Chromomagnetic mechanism for the X(3872) resonance

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The chromomagnetic interaction, with proper account for flavor-symmetry breaking, is shown to explain the mass and coupling properties of the X(3872) resonance as a $J^{PC} = 1^{++}$ state consisting of a heavy quark-antiquark pair and a light one. It is crucial to introduce all the spin-color configurations compatible with these quantum numbers and diagonalize the chromomagnetic interaction in this basis. This approach thus differs from the molecular picture $D\bar{D}^*$ and from the diquark–anti-diquark picture.

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In recent months, several intriguing new hadron states have been announced. Some of them are rather controversial since they were tentatively seen in some experiments and not in others. On the other hand, the X(3872) reported by the Belle Collaboration [1] can be considered as well established, since it has been confirmed at *BABAR* and at the Fermilab collider [2]. There are some indications that it could have $J^{PC} = 1^{++}$ quantum numbers [3], in particular, it is not seen in a $\gamma\gamma$ search at CLEO [4].

Several theoretical explanations have been proposed for the X(3872). It could be mainly a charmonium excitation $(c\bar{c})$, though none of the partial-wave assignments $^{2s+1}\mathrm{L}_J$ actually matches the predictions of charmonium models tuned to fit the known levels [3,5].

A hybrid scenario also has been suggested for this state, or for the other states discovered in this region: X(3940) [6] or Y(4260) [7]. Excitations of the string linking the quark to the antiquark, or, in the QCD language, of the gluon field, were proposed long ago on the basis of some models and confirmed by Lattice QCD. A signature of this would be a decay with at least one orbitally excited meson, for instance $D^{**} + \bar{D}$ [8].

The Yukawa mechanism is not restricted to the nucleon-nucleon system and holds for any pair of hadrons containing light quarks. In particular, pion exchange, if allowed and attractive, can be just strong enough to bind heavy hadrons to form a deuteronlike compound. Remarkably, this mechanism led some authors to predict the existence of $D\bar{D}^*$ + c.c. states and when the X(3872) was found very close to the $D\bar{D}^*$ threshold, it was considered as a very natural candidate [9]. However, some uncertainties remain: though the pion-coupling is deduced from the nucleon-nucleon case, the $DD^*\pi$ form factor is not known accurately as well as the short-range part of the interaction needed to supplement pion exchange. Also, due the mass difference between D and D^*, the Yukawa potential might be of shorter range in $D\bar{D}^*$ than in the nucleon-nucleon case, and hence be less effective [10].

More generally, several approaches are based on the X(3872) having mainly a $(cq\bar{c}\,\bar{q})$ quark content, where q denotes a light quark. Besides the nuclear-physics approach, schematically noted $(c\bar{q}) - (\bar{c}q)$, an interesting comeback of the diquark concept has been observed in recent literature. In particular, Maiani *et al.* proposed to describe simultaneously the X(3872) and X(3940) as $(cq)(\bar{c}\,\bar{q})$ states and Y(4260) as a $(cs) - (\bar{c}\,\bar{s})$ state with an orbital momentum $\ell = 1$ between the diquark and the anti-diquark [11]. This is a rather elegant picture, but the mass of the diquark is not known and has to be adjusted empirically.

None of the available models has won an overall consensus yet, and the door remains open for another binding mechanism. This is the aim of the present paper. More details, as well as applications to other spin-flavor combinations will be presented in a forthcoming article. The starting point is the chromomagnetic interaction, inspired by the one-gluon-exchange contribution [12], but covering a wider class of model with a spin-spin interaction that bears the color dependence of a color-octet exchange. The chromomagnetic interaction gives a convincing explanation of the mass splittings of ordinary hadrons and has been decisive in promoting the possibility of hadron states with a multiquark content. In particular, some *S*-wave $(q^2\bar{q}^2)$ states can well be lighter than the *P*-wave excitations of the $(q\bar{q})$ system, to explain why supernumerary scalar states are observed with a low mass [13]. Exotic configurations can also occur, due to a coherent chromomagnetic attraction that is larger than the sum of chromomagnetic effects in the decay products [14].

