Heavy Quarkonium and non-perturbative corrections

A. Pineda and J. Soto

Departament d’Estructura i Constituents de la Matèria

and

Institut de Física d’Altes Energies

Universitat de Barcelona

Diagonal, 647

E-08028 Barcelona, Catalonia, Spain.

e-mails: pineda@ecm.ub.es, soto@ecm.ub.es

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Abstract

We analyse the possible existence of non-perturbative contributions in heavy \( \bar{Q}Q \) systems (\( \bar{Q} \) and \( Q \) need not have the same flavour) which cannot be expressed in terms of local condensates. Starting from QCD, with well defined approximations and splitting properly the fields into large and small momentum components, we derive an effective lagrangian where hard gluons (in the non-relativistic approximation) have been integrated out. The large momentum contributions (which are dominant) are calculated using Coulomb type states. Besides the usual condensate corrections, we see the possibility of new non-perturbative contributions. We parametrize them in terms of two low momentum correlators with Coulomb bound state energy insertions \( E_n \). We realize that the Heavy Quark Effective lagrangian can be used in these correlators. We calculate the corrections that they give rise to in the decay constant, the bound state energy and the matrix elements of bilinear currents at zero recoil. We study the cut-off dependence of the new contributions and we see that it matches perfectly.
with that of the large momentum contributions. We consider two situations in
detail: i) $E_n \gg \Lambda_{QCD}$ ($M_Q \to \infty$) and ii) $E_n \ll \Lambda_{QCD}$, and briefly discuss
the expected size of the new contributions in $\Upsilon$, $J/\Psi$ and $B_c$ systems.

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1 Introduction

The study of heavy quark bound state systems remains one of the more promising topics in order to test both perturbative and non-perturbative aspects of QCD, as it is clear from the steady activity in the field [1]-[6]. These systems can be understood in a first approximation as non-relativistic bound states which occur due to a Coulomb type interaction predicted by perturbative QCD. In order to improve this basic picture one has to deal on one side with perturbative relativistic and radiative corrections, and on the other side with non-perturbative corrections (power corrections).

In this paper we shall only be concerned with non-perturbative corrections. Usually, the latter have been parametrized using both the multipole expansion and the adiabatic approximation in terms of the gluon condensate [7, 8]. Corrections to the Coulomb potential due to condensates can also be considered, although these are sub-leading [3, 9]. We have argued before [6] that new non-perturbative contributions could arise which cannot be expressed in terms of local condensates, and hence a convenient parametrization for them is required. This kind of nonperturbative contributions has been discussed in [10] in a different context and, in fact, the various Isgur-Wise functions extensively used in the Heavy Quark Effective Theory (HQET) may be regarded as such [11].

Let us recall the main idea behind the possibility of new non-perturbative contributions in heavy quarkonium\footnote{We use 'heavy quarkonium' to denote a general heavy quark-antiquark bound state. The quark and the antiquark need not have the same flavour.}. When the relative three momentum in the bound state is big enough the dominant interaction must be the perturbative Coulomb potential, but for small relative three momentum this need not longer be true. Therefore, heavy quarks in the latter kinematical situation should better be kept as low energy degrees of freedom. It turns out that a convenient parametrization of this kinematical region may be given in terms of the HQET for quarks and antiquarks [6, 12].

The HQET for quarks and antiquarks enjoys rather peculiar features, which make it quite different from the usual HQET describing either quarks or antiquarks, which has been so popular in the study of $Q\bar{q}$ and $Qqq$ systems in recent years [13] (see [14] for reviews). For instance, it enjoys a symmetry, which is larger than the well-known spin and flavor symmetry, that breaks spontaneously down to the latter giving rise to...
quark-antiquark states as Goldstone modes [12]. Its peculiarities concerning radiative corrections have recently been illustrated in [15].

The main aim of this paper is to work out a controlled derivation from QCD of the effective lagrangian describing the small relative momentum regime of heavy quarks in quarkonium. Whereas the basic ideas above have already been elaborated in [6], a complete and systematic derivation is lacking, and hence worth being presented. Within this new framework we recalculate the non-perturbative contributions of this region to the energy levels, the decay constant and the matrix elements of bilinear currents at zero recoil. We find a few corrections to the formulas given in [6]. For all these observables it is enough to work in the center of mass frame (CM), which we shall do in most of the paper.

