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# Hadronic resonances enhanced by thresholds

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ABSTRACT

We present a neat example of a meson–baryon system where the vicinity of two different thresholds enhances the binding of a hadronic resonance, a pentaquark. As a consequence the pattern of states may change when moving among different flavor sectors, what poses a warning on naive extrapolations to heavy flavor sectors based on systematic expansions. For this purpose we simultaneously analyze the $N\bar{D}$ and $NB$ two-hadron systems looking for possible bound states or resonances. When a resonance is controlled by a coupled-channel effect, going to a different flavor sector may enhance or diminish the binding. This effect may, for example, generate significant differences between the charmonium and bottomonium spectra above open-flavor thresholds or pentaquark states in the open-charm and open-bottom sectors.



The interpretation of the recently discovered baryonic states at the LHCb, $P_c(4380)^+$ and $P_c(4450)^+$ [1], as well as some of the exotic mesonic states discovered in the hidden-charm or hidden-beauty sector is still puzzling [2,3]. A common feature of all these states seems to be the proximity of their masses to two-hadron thresholds. Their naive description as simple baryon–meson or meson–meson resonances gave rise to predictions of bound states in heavier flavor sectors by different spectroscopic models, like those based on the traditional meson theory of the nuclear forces or resorting to heavy quark symmetry arguments [4,5].

In a recent paper [6] a mechanism to explain the stability or metastability of the exotic mesonic states discovered in the hidden-charm or hidden-beauty sector was proposed. It was pointed out how two effects have to come together to allow for the formation of a bound state above open-flavor thresholds: the presence of two nearby thresholds and a strong coupling between them, in spite of the fact that the diagonal interactions contributing to this state are not strong. Thus, such mechanism, as long as it is possible avoids the risk of proliferation of states appearing in some quark-model calculations.

For the $X(3872)$ this is well plausible [7]. It was pointed out that two of the possible dissociation thresholds are almost exactly degenerate, the one corresponding to spin-singlet charmed meson plus a spin-triplet anti-charmed meson (or conjugate), and the one made of a light vector-meson and a charmonium vector-meson. For instance, in the flavor-SU(3) limit, the $H$ dibaryon benefits of the degeneracy of the $\Lambda\Lambda$ and $N\Xi$ thresholds, and is found stable in some model calculations, whereas for broken SU(3) the degeneracy is lost and the $H$ dibaryon becomes unstable in the same models [8].

This idea had been already anticipated in a qualitative analysis of the possible dissociation thresholds of four-quark systems with a $Q\bar{Q}n\bar{n}$ structure (in the following $n$ stands for a light quark and $Q$ for a heavy $c$ or $b$ quark), making stringent predictions as the non-existence of a bottom partner for the $X(3872)$ or the existence of exotic doubly heavy mesons [9]. While the first prediction seems to survive experiment in contrast to those of other theoretical models [4,5], the second is still awaiting for an experimental effort [10]. We wonder if there could be a neat example where one could think of some degeneracy of two baryon–meson thresholds leading to exotic or crypto-exotic baryons, as it may be the case for some of the exotic meson states [9]. The existence of an exotic state in a given flavor sector can not be naively generalized to other flavor sectors in case of loss of the vicinity of the thresholds. In a similar manner, its non-existence in a particular flavor sector does not exclude its presence in different flavor sectors.

In this letter we discuss a relevant example of a five-quark state in the $NB$ two-hadron system that clearly exemplifies the importance of the mechanism we have previously presented. It should be considered in the phenomenological analysis of the recently reported pentaquark states and may serve as a guideline for the study of the pattern of exotic states in the baryon and meson sectors [2]. Our findings come up in the shadow of a previous study

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**Table 1**
Interacting baryon–meson channels in the isospin–spin $(T, J)$ basis.

| | $T=0$ | $T=1$ | $T=2$ |
|---|---|---|---|
| $J=1/2$ | $NB-NB^*$ | $NB-NB^*-\Delta B^*$ | $\Delta B^*$ |
| $J=3/2$ | $NB^*$ | $NB^*-\Delta B-\Delta B^*$ | $\Delta B-\Delta B^*$ |
| $J=5/2$ | – | $\Delta B^*$ | $\Delta B^*$ |

of a different two-hadron system, $N\bar{D}$, whose generalization to the bottom case gave rise to an a priori unexpected result that remarks the effect of the almost degeneracy of two different baryon–meson thresholds.

In Ref. [11] we studied the $N\bar{D}$ system by means of a chiral constituent quark model. Our main motivation at that time was the study of the interaction of $D$ mesons with nucleons which is a goal of the P̄ANDA Collaboration at the European facility FAIR [12]. Thus, our theoretical study was a challenge to be tested at the future experiments. In this letter we perform a parallel study of the $NB$ system (with a similar quark structure, $nnnn\bar{Q}$) looking for similarities and differences with respect to the $N\bar{D}$ system. Our objective is to highlight a particular case where the vicinity of thresholds will enhance the binding of the baryon–meson system disrupting the number and the ordering of states obtained in the charm sector. Although the conclusions of this study aim to be independent of the particular details of the interacting model used, we for instance made use of the chiral constituent quark model (CCQM) of Ref. [13]. It was proposed in the early 90's in an attempt to obtain a simultaneous description of the nucleon–nucleon interaction and the baryon spectra [14]. It was later on generalized to all flavor sectors giving a reasonable description of the meson and baryon spectra. The model is based on the assumption that the light-quark constituent mass appears because of the spontaneous breaking of the original $SU(3)_L \otimes SU(3)_R$ chiral symmetry at some momentum scale. In this domain of momenta, quarks interact through Goldstone boson exchange potentials. QCD perturbative effects are taken into account through the one-gluon-exchange potential. Finally, it incorporates confinement as dictated by unquenched lattice QCD calculations. A detailed discussion of the model can be found in Refs. [13,14].