In pioneering papers on chromomagnetic effects applied to multiquark states, ordinary $(q = u, d)$ and strange (s) quarks were treated in the limit of $\mathrm{SU}(3)_\mathrm{F}$ flavor symmetry. However, when its breaking is introduced, the chromomagnetic attraction of the Λ baryon is not changed, while that

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of $H = (ssuudd)$ decreases. Hence the stability of the H-dibaryon is weakened by $\mathrm{SU(3)_F}$ breaking. A similar effect is observed for the 1987-vintage pentaquark, $\mathrm{P} = (\bar{Q}sqqq)$. See, e.g., Ref. [15]. Another difficulty is that the strength of the chromomagnetic force, related to the quark-quark short-range correlation, is probably smaller in H or P than in ordinary hadrons. It thus seems necessary to refine the treatment of chromomagnetic effects.

The present study takes full account of flavor-symmetry breaking when estimating the chromomagnetic interaction of multiquarks, and it happens that this treatment provides a very good candidate for the X(3872), with about the right mass, and the right coupling patterns, namely $\mathrm{D}\bar{\mathrm{D}}^*$ and J/ψ plus a light vector meson.

The interaction Hamiltonian acting on the color and spin degrees of freedom reads

$$H_{\mathrm{CM}} = -\sum_{i,j} C_{ij}\tilde{\lambda}_i \cdot \tilde{\lambda}_j \boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j, \tag{1}$$

where the coefficients C_{ij} depend on the quark masses and properties of the spatial wave function. In absence of a complete theory, this Hamiltonian leads to a mass formula

$$\mathcal{M} = \sum_i m_i - \left\langle \sum_{i,j} C_{ij}\tilde{\lambda}_i \cdot \tilde{\lambda}_j \boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j \right\rangle, \tag{2}$$

with effective masses m_i which include constituent masses and their chromoelectric energy (binding effect). This formula reflects the basic symmetry principles which govern the ground-state hadron masses. The solution of the eigenvalue problem for the chromomagnetic term is thus of interest, not only in spectroscopy, but in all the reactions where a quark or an antiquark interacts with a system of other quarks, for instance, final-state interaction in weak decays.

A useful phenomenology can be developed on the basis of mass formulas such as (2). See, for instance, Ref. [16]. The mesons being more tightly bound than baryons, the fits usually lead to lighter values of the effective masses m_i for mesons, and larger correlation coefficients. This is in qualitative agreement with model calculations which can be performed within the harmonic oscillator model, or with more general inequalities relating mesons to baryons [17]. A fit within ± 10 MeV of charmed baryons gives the set of masses

$$m_c = 1550\ \mathrm{MeV}, \qquad m_q = 450\ \mathrm{MeV}, \qquad m_s = 590\ \mathrm{MeV}, \tag{3}$$

and strength factors

$$C_{qq} = 20\ \mathrm{MeV}, \quad C_{qc} = 5\ \mathrm{MeV}, \quad C_{qs} = 15\ \mathrm{MeV}, \quad C_{ss} = 10\ \mathrm{MeV}, \quad C_{cs} = 4\ \mathrm{MeV}, \quad C_{c\bar{c}} = 4\ \mathrm{MeV}. \tag{4}$$

For $(q\bar{q})$ mesons, $H_{\mathrm{CM}} = 16C_{12}\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2/3$, and for $(q_1q_2q_3)$ baryons with three valence quarks, $\langle \tilde{\lambda}_i \cdot \tilde{\lambda}_j \rangle = -8/3$ factors out for all pairs. Estimation of the value of $\sum \tilde{\lambda}_i \cdot \tilde{\lambda}_j \boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j$ for more complicated systems, once an overall strength has been factored out (i.e., in the flavor symmetric limit), has been carried out in [18] and further developed by several authors. The formulas involve the Casimir operators of SU(2) (spin), SU(3) (color or flavor), and SU(6) (spin-color). As we are dealing with states combining heavy and light quarks, and even account for $\mathrm{SU(3)_F}$ breaking in the light sector, we cannot assume that all C_{ij} are equal. Hence for any given J^{PC} set of quantum numbers, we list all possible color-spin states and write down explicitly H_{CM} in this basis.

