$D^{(*)}B^{(*)}$ states in heavy meson effective theory

Luciano M Abreu
E-mail: luciano.abreu@ufba.br
Instituto de Física, Universidade Federal do Brazil, 40210-340, Salvador, BA, Brazil

Abstract.
In this work we investigate deuteron-like molecules with both open charm and bottom, within the framework of the Heavy Meson Effective Theory. Using the available information of some exotic hadronic states as inputs to fix the coupling constants of the theory, we analyze the formation of loosely-bound $D^{(*)}B^{(*)}$-states and estimate their relevant properties.

1. Introduction
In recent years, several experiments have been reported the existence of unconventional states in hadron spectroscopy with unusual properties, as unexpected decay modes (see Refs. [1, 2] for a review).

In order to understand the structure of these exotic states, several descriptions have been proposed [1]. Among them, the loosely bound molecule interpretation is one of the most widely used, mainly because the proximity of the masses to some hadronic thresholds. In this sense, these exotic hadrons are interpreted as bound states of heavy mesons if they are below the threshold and in the first Riemann sheet of the scattering amplitude [3, 4, 5]. Some emblematic examples are the $X(3872)$ in charmonium sector, suggested to be a loosely bound state of $D\bar{D}^{*}$ [6, 7, 8, 9, 10, 11]; and the $Z_b(10610)$ in bottomonium sector, considered as a $BB^{*}$ molecule [10, 12, 13].

Notice that there is not yet any experimental evidence of unconventional states in the intermediate region between charmonium and bottomonium sectors. Nevertheless, questions about the existence and properties of exotic hadrons in the $B_c$ sector have been raised. Attempts to address this issue in the context of loosely bound molecule interpretation can be found in Refs. [10, 14, 15, 16, 17].

Thus, taking the scenario above as motivation, in this work we analyze the masses and other properties of loosely-bound $D^{(*)}B^{(*)}$-states within the framework of the Heavy Meson Effective Theory (HMET). We use the information available of observed exotic hadronic states as inputs to fix the coupling constants of the theory.

We introduce the formalism in Section 2. In Section 3 we briefly describe the determination of the coupling constants and present the obtained results. Finally, Section 4 is devoted to discussion and concluding remarks.

2. Formalism
Let us consider the HMET in the analysis of bound states between heavy mesons [7, 8, 17, 18, 19]. At lowest order in the $1/m_Q$ expansion ($m_Q$ is the heavy quark mass), the four-body interaction
Lagrangian respecting heavy-quark spin, heavy-quark flavor and light-quark flavor symmetries can be written as

\[ L_4 = -\frac{D_1}{4} \text{Tr} \left[ H^{(Q)a} H_a^{(Q)} \gamma^\mu \right] \text{Tr} \left[ H^{(\bar{Q})a} \bar{H}_a^{(\bar{Q})} \gamma_\mu \right] + \frac{D_2}{4} \text{Tr} \left[ H^{(Q)a} H_a^{(Q)} \gamma^{\mu\nu} \gamma_5 \right] \text{Tr} \left[ H^{(\bar{Q})a} \bar{H}_a^{(\bar{Q})} \gamma_\mu \gamma_5 \right] - \frac{E_1}{4} \text{Tr} \left[ H^{(Q)a} (\lambda^A)^b H_b^{(Q)} \gamma_\mu \right] \text{Tr} \left[ H^{(\bar{Q})a} (\lambda^A)^b \bar{H}_b^{(\bar{Q})} \gamma_\mu \right] - \frac{E_2}{4} \text{Tr} \left[ H^{(Q)a} (\lambda^A)^b H_b^{(Q)} \gamma_\mu \gamma_5 \right] \text{Tr} \left[ H^{(\bar{Q})a} (\lambda^A)^b \bar{H}_b^{(\bar{Q})} \gamma_\mu \gamma_5 \right]. \tag{1} \]