The systems under study, $N\bar{D}$ and $NB$, do not present quark–antiquark annihilation complications that may obscure the predictions of a particular model under some non-considered dynamical effects. They contain a heavy antiquark, what makes the interaction rather simple. The quark-model used provides parameter-free predictions for the interaction in a baryon–meson system with charm $-1$ or bottom $+1$. Besides, the existence of identical light quarks in the two hadrons generates quark-Pauli effects in some particular channels, what gives rise to an important short-range repulsion due to lacking degrees of freedom to accommodate the light quarks [14].

To study the possible existence of exotic states made of a light baryon, $N$ and $\Delta$, and a charmed meson, $\bar{D}$ and $\bar{D}^*$, or a bottom meson, $B$ and $B^*$, we solve the Lippmann–Schwinger equation for negative energies looking at the Fredholm determinant $D_F(E)$ at zero energy [15]. If there are no interactions then $D_F(0)=1$, if the system is attractive then $D_F(0)<1$, and if a bound state exists then $D_F(0)<0$. This method permitted us to obtain robust predictions even for zero-energy bound states, and gave information about attractive channels that may lodge a resonance [7]. We consider a baryon–meson system $Q_iR_j$ ($Q_i=N$ or $\Delta$ and $R_j=\bar{D}$ or $\bar{D}^*$ for charm $-1$ and $R_j=B$ or $B^*$ for bottom $+1$) in a relative $S$ state interacting through a potential $V$ that contains a tensor force. Then, in general, there is a coupling to the $Q_iR_j$ $D$ wave. Moreover, the baryon–meson system can couple to other baryon–meson states, $Q_kR_m$. We show in Table 1 the coupled channels in the isospin–spin $(T, J)$ basis for the $NB$ system (for the $N\bar{D}$ system one would replace $B$ by $\bar{D}$ and $B^*$ by $\bar{D}^*$). Let us briefly sketch the method to look for bound state solutions using the Fredholm determinant. If we denote the different baryon–meson systems as channel $Q_iR_j \equiv A_n$, the Lippmann–Schwinger equation for the baryon–meson scattering becomes

$$
\begin{aligned}
& t^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}_{\alpha\beta;TJ}(p_\alpha,p_\beta;E) \\
& = V^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}_{\alpha\beta;TJ}(p_\alpha,p_\beta) \\
& + \sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\int_0^\infty p_\gamma^2 dp_\gamma V^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}_{\alpha\gamma;TJ}(p_\alpha,p_\gamma) \\
& \times G_\gamma(E;p_\gamma) t^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}_{\gamma\beta;TJ}(p_\gamma,p_\beta;E) \ , \ \alpha,\beta=A_1,A_2,\cdots \ ,
\end{aligned} \tag{1}
$$

where $t$ is the two-body scattering amplitude, $T$, $J$, and $E$ are the isospin, total angular momentum and energy of the system, $\ell_\alpha s_\alpha$, $\ell_\gamma s_\gamma$, and $\ell_\beta s_\beta$ are the initial, intermediate, and final orbital angular momentum and spin, respectively, and $p_\gamma$ is the relative momentum of the two-body system $\gamma$. The propagators $G_\gamma(E;p_\gamma)$ are given by

$$
G_\gamma(E;p_\gamma) = \frac{2\mu_\gamma}{k_\gamma^2 - p_\gamma^2 + i\epsilon} \,, \tag{2}
$$

with

$$
E = \frac{k_\gamma^2}{2\mu_\gamma} \,, \tag{3}
$$

where $\mu_\gamma$ is the reduced mass of the two-body system $\gamma$. For bound-state problems $E<0$ so that the singularity of the propagator is never touched and we can forget the $i\epsilon$ in the denominator. If we make the change of variables

$$
p_\gamma = d\,\frac{1+x_\gamma}{1-x_\gamma} , \tag{4}
$$

where $d$ is a scale parameter, and the same for $p_\alpha$ and $p_\beta$, we can write Eq. (1) as

$$
\begin{aligned}
& t^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}_{\alpha\beta;TJ}(x_\alpha,x_\beta;E) \\
& = V^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}_{\alpha\beta;TJ}(x_\alpha,x_\beta) \\
& + \sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\int_{-1}^{1} d^2\left(\frac{1+x_\gamma}{1-x_\gamma}\right)^2 \frac{2d}{(1-x_\gamma)^2}dx_\gamma \\
& \times V^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}_{\alpha\gamma;TJ}(x_\alpha,x_\gamma)\, G_\gamma(E;p_\gamma)\, t^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}_{\gamma\beta;TJ}(x_\gamma,x_\beta;E) \,.
\end{aligned} \tag{5}
$$

We solve this equation by replacing the integral from $-1$ to 1 by a Gauss–Legendre quadrature which results in the set of linear equations

$$
\begin{aligned}
& \sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\sum_{m=1}^{N} M^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}_{\alpha\gamma;TJ}(E)\, t^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}_{\gamma\beta;TJ}(x_m,x_k;E) \\
& = V^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}_{\alpha\beta;TJ}(x_n,x_k) \,,
\end{aligned} \tag{6}
$$

with