In the case of color-singlet, $J^{PC} = 1^{++}$, a basis can be built with (1,3) and (2,4) subsystems having a well-defined color (superscript 1 for singlet and 8 for octet) and spin (0 or 1 in subscript)

$$\begin{aligned}
\alpha_1 &= (q_1\bar{q}_3)_0^1 \otimes (q_2\bar{q}_4)_1^1, & \alpha_2 &= (q_1\bar{q}_3)_1^1 \otimes (q_2\bar{q}_4)_0^1,\\
\alpha_3 &= (q_1\bar{q}_3)_1^1 \otimes (q_2\bar{q}_4)_1^1, & \alpha_4 &= (q_1\bar{q}_3)_0^8 \otimes (q_2\bar{q}_4)_1^8,\\
\alpha_5 &= (q_1\bar{q}_3)_1^8 \otimes (q_2\bar{q}_4)_0^8, & \alpha_6 &= (q_1\bar{q}_3)_1^8 \otimes (q_2\bar{q}_4)_1^8.
\end{aligned} \tag{5}$$

For $(cq\bar{c}\,\bar{q}) = (1, 2, 3, 4)$ states, the calculation is simplified since $C_{14} = C_{23}$ and $C_{12} = C_{34}$ by charge conjugation symmetry.[1] The Hamiltonian $-H_{\mathrm{CM}}$ acting on the basis (5) is represented by the matrix given in Table I.

It is immediately seen from this matrix that in the case where the chromomagnetic interaction is the same for a quark-quark pair as for the quark-antiquark pair, i.e., $C_{12} = C_{23}$, there is an eigenvector with eigenvalue $-(8C_{12} + 28C_{23} + 2C_{13} + 2C_{24})/3$ for the colormagnetic Hamiltonian, which is a pure color octet $\otimes$ octet, spin $(s = 1) \otimes (s = 1)$ state, $\alpha_6 = (c\bar{c})_1^8 \otimes (q\bar{q})_1^8$. This state therefore cannot freely dissociate into a charmonium state and a light meson.

This eigenstate of the chromomagnetic Hamiltonian can fall freely apart in two mesons carrying charm. However, if the state is rewritten in the (1,4) (2,3) basis, corresponding to charmed mesons, i.e., $(c\bar{q})(q\bar{c})$,

$$\begin{aligned}
\beta_1 &= (q_1\bar{q}_4)_0^1 \otimes (q_2\bar{q}_3)_1^1, & \beta_2 &= (q_1\bar{q}_4)_1^1 \otimes (q_2\bar{q}_3)_0^1,\\
\beta_3 &= (q_1\bar{q}_4)_1^1 \otimes (q_2\bar{q}_3)_1^1, & \beta_4 &= (q_1\bar{q}_4)_0^8 \otimes (q_2\bar{q}_3)_1^8,\\
\beta_5 &= (q_1\bar{q}_4)_1^8 \otimes (q_2\bar{q}_3)_0^8, & \beta_6 &= (q_1\bar{q}_4)_1^8 \otimes (q_2\bar{q}_3)_1^8,
\end{aligned} \tag{6}$$

using the crossing matrix from the basis in Eq. (5) to the basis in Eq. (6),

[1] The first calculation relevant for this case was made by G. Gelmini [19].

TABLE I. Colormagnetic Hamiltonian $-H_{\rm CM}$ in the basis (5).

$16C_{13} - \frac{16}{3}C_{24}$	0	0	0	$\frac{8\sqrt{2}}{3}(C_{23} + C_{12})$	0
0	$-\frac{16}{3}C_{13} + 16C_{24}$	0	$\frac{8\sqrt{2}}{3}(C_{23} + C_{12})$	0	0
0	0	$-\frac{16}{3}(C_{13} + C_{24})$	0	0	$\frac{8\sqrt{2}}{3}(C_{23} - C_{12})$
0	$\frac{8\sqrt{2}}{3}(C_{23} + C_{12})$	0	$\frac{2}{3}C_{24} - 2C_{13}$	$\frac{28}{3}C_{23} - \frac{8}{3}C_{12}$	0
$\frac{8\sqrt{2}}{3}(C_{23} + C_{12})$	0	0	$\frac{28}{3}C_{23} - \frac{8}{3}C_{12}$	$-2C_{24} + \frac{2}{3}C_{13}$	0
0	0	$\frac{8\sqrt{2}}{3}(C_{23} - C_{12})$	0	0	$\frac{2}{3}(4C_{12} + 14C_{23} + C_{13} + C_{24})$