In Eq. (1), \( Q = c, b \) represents the index with respect to the heavy-quark flavor group \( SU(2)_{HF} \), and \( a \) the triplet index of the light-quark flavor group \( SU(3)_V \). The fields \( H_a^{(Q)} \) and \( H^{(\bar{Q})a} \) are given by

\[ H_a^{(Q)} = \left( 1 + \frac{v_\mu \gamma^\mu}{2} \right) \left( P_{a\mu}^{(Q)} \gamma^\mu - P_{a}^{(Q)} \gamma_5 \right), \]

\[ H^{(\bar{Q})a} = \left( P_{a\mu}^{(\bar{Q})} \gamma^\mu - P_{\bar{a}}^{(\bar{Q})} \gamma_5 \right) \left( \frac{1 - \frac{v_\mu \gamma^\mu}{2}}{2} \right), \tag{2} \]

where \( P_{a\mu}^{(Q/\bar{Q})} \) and \( P_{a\mu}^{(Q/\bar{Q})} \) are the pseudoscalar and vector heavy-meson fields forming a \( \bar{3} \) representation of \( SU(3)_V \):

\[ P_{a}^{(c)} = (D^0, D^+, D_s^+) \],

\[ P_{a}^{(c)} = (D^0, D^-, D_s^-), \tag{3} \]

for the charmed meson field, and

\[ P_{a}^{(b)} = (B^-, \bar{B}^0, \bar{B}_s^0), \]

\[ P_{a}^{(b)} = (B^+, B^0, B_s^0), \tag{4} \]

for the bottomed meson field (and similar expressions for the vector case). The fields \( \bar{H}^{(Q)a} \) and \( \bar{H}_a^{(Q)} \) are the hermitian conjugate fields:

\[ \bar{H}^{(Q)a} = \gamma^\gamma H_a^{(Q)} \gamma^0, \]

\[ \bar{H}_a^{(Q)} = \gamma^0 H_a^{(Q)} \gamma^\gamma. \tag{5} \]

They are necessary to construct invariant quantities under the symmetries above mentioned. Finally, notice that \( \lambda_A \) in Eq. (1) are the Gell-Mann matrices and \( D_1, D_2, E_1 \) and \( E_2 \) are the four parameters describing the interaction strength.

Since we will work at the leading order in the \( 1/m_Q \) expansion, relativistic effects are suppressed, and we can make use of the non-relativistic version of the theory by specifying the velocity parameter as \( v = \{1, 0\} \) [17, 18, 19]. This choice allows us to work only with the euclidian part of the vector meson fields, due to the conditions \( v \cdot P_{a}^{(Q)} = 0 \), and \( v \cdot P_{a}^{(Q)a} = 0 \). Also, the following normalization is employed:

\[ \sqrt{2} P_{a}^{(s\mu)} \to P_{a}^{(s\mu)}. \tag{6} \]
It is interesting to notice that the Lagrangian density $L_4$ in Eq. (1) is leading order $O(1)$ in the $1/m_Q$ expansion as well as in the chiral expansion. Corrections to lowest order appear from higher derivatives or mass terms and from loop diagrams [18, 21]. In addition, pion-exchange contributions are perturbative over the expected range of applicability of HMET and are suppressed, as discussed in Refs. [8, 10, 17, 19]. Therefore, we consider the lowest-order potential only with contact interactions and the region where the pion-exchange contributions are irrelevant.

At this point we are able to obtain the non-relativistic interaction potential, $V$, relating it to the scattering amplitude $iM(D^{(*)}B^{(*)} \rightarrow D^{(*)}B^{(*)})$ via the Breit approximation:

$$ V(p) = -\frac{1}{\sqrt{12m_i 2m_f}} M(D^{(*)}B^{(*)} \rightarrow D^{(*)}B^{(*)}), \quad (7) $$

where $m_i$ and $m_f$ are respectively the masses of initial and final states, and $p$ is the momentum exchanged between the particles in Center-of-Mass.

For our purposes, it is more appropriate to arrange the $D^{(*)}B^{(*)}$-states in the basis of states $B \equiv \{|DB\rangle, |D^*B\rangle, |DB^*\rangle, |D^*B^*\rangle\}$. Then, using the four-body Lagrangian $L_4$ in Eq. (1) and the relation (7), we get the effective potential $V$ at tree-level approximation in the basis $B$:

$$ V(p) = \begin{pmatrix} C_1 & 0 & 0 & -C_2 \varepsilon_1 \cdot \varepsilon_2 \\ 0 & C_1 \varepsilon_3^3 \cdot \varepsilon_1 & -C_2 \varepsilon_3^3 \cdot \varepsilon_2 & -C_2 \varepsilon_2 \cdot \tilde{S}_1 \\ 0 & -C_2 \varepsilon_3^3 \cdot \varepsilon_1 & C_1 \varepsilon_3^3 \cdot \varepsilon_2 & C_2 \varepsilon_1 \cdot \tilde{S}_2 \\ -C_2 \varepsilon_3^3 \cdot \varepsilon_4 & -C_2 \varepsilon_4^3 \cdot \tilde{S}_1 & C_2 \varepsilon_3^3 \cdot \tilde{S}_2 & C_1 \varepsilon_3^3 \cdot \varepsilon_1 \varepsilon_4 \cdot \tilde{S}_2 + C_2 \tilde{S}_1 \cdot \tilde{S}_2 \end{pmatrix}, \quad (8) $$