In order to deal with heavy quarkonia systems we keep the relevant degrees of freedom in the QCD lagrangian. In fact, since virtual heavy quark creation is very much suppressed, we could safely start from non-relativistic QCD (NRQCD). The derivation of NRQCD from QCD is well understood and a technique to incorporate relativistic corrections to it has also been developed [16]. First of all, we split the gluon field in hard and soft by a three momentum cut-off. From the hard gluon fields we only keep the zero component and disregard the spacial components. This is legitimate as far as we are not interested in relativistic corrections. We next integrate out the zero component of the hard gluon field to obtain the Coulomb potential between heavy quark currents. The Coulomb potential has an infrared momentum cut-off since the zero component of the soft gluon field has not been integrated out. At this point we have an effective lagrangian formally equal to the one used by Voloshin and Leutwyler (VL) [7, 8], except for the IR cut-off in the Coulomb potential. After introducing CM and relative momentum for the bound states we are interested in, we further split the quark fields in large and small relative three-momentum regimes\(^2\). The resulting lagrangian can then be separated in three pieces: \(L^\mu\) which contains small relative momentum quark fields only, \(L_\mu + L^I_\mu\) contains large relative momentum quark fields only and \(L^{I\mu}_\mu\) which contains both small and large relative momentum quark fields. For \(L^\mu\) we can approximate the lagrangian to the HQET lagrangian, where eventually all its powerful symmetries can be used. No Coulomb term remains in this part of the lagrangian. For \(L_\mu + L^I_\mu\) we obtain again the VL starting point lagrangian except for two facts: both

\(^2\)The large and small relative momentum regions were denoted as off- and on-shell regions respectively in [6].
the Coulomb potential and the Hilbert space are restricted to three-momenta larger than a certain cut-off. Keeping the cut-off much higher than $\Lambda_{QCD}$ but much smaller than the inverse Bohr radius we may safely assume that the multipole expansion holds for this part of the lagrangian. If we further assume that the adiabatic approximation also holds, we may proceed in total analogy to VL. The hypothesis above on the cut-off also allows us to treat $L^I_\mu$ as a perturbation. The various contributions from this perturbation to the different observables can be eventually expressed as correlators of the HQET.

We would like to stress that our formalism is less restrictive than the one used by VL since neither the adiabatic approximation nor the multipole expansion are assumed to hold in the small relative momentum region of the heavy quark fields. Indeed we may always recover VLs results by putting to zero the cut-off which separates large and small relative momentum.

We distribute the paper as follows. In sect. 2 we derive the effective action for the small relative momentum fields. In sect. 3 we calculate the decay constant, the bound state mass and the matrix element of any bilinear heavy quark current between quarkonia states at zero recoil. The latter is relevant for the study of semileptonic decays at zero recoil. In sect. 4 we prove the cut-off independence of our results. In section 5 we study the low momentum correlators in two situations: the asymptotic limit ($M_Q \to \infty$) $E_n \gg \Lambda_{QCD}$, where, using OPE techniques, we see that no new corrections arise, and ii) $E_n \ll \Lambda_{QCD}$, where the low momentum contributions are evaluated using an effective 'chiral' lagrangian which incorporates the relevant symmetries of the HQET for quarks and antiquarks. Working in this way we find new non-perturbative contributions which are parametrized by a single non-perturbative constant. We also give preliminary estimations of their size. Section 6 is devoted to the conclusions.

2 Effective action

In this section we derive the effective lagrangian for heavy quarks and antiquarks in the small relative momentum regime from QCD.

The QCD lagrangian reads

$$\mathcal{L} = -\frac{1}{4} F^2 + \sum_a \bar{Q}_a(i\not{D} - m_a)Q_a$$

(2.1)
where
\[ D_\mu = \partial_\mu - igV_\mu, \quad V = V^{r} T^{r} \]  \hspace{1cm} (2.2)
\[ F_{\mu\nu}^{r} = \partial_\mu V_\nu^{r} - \partial_\nu V_\mu^{r} + gf^{rst}V_\mu^{s}V_\nu^{t} \] \hspace{1cm} (2.3)