$$\begin{bmatrix} \frac{1}{6} & \frac{1}{6} & \frac{1}{3\sqrt{2}} & \frac{2}{3\sqrt{2}} & \frac{2}{3\sqrt{2}} & \frac{2}{3} \\ \frac{1}{6} & \frac{1}{6} & -\frac{1}{3\sqrt{2}} & \frac{2}{3\sqrt{2}} & \frac{2}{3\sqrt{2}} & -\frac{2}{3} \\ \frac{1}{3\sqrt{2}} & -\frac{1}{3\sqrt{2}} & 0 & \frac{2}{3} & -\frac{2}{3} & 0 \\ \frac{2}{3\sqrt{2}} & \frac{2}{3\sqrt{2}} & \frac{2}{3} & -\frac{1}{6} & -\frac{1}{6} & -\frac{1}{3\sqrt{2}} \\ \frac{2}{3\sqrt{2}} & \frac{2}{3\sqrt{2}} & -\frac{2}{3} & -\frac{1}{6} & -\frac{1}{6} & \frac{1}{3\sqrt{2}} \\ \frac{2}{3} & -\frac{2}{3} & 0 & -\frac{1}{3\sqrt{2}} & \frac{1}{3\sqrt{2}} & 0 \end{bmatrix}, \quad (7)$$

it is immediately realized that there is little octet-octet content for this eigentstate in this crossed basis, and that the coulor singlet-singlet part just corresponds to the charmed mesons D and $\bar{\rm D}^*$ (or c.c.). Because of the lack of phase space, this decay is strongly suppressed.

If the condition $C_{12} = C_{23}$ is relaxed and a different interaction strength is allowed for the quark-quark and quark-antiquark pairs, the interesting eigenvector of the chromomagnetic Hamiltonian is readily seen to acquire a small component on the state α_3, but not on α_1 or on α_2. This means that this eigenstate of $H_{\rm CM}$ will choose to disintegrate into a J/ψ and an ordinary vector meson, just as the X(3872) does. There is no amplitude for dissociation into charmonium and a pseudoscalar meson, at least at the level of the mere quark rearrangement.

Instead of an analytical proof which involves some tedious 6×6 linear algebra, a numerical illustration will be given. If the parameters (4) are adopted and if $C_{\bar{q}c} = C_{qc}$ is further assumed, the eigenstate α_6 receives a chromomagnetic energy -76 MeV. If a value $C_{\bar{q}c} = 6.5$ MeV is adopted instead, an eigenvector $\sum_i a_i \alpha_i$ is obtained, with

$$\{a_i\} = \{0, 0, \epsilon = 0.026, 0, 0, \sqrt{1 - \epsilon^2}\}, \quad (8)$$

i.e., a very small $J/\psi + \rho$ or $J/\psi + \omega$ component, and its eigenvalue is now -90 MeV. If inserted in Eq. (2), it corresponds to a mass $\mathcal{M}(X) = 3910$ MeV with the parameters (3), close to the observed mass 3872 MeV. Several corrections can be anticipated, for instance a coupling to the $\rm D\bar{D}^*$ channel.

It is worth stressing that the above state has *not* the lowest eigenvalue for $H_{\rm CM}$. Another eigenstate exists with a much lower eigenvalue, about -220 MeV. This state, $\sum_i b_i \alpha_i$ with

$$\{b_i\} = \{-0.0026, -0.989, 0, -0.146, -0.021, 0.0\}, \quad (9)$$

is seen to be coupled almost completely to the channel consisting of J/ψ and a light pseudoscalar and therefore is probably very broad and is just a part of the continuum.

The remarkable eigenstate of the chromomagnetic Hamiltonian actually consists of four states, namely, $X_+ = (cu\bar{c}\,\bar{d})$, $X_- = (cd\bar{c}\,\bar{u})$, $Y_1 = (cu\bar{c}\,\bar{u})$, and $Y_2 = (cd\bar{c}\,\bar{d})$. They all receive a contribution from the QCD version of the Pirenne annihilation potential [20] acting on $(c\bar{c})$. In addition, the two neutral states mix through annihilation of the $(u\bar{u})$ and $(d\bar{d})$, color-octet, spin 1, components.[2]

With this mixing and the mass difference between u and d quarks, the isospin zero state $(Y_1 + Y_2)$ and the neutral isospin one state $(Y_1 - Y_2)$ are not anymore eigenstates of the Hamiltonian. The mass matrix governing the physical states is

$$\begin{bmatrix} -a & -a \\ -a & 2(m_d - m_u) - a \end{bmatrix}, \quad (10)$$

where a is the annihilation potential term.