where $\varepsilon_i$ means the polarization of incoming or outgoing vector heavy meson; $\tilde{S}_i$ is the spin-1 operator, whose matrix elements are equivalent to the vector product of polarizations: $\tilde{S}_1 \equiv (\varepsilon_3^3 \times \varepsilon_1)$ and $\tilde{S}_2 \equiv (\varepsilon_4^3 \times \varepsilon_2)$, and $C_i = D_i + E_i \lambda A^i$; $i = 1, 2$.

Notice that each state of the basis $B$ has nine light-quark flavor states of $SU(3)_V$ basis, one singlet and one octet [14, 17]. Specifically, there are two isosinglets, one isotriplet, and two isodoublets: $\xi \equiv \{s, 2, t, d, 12\}$. In this sense, we have

$$ |H\bar{H}_{\xi}\rangle \equiv \{|H\bar{H}_{s1}\rangle, |H\bar{H}_{s2}\rangle, |H\bar{H}_{t1}\rangle, |H\bar{H}_{d1}\rangle, |H\bar{H}_{d2}\rangle\} \quad (9) $$

where $H\bar{H} = DB, D^*B, DB^*, D^*B^*$. Thus, the coupling constants $C_1$ and $C_2$ appearing in Eq. (8) are given by the following expressions with respect to light-quark flavor basis:

$$ |H\bar{H}_{s1}\rangle : \quad C_i = 2D_i + \frac{20}{3} E_i, $$
$$ |H\bar{H}_{s2}\rangle : \quad C_i = 2D_i + \frac{8}{3} E_i, $$
$$ |H\bar{H}_{t}\rangle : \quad C_i = 2D_i - \frac{4}{3} E_i, $$
$$ |H\bar{H}_{d1}\rangle : \quad C_i = 2D_i - \frac{4}{3} E_i, $$
$$ |H\bar{H}_{d2}\rangle : \quad C_i = 2D_i - \frac{4}{3} E_i. \quad (10) $$

We are interested in dynamically generated poles in the amplitudes. So, we make use of solutions of Lippmann-Schwinger equation for a specific channel,

$$ T^{(\alpha)} = \frac{V^{(\alpha)}(\omega)}{1 - V^{(\alpha)}(\omega) G^{(\alpha)}}, \quad (11) $$

where $T^{(\alpha)}$ are the transition amplitudes, $\alpha = |H\bar{H}\xi\rangle$, and

$$
G \equiv \int \frac{d^4q}{(2\pi)^4} \frac{1}{2m_{D(s)}^2 + q_0^2 - \tilde{q}^2} + \frac{i\epsilon}{2m_{\pi(s)}^2 + q_0 - \tilde{q}^2} + i\epsilon.
$$

(12)

In the case of vector mesons, we must replace $G \rightarrow G^{\mu\nu}$.

Bound-state solutions, i.e. poles located below the threshold (in the first Riemann sheet), can be obtained by applying in Eq. (11) the residue theorem and dimensional regularization [7, 17]. Therefore, some relevant physical quantities can be obtained, as the binding energy and scattering length, respectively:

$$
E_b^{(\alpha)} = \frac{32\pi^2}{\left(\tilde{V}^{(\alpha)}\right)^2 \mu^3},
$$

$$
a_s^{(\alpha)} = \frac{\mu \tilde{V}^{(\alpha)}}{8\pi},
$$

(13)

where $\tilde{V}^{(\alpha)}$ is the renormalized potential, i.e. the renormalized contact interaction, and $\mu$ is reduced mass of $D^{(s)}B^{(s)}$ system.

We remark that the potential $\tilde{V}^{(\alpha)}$ is a quantity dependent of renormalization scheme, because possible divergences contained in the Lippmann-Schwinger equation may be absorbed in renormalization of the coupling constants. However, physical quantities such as binding energy and scattering length are observables, so they are renormalization-independent.