We split the gluon field \( V \) in large \( A \) and small \( B \) momentum modes \( V(x) = A(x) + B(x) \). Next we exactly integrate \( A_0 \) and neglect \( A_1 \). The latter would give rise to relativistic corrections. Consistently, at the same point we perform a Foldy-Wouthuysen transformation and keep terms up to \( 1/m \). We obtain
\[
L = \frac{1}{4} \int d^3xF_B + \sum_a \int d^3x \left( \tilde{Q}_a(i\gamma^0 D_0^B - m_a)Q_a + \tilde{Q}_a\tilde{D}_B^{2m_a}Q_a + \tilde{Q}_a\frac{gS_{B}Q_a}{2m_a} \right) \\
+ O\left( \frac{1}{m_a^2} \right) - \frac{g^2}{2} \sum_{aa'} \int d^3x \int d^3y \tilde{Q}_a\gamma^0 T^{r} Q_a(x) \left( \frac{1}{D_B^2} \right)^{r_y} \tilde{Q}_a\gamma^0 T^{s} Q_a(y) 
\]  \hspace{1cm} (2.4)

which is manifestly gauge invariant \(^3\). Although, in principle, we could attempt to carry out an explicitly gauge invariant calculation, in practise, it is most convenient to choose a slightly modified Schwinger gauge for the small momentum gluons
\[ \left( \tilde{z} - \frac{m_a}{m_a + m_d} \tilde{y} \right) \tilde{B}(z) = 0 \] \hspace{1cm} (2.5)

In this gauge \( \tilde{B} \) in the kinetic and Coulomb terms gives rise to subdominant contributions when the multipole expansion is carried out, which greatly simplifies the calculation. In particular, recall that the propagator in the Coulomb term always carries large momentum (we have not integrated out the small momentum \( V_0 \) which is kept in \( B_0 \)). Hence the multipole expansion is always legitimated in the Coulomb term. This allows to drop \( \tilde{B} \) in the Coulomb term straight away. As long as we are interested in quark-antiquark bound states only, we may also safely neglect the four-fermion interaction terms involving only quarks or only antiquarks. We next rearrange the quark-antiquark interaction term in a convenient way in order to describe the bound state dynamics . Finally, the effective lagrangian reads
\[
L = \frac{1}{4} \int d^3xF_B + \sum_a \int d^3x \left( \tilde{Q}_a(i\gamma^0 D_0^B - m_a)Q_a + \tilde{Q}_a\tilde{D}_B^{2m_a}Q_a \right) \\
+ O\left( \frac{1}{m_a^2} \right) - \frac{1}{2} \sum_{aa'} \frac{m_a^3 N_A}{(2\pi)^3} \sum_s \int \frac{d^3q}{(2\pi)^3} \int \frac{d^3q'}{(2\pi)^3} V^{A}(q' - q) \\
\times \tilde{Q}_a(-m_a\tilde{v} + \tilde{q}, t)\Gamma_s\tilde{Q}_{a'}(m_{a'}\tilde{v} + \tilde{q'}, t) \left[ \tilde{Q}_{a'}(m_{a'}\tilde{v} + \tilde{q'}, t)\Gamma_s\tilde{Q}_a(-m_a\tilde{v} + \tilde{q'}, t) \right] 
\] \hspace{1cm} (2.6)

\(^3\)Similar approaches can be found in the literature \([4]\).
where $A = 0, r$ denotes colour (0 singlet and $r$ octet, $r = 1...8$), $|\vec{q} - \vec{q}'| > \mu$, $\mu$ being the cut-off which separates small and large momenta, and

$$m_{aa'} = m_a + m_{a'}, \quad \Gamma_s = i\gamma_5 p_-, i\gamma^i p_-,$$

$$p_\pm := \frac{1 \pm \gamma_0}{2}, \quad (2.7)$$

$$N_A = \frac{1}{\sqrt{N_c}}, \sqrt{2}, \quad T^A = 1, T^r.$$  

While, the potential reads

$$V^0(\vec{p}) = -\frac{C_F g^2}{\vec{p}^2}, \quad V^r(\vec{p}) = \frac{g^2}{2N_c}\vec{p}^2. \quad (2.8)$$

where $C_F = (N_c^2 - 1)/2N_c$ and $|\vec{p}| > \mu$ must be understood due to the cut-off coming from soft gluons.\footnote{Several aspects related to this cut-off dependence have been studied in [15].}

Written in this way, we can understand the four-fermion Coulomb interaction term as one which creates a quark-antiquark state with central velocity $\vec{v}$ and relative momentum $\vec{q}'$ and annihilates a quark-antiquark state with the same center of masses velocity $\vec{v}$ and relative momentum $\vec{q}$. Obviously $\vec{v}$ is a conserved quantity in this non-relativistic approximation. We consider the spin breaking term as subleading and we will neglect it in the following. Therefore, spin symmetry for both low and high momentum is implicit in the rest of the paper.