In the one-gluon-exchange model with free gluons in the intermediate state, the strength $C_{q\bar{q}}$ of t-channel exchange and a of s-channel exchange are related by $a = 6C_{q\bar{q}}$. However, perturbation theory with confined gluons suggests $a \simeq C_{q\bar{q}}$ [19]. If a value $a = 15$ MeV is taken for the annihilation term and $m_d - m_u = 3.5$ MeV, the "mostly $I = 1$" state lies 31 MeV above the "mostly $I = 0$" state. The lowest state, mostly $I = 0$, has an amplitude for $J/\psi + \rho$ decay which is about 0.11 times the amplitude for $J/\psi + \omega$ decay: this is roughly what is needed to explain the branching ratio of X(3872) for the two different final states. The observed branching ratios are about equal although phase space strongly favors $J/\psi + \rho$ decay, since only the low-mass tail of the ω is kinematically allowed (see [21] for a more detailed discussion on this point). A further shift of the $I = 0$ and $I = 1$ states is induced by the nuclear forces acting on the long-range $(c\bar{q})(q\bar{c})$ part of the wave function, favoring $I = 0$. The effect is there, even if this is not the main binding mechanism in our approach. The state with mostly $I = 1$ isospin content, should be seen as a broad resonance decaying into $\rm D\bar{D}^*$ or $J/\psi + \pi\pi$.

It is natural to ask what happens if the flavor content of $(cq\bar{c}\,\bar{q})$ is modified, while keeping the $J^{PC} = 1^{++}$ quantum numbers. For $(qq\bar{q}\,\bar{q})$ and $(sq\bar{s}\,\bar{q})$, the state is well above the threshold of two mesons. This appears also to be

[2]Mixing with glueballs, hybrids and high-mass tetraquark states is neglected.

the case for the $(cs\bar{c}\,\bar{s})$, $(bs\bar{b}\,\bar{s})$, and $(bc\bar{b}\,\bar{c})$ configurations. (For this last configuration, since the spin excitation of the B_c meson is not known, it has been necessary to extract the value of the coefficient $C_{b\bar{c}}$ from theoretical calculations [22].) The situation is different for the $(bq\bar{b}\,\bar{q})$ states: if the parameters are tuned to fit the measured values of the masses of B, B*, Υ, and Λ_b hadrons, the $(bq\bar{b}\,\bar{q})$ states appears as stable against dissociation into $B\bar{B}^*$. It can decay into $\Upsilon + \omega$. To summarize, the chromomagnetic interaction, acting on the configurations $(cq\bar{c}\,\bar{q})$ with hidden charm, has been shown to single out a remarkable state which is an almost pure octet-octet state in the $(c\bar{c}) + (q\bar{q})$ channel. It has a large singlet-singlet component of the type $D + \bar{D}^*$ in the crossed $(c\bar{q}) + (\bar{c}q)$ channel. However, this decay is kinematically strongly suppressed, as the state is at about the same mass as this threshold. A small impurity gives a small branching ratio into $J/\psi + \rho$ and $J/\psi + \omega$, the former being favored by phase space, while $J/\psi + \pi$ is suppressed. This hadron is thus rather narrow, a remarkable property for a multiquark without internal orbital momentum between clusters. This state is therefore a most natural candidate for describing the X(3872).

Since the time of baryonium, "color chemists" thought that color will show up as a new spectroscopic degree of freedom, and states such as "mock-baryonium," "mesobaryons," or "pseudomesonium" were proposed, with color-triplets, sextets, or octets at both ends of a rotating colorelectric string [23]. However, it was never convincingly explained how such a clustering could occur from the dynamics of confinement. Our state should be more easily accepted, since the two quarks and the two antiquarks are in an overall S-wave.

Further measurements of the properties of the X(3872) will help to test the chromomagnetic mechanism, which furthermore predicts other interesting states, especially in configurations combining heavy and light flavors. An example is $(bc\bar{q}\,\bar{q})$ with $J^P = 1^+$. This will be studied in a forthcoming paper. It is simply stressed here that the mechanism proposed for the X(3872) requires very specific spin and flavor configurations. This explains why multiquark states are so elusive in the hadron spectrum.

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