3. Results

Now we discuss the bound-state solutions for $D^{(s)}B^{(s)}$ systems in the framework introduced above. As remarked, we believe that this analysis taking into account the contact interaction makes sense in the region where pion-exchange contributions are irrelevant. Then, an appropriate constraint is to consider the region where bound states obey the requirement

$$
a_S \gtrsim 3\lambda_\pi,
$$

(14)

where $\lambda_\pi = 1/m_\pi \sim 1$ fm is the pion Compton wavelength.

We need a strategy for fixing the four coupling constants introduced previously. The point is that there is no yet experimental information of exotic states in $B_c$ sector. Nevertheless, there are available experimental data in charmonium and bottomonium sectors, and although renormalization procedure should be performed in each sector, a first attempt can be the analysis of the parameter space in the light of information available. Thus, we choose a specific renormalization scheme, which allows us to relate the results for different sectors, and determine the mass, binding energy and scattering length of $S$-wave bound states as functions of interaction strengths.

Suitably, we use the information available for channels that have a different dependence on the coupling constants. In agreement with Refs. [9, 10, 17], we employ the data about $X(3872)$, $Z_b(10610)$, $X(3915)$ and $Y(4140)$ states as inputs to fix the four couplings. Specifically,

(i) $X(3872)$ state: interpreted here as a $S$-wave isoscalar ($D^0\bar{D}^{*0} + c.c.$)-molecule ($J^{PC} = 1^{++}$) [24]. This channel has the leading-order potential given by

$$
V(D^0\bar{D}^{*0} + c.c.)(\beta S_1) = (D_1 + 3E_1) + (D_2 + 3E_2).
$$

(15)

We consider the following values of $X(3872)$ mass, threshold and binding energy: $M_X = 3871.69$ MeV [2], $m_{D^0} + m_{\bar{D}^{*0}} = 3871.8$ MeV [2], and $E_b = 0.11$ MeV [24], respectively.
Z_b(10610): understood as S-wave isovector \((B\bar{B}^* + c.c.)\) molecule \((J^{PC} = 1^{+-})\), with the leading-order potential being given by

\[
V^{(B\bar{B}^*+c.c.)}(3S_1) = (D_1 - E_1) - (D_2 - E_2).
\] (16)

The binding energy of \(Z_b(10610)\) is assumed to be \(2.0 \pm 2.0\) MeV [10, 17], the mass and threshold are \(M_{Z_b} = 10602.6 \pm 2.0\) MeV and \(m_B + m_{\bar{B}} = 10604.6\) MeV [2], respectively.

(iii) \(X(3915)\) state: it is interpreted as a \(0^{++}\) isoscalar \((D^*\bar{D}^*)\) molecule, with

\[
V^{(D^*\bar{D}^*)}(1S_0) = (D_1 + 3E_1) - 2(D_2 + 3E_2)
\] (17)

The masses and threshold are \(M_{X(3915)} = 3917\) MeV, \(m_{D^*} + m_{\bar{D}^*} = 4017.2\) MeV.

(iv) \(Y(4140)\) state: it is a \(0^{++} (D_s^*\bar{D}_s)\) molecule, with the leading order potential given by:

\[
V^{(D_s^*\bar{D}_s)}(1S_0) = (D_1 + E_1) - 2(D_2 + E_2).
\] (18)

The masses and threshold are \(M_{Y(4140)} = 4140\) MeV and \(m_{D^*} + m_{\bar{D}^*} = 4224.6\) MeV.

Note that contact interactions in Eqs. (15), (16), (17) and (18) are renormalized. However, as a means to simplify the notation we continue to denote henceforth the renormalized parameters as \(D_1, D_2, E_1\) and \(E_2\).

Therefore, using Eqs. (15)-(18) in Eq. (13) to reproduce the data available of the \(X(3872), Z_b(10610), X(3915)\) and \(Y(4140)\) states, we get

\[
\begin{align*}
D_1 &= -3.97 \times 10^{-4}\ \text{MeV}^{-2}, \\
E_1 &= 2.70 \times 10^{-4}\ \text{MeV}^{-2}, \\
D_2 &= -1.70 \times 10^{-4}\ \text{MeV}^{-2}, \\
E_2 &= -1.35 \times 10^{-4}\ \text{MeV}^{-2}.
\end{align*}
\] (19)

Hence, the quantities characterizing the \(D(s)^*B(s)^*\) states can be estimated in the specific renormalization scheme explicit in Eq. (19). The results are shown in Table 1.

