by the propagation of \(D\) mesons in nuclear matter as an \(S\) wave \(\Delta D^*\) system and it thus constitutes a challenge for the \(P\)ANDA Collaboration.

**CONCLUSIONS**

To summarize, we have presented a plausible mechanism for the origin of the \(XYZ\) mesons in the heavy meson spectra within a standard quark-model picture. Its generalization to the heavy baryon sector has been analyzed. The existence of open flavor two–hadron thresholds is a feature of the heavy hadron spectra that needs to be considered as a relevant ingredient into any description of the plethora of new states reported in heavy meson and baryon spectroscopy. They might be, at a first glance, identified with simple quark-antiquark or three-quark states, however in some cases their energies and decay properties do not match such oversimplified picture. Our results prove the relevance of higher or-
der Fock space components through the allowed two-hadron thresholds. On the one hand, one has the lower $(Q\bar{Q})(n\bar{n})$ and $(nn)(Q\bar{n})$ systems, made by almost noninteracting hadrons, that constitutes the natural breaking apart end-state. On the other hand, the higher $(Q\bar{n})(n\bar{Q})$ and $(nnQ)(n\bar{n})$ systems appear. When there is an attractive interaction characterizing the upper systems combined with a strong enough coupling, together with the vicinity of the two allowed thresholds, a multiquark bound state may emerge. As one can see the mechanism proposed is restrictive enough as not predicting a proliferation of bound states when explaining the existence of an hypothetical molecular structure. Once this is performed, the present experimental effort with ongoing experiments at BESIII, current analyses by the LHC collaboration and future experiments at Belle II and Panda together with the very impressive results that are being obtained by lattice gauge theory calculations [25] may confirm the theoretical expectations of our quark-model calculation pattern that will provide with a deep understanding of low-energy realizations of QCD.

ACKNOWLEDGMENTS

This work has been partially funded by the Spanish Ministerio de Educación y Ciencia and EU FEDER under Contract No. FPA2013-47443, by the Spanish Consolider-Ingenio 2010 Program CPAN (CSD2007-00042) and by Generalitat Valenciana Prometeo/2009/129. A.V. thanks financial support from the Programa Propio I of the University of Salamanca.

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