If we stopped at this point we would obtain the standard VL results. However, we would like to go beyond and look for new non-perturbative contributions. We observe that quarks with small relative three momentum only feel the Coulomb interaction of quarks with large relative momentum. This suggests to perform a splitting of the physical quark and antiquark fields into small and large relative momentum in the bound state. The physical picture behind is that if the relative three momentum in the bound state is big enough we can understand it as a perturbative Coulomb type bound state. But for small relative three momentum that is no longer true. For that momentum regime the quark and antiquark fields should be kept as low momentum degrees of freedom. That is, in fact, the main idea of the paper. Therefore, let us write down the currents related to the physical quark-antiquark states in momentum space

$$J^{A,\alpha'}_{\Gamma}(x) = \bar{Q}_{\alpha'} T^A \Gamma Q_{\alpha}(x) =$$

$$m_{aa'}^3 \int \frac{d^3q}{(2\pi)^3} \tilde{\epsilon} m_{aa'} \vec{v}, \frac{d^3q}{(2\pi)^3} \tilde{\epsilon} m_{aa'} (-m_{aa'} \vec{v} - \vec{q}, t) T^A \Gamma \bar{Q}_{\alpha}(m_{aa'} \vec{v} - \vec{q}, t). \quad (2.9)$$
The matrix $\Gamma$ should be such that it projects over quark-antiquark states according to our non-relativistic picture. Notice that the time dependence is kept explicit. Furthermore, we split the relative three momentum with the same cut-off $\mu$ as above. Thus, (2.9) reads

$$
J^A_{\Gamma}(x) = J^A_{l_\Gamma}(x) + J^A_{h_\Gamma}(x)
$$

$$
= m^3_{aal} \int \frac{d^3q}{(2\pi)^3} e^{im_{aal}v_\alpha x} \int_\mu \frac{d^3q}{(2\pi)^3} \bar{h}_a^v(-q_\alpha t)T^A \Gamma h_a^v(-\bar{q}_\alpha t)
+ m^3_{aal} \int \frac{d^3q}{(2\pi)^3} e^{im_{aal}v_\alpha x} \int_\mu \frac{d^3q}{(2\pi)^3} \bar{Q}_a(-m_{aal}v_\alpha - \bar{q}_{\alpha} t)T^A \Gamma Q_a(m_{aal}v_\alpha - \bar{q}_{\alpha} t),
$$

(2.10)

where $\bar{Q}_a(m_{aal}v_\alpha + \bar{k}) = h_a^v(\bar{k})$.

After that we may divide the lagrangian in three pieces.

$$
L = L_\mu + L^\mu + L^I
$$

(2.11)

$L_\mu$ is the piece of the effective lagrangian containing large momenta only. It reads

$$
L_\mu = \sum_a \int d^3x \left( Q_a(ie^0_\alpha \partial_0 - m_a)Q_a + Q_a \bar{\nabla}_2^2 Q_a - \frac{1}{2} \sum_{aa'} m_{aal} N^2_a s \int \frac{d^3q}{(2\pi)^3} \int_\mu \frac{d^3q'}{(2\pi)^3} \int_\mu \frac{d^3q''}{(2\pi)^3} V^A(q'' - q') \times \{ \bar{Q}_a(-m_{aal}v_\alpha - \bar{q}_{\alpha} t)T^A \Gamma_s \bar{Q}_a(m_{aal}v_\alpha + \bar{q}_{\alpha} t) \}ight)

$$

(2.12)

where $|\bar{q} - q'| > \mu$.