Table 1: Quantities obtained for the \(D(s)^*B(s)^*\) systems, with respect to the \(B\) and light-quark flavor \(SU(3)_V\) bases (see Eq. (9)). The couplings \((D_1, E_1, D_2, E_2)\) are chosen as in the situation manifested in Eq. (19); \((m_H + m_R)\), \(M\), B.E., and \(a_s\) mean threshold, mass, binding energy and scattering length of each respective state.
It can be noticed that the errors in the quantities in Table 1 are computed taking into account the partial errors: in $Z_b(10610)$ binding energy, and due to violations of heavy-quark spin and light-quark flavor symmetries. In the case of violation of heavy-quark spin symmetry, we expect a relative uncertainty of the order of $\Lambda_{QCD}$. Hence, the total errors in $\tilde{\alpha}$ by considering the ratio between the kaon and pion decay constants, $f_K/f_\pi$, which yields a relative error of 20% in states containing strange quarks. Assuming $\Lambda_{QCD} \sim 200$ MeV and the quark charm mass $m_c \sim 1.5$ GeV [9], this gives 15% error in leading-order contact interactions. Deviations due to $SU(3)_V$-breaking effects are evaluated by considering the ratio between the kaon and pion decay constants, $f_K/f_\pi \sim 1.2$, which yields a relative error of 20% in states containing strange quarks. Hence, the total errors in $\tilde{\alpha}$ and relevant quantities are computed by adding the partial errors in quadrature.

In general, the outcomes shown in Table 1 reveal a pattern of loosely bound states. It is worthy mentioning, however, that the choice of $X(3915)$ and $Y(4140)$ states as inputs to fix our renormalization scheme present some shortcomings, as larger binding energies and experimental outcomes, as remarked in Ref. [9].

4. Discussion and Concluding Remarks
In this work we have analyzed the formation of mesonic molecules with both open charm and bottom in the framework of Heavy-Meson Effective Theory. The parameter space of the coupling constants has been explored in the specific region in which loosely bound $D^{(*)}B^{(*)}$-state solutions are allowed.

We have chosen a specific renormalization scheme, in which the relation between the bare coupling constants is conserved after renormalization procedure. This approach has allowed to relate the results for different sectors, taking into account the fact that there is not yet experimental evidence of hadronic molecules with both open charm and open bottom.

The couplings have been fixed by considering the available experimental data in charmonium and bottomonium sectors. In particular, we have used the location of $X(3872)$, $Z_b(10610)$, $X(3915)$ and $Y(4140)$ states, interpreted in our context as heavy-meson molecules. Then, the masses, binding energies and scattering lengths of $D^{(*)}B^{(*)}$-states have been estimated as
functions of interaction strength in a specific renormalization scheme, examining each state of light-quark flavor $SU(3)_V$ basis. The results obtained with this specific choice exhibit in general a pattern of loosely bound states.

We stress that the present approach is valid in the region of relevance of contact-range interaction, in which the pion-exchange contribution is not relevant. In this sense, an appropriate condition of validity is to consider bound states which obey the requirement in Eq. (14).

The results have indicated that, except for the case $|D^*B^*(1S_0)_{\xi}\rangle$, the $t, d, l_1, d_2$ channels present greater binding energies than the isosinglets $s_1, s_2$ in the specific choice of parameters. Moreover, the mentioned case $|D^*B^*(1S_0)_{\xi}\rangle$ is the only one in this specific choice that does not obey the condition (14), which does not assure the domain of validity of contact-range interactions for this state.

Another point is that the choice given by Eq. (19) generates isosinglets with some properties (very small binding energies, for example) that suggest the impossibility of their formation.

Some results reported above can be compared with other ones available in literature. For example, we see that our findings in Table 1 for the $|D^*B^*(3S_1)_{t}\rangle$ state are in the range of the results in Ref. [10], but in the case of $|D^*B^*(5S_2)_{s}\rangle$ our estimation yields a bound state with much smaller binding energy.

Besides, bound-state solutions identified in Ref. [14] can be discussed in some way under our framework, with the renormalized coupling constants in different renormalization schemes, as remarked in Ref. [17]. Nevertheless, this comparison with Ref. [14] must be done carefully, since we do not consider the S-D mixing effect.

The present approach deserves some improvements, as the inclusion of pion-exchange contributions to increase the range of validity beyond the contact interaction. Also, analyses of hypothetical decays of the $D^*B^*$ states into $B_c$ mesons and light mesons may provide some insight into this sector of exotic hadron spectroscopy.

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