In fact it is nothing but the standard Coulomb lagrangian, except for the cut-offs. $L^\mu$ is the piece of the effective lagrangian containing small momenta only. It reads

$$
L^\mu = -\frac{1}{4} F_B^2 + \sum_a \left\{ T^a_v(ie^0_\alpha D_0^B - m_a)h^v_a + \bar{T}_a^v \bar{D}_B^2 h^v_a \right\}.
$$

(2.13)

Notice that (2.13) does not have the four-fermion Coulomb term. It contains the whole soft gluon lagrangian as well as the heavy quark and antiquark fields with small three relative momentum. All the fields in (2.13) are in the non-perturbative regime of QCD. Notice that if we drop the term in $1/m_a$ and make $h^v_a \to e^{-ie^0_\alpha m_a x_0} h^v_a$ (2.13) becomes the HQET lagrangian in the rest frame. Although the $1/m_a$ term is naively subleading for small relative momentum, it plays a crucial role in certain circumstances, as we shall see in Section 4. Nonetheless let us advance that for the correlators we will be interested in one can safely neglect it and work with the HQET lagrangian.
\( L^I \) mixes small and large momenta

\[
L^I = L^I_\mu + L^I_{\mu} \tag{2.14}
\]

The first term reads

\[
\mathcal{L}^I_\mu(x) = gQ_\alpha \gamma^0 B^T_\alpha T^\nu Q_\alpha(x) \tag{2.15}
\]

which gives the leading contribution to the multipole expansion. We will not discuss these contributions (2.15) here since they have been extensively studied in the literature [3, 7, 8]. Let us focus on the second term. It reads

\[
L^I_{\mu} = -\frac{1}{2} \sum_{a \sigma A} m^3_{a \sigma} N^2_A \sum_{s} \int \frac{d^3v}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} V^A(q - q') \times \left[ \tilde{\psi}_\alpha'(\vec{q}', t) T^A \Gamma_s \tilde{\psi}_\alpha'(\vec{q}, t) \right] \left[ \bar{Q}_\alpha(m_\sigma \vec{v} + \vec{q}', t) T^A \bar{Q}_\alpha(-m_\sigma \vec{v} + \vec{q}', t) \right] + (h.c.).
\]  

In this expression the Coulomb potential is the only piece which mixes small and large relative momentum. We can perform a derivative expansion since \( q \) and \( q' \) belong to different momentum regimes \( (q \sim \Lambda_{QCD} << q' \sim \frac{m_\sigma}{\Lambda}) \) and keep only the leading term (further orders would give subleading corrections). It turns out that the small relative momentum term decouples from the Coulomb potential and can be written like a local current. Finally, we obtain

\[
L^I_{\mu} = -\frac{1}{2} \sum_{a \sigma A} m^3_{a \sigma} N^2_A \sum_{s} \int \frac{d^3v}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} V^A(q) \times \int d^3x e^{-im_\sigma \vec{v} \cdot \vec{x}} f_{A, \sigma}^{A, a \sigma} (x, t) \left[ \bar{Q}_\alpha(m_\sigma \vec{v} + \vec{q}', t) T^A \bar{Q}_\alpha(-m_\sigma \vec{v} + \vec{q}', t) \right] + (h.c.).
\]  

The formalism developed in [6] was not powerful enough as to uncover the interaction lagrangian (2.17). This interaction lagrangian is indeed the responsible for the differences between the results presented there and the ones obtained in the next section.

If we assume that small momentum terms are small in comparison with the large momentum terms we can treat the interaction lagrangian (2.17) as a perturbation. This is so for the lower energy levels of heavy quark bound states. In the next sections we focus on the nonperturbative contributions coming from (2.13) and (2.17).
3 Physical Observables

In this section we work out the non-perturbative corrections from the small relative momentum region to the decay constant, the bound state mass and the matrix elements of bilinear currents at zero recoil. We take the bound state velocity small or zero.

Consider first the eigenvalues and eigenstates of $H_\mu$, the hamiltonian associated to $L_\mu$. They read

$$\begin{align*}
| (ab,n,s,A) ; \vec{v} \rangle &= \frac{N_A}{\sqrt{2}} \int \frac{d^3k}{(2\pi)^3} \hat{\Phi}_{ab,n}^A (\vec{k} ; \mu) \vec{v}^\alpha (\vec{p}_1) \Gamma_\epsilon \nu^\beta (\vec{p}_2) T_{i_1,i_2}^{A(\alpha_1\beta_1)} (\vec{p}_1) \delta_{i_1,i_2}^{b_1} (\vec{p}_2) | 0 \rangle, \\
E_{ab,n}^A (\mu) &=
\end{align*}$$

(3.1)

(3.2)

where

$$\begin{align*}
\vec{p}_1 &= m_a \vec{v} + \vec{k} , \quad \vec{p}_2 = m_b \vec{v} - \vec{k} , \\
E_{ab,n}^A , \Phi_{ab,n}^A (\vec{x}) \text{ and } \hat{\Phi}_{ab,n}^A (\vec{k}) \text{ are the energy, the coordinate space wave function and} \\
\text{the momentum space wave function of a Coulomb-type state with quantum number} \\
n = (n,l,m) . \quad \vec{v} \text{ is the bound state 3-vector velocity. } a \text{ and } b \text{ are flavour indices} \text{ and } s \text{ denotes spin. } b^\dagger \text{ and } d^\dagger \text{ are creation operators of particles and anti-particles} \text{ respectively.}
\end{align*}$$

(3.3)

$$\begin{align*}
\{ b^\dagger_{\alpha_1,i_1} (\vec{p}_1) , b^\dagger_{\beta,i_2} (\vec{p}_2) \} &= (2\pi)^3 \delta^{ab} \delta_{\alpha_\beta} \delta_{i_1i_2} \delta^3 (\vec{p}_1 - \vec{p}_2) \\
\{ d^\dagger_{\alpha_1,i_1} (\vec{p}_1) , d^\dagger_{\beta,i_2} (\vec{p}_2) \} &= (2\pi)^3 \delta^{ab} \delta_{\alpha_\beta} \delta_{i_1i_2} \delta^3 (\vec{p}_1 - \vec{p}_2) \\
\{ b^\dagger_{\alpha_1,i_1} (\vec{p}_1) , d^\dagger_{\beta,i_2} (\vec{p}_2) \} &= 0
\end{align*}$$

(3.4)

(3.5)

(3.6)

$u^\alpha (\vec{p}_1)$ and $\nu^\beta (\vec{p}_2)$ are spinors normalized in such a way that in the large $m$ limit the following holds

$$\begin{align*}
\sum_\alpha u^\alpha (\vec{p}_1) \bar{u}^\alpha (\vec{p}_1) = p_+ , \quad \sum_\alpha \nu^\alpha (\vec{p}_1) \bar{\nu}^\alpha (\vec{p}_1) = -p_- .
\end{align*}$$

(3.7)

(3.1) has the non-relativistic normalization

$$\langle (ab,n,s,A) ; \vec{v} | (a'b',n',s',A') ; \vec{v}' \rangle = (2\pi)^3 \delta^{(3)} (m_{ab} (\vec{v} - \vec{v}') ) \delta_{nn'} \delta_{ss'} \delta_{AB} \delta_{AA'} ,$$

(3.8)

where we have used

$$\begin{align*}
\text{tr} \left( p_+ \Gamma^s p_- \Gamma^{s'} \right) &= -2 \delta_{s,s'} .
\end{align*}$$

(3.9)
From (3.8) it follows the wave function normalization

$$\int d^3\bar{q} \tilde{\Phi}^A_{ab,n}(\bar{q}; \mu) \tilde{\Phi}^A_{ab,n}(\bar{q}; \mu) = \delta_{n,n'} \quad (3.10)$$

where there is no sum over A. The wave function and the energy fulfill the equation

$$\frac{p^2}{2\mu_{ab}} \Phi^A_{ab,n}(\vec{p}, \mu) + \int \frac{d^3\bar{q}}{(2\pi)^3} \tilde{\Phi}^A_{ab,n}(\bar{q}; \mu) V^A(\vec{p} - \bar{q}) = E^A_{ab,n}(\mu) \tilde{\Phi}^A_{ab,n}(\vec{p}, \mu), \quad (3.11)$$

$$p > \mu, \quad \mu_{ab} = \frac{ma mb}{ma + mb}.$$

From (2.8) it trivially follows that the eight components of the octet wave function fulfill the same equation and hence they are the same. Notice that the wave function normalization and the differential equation above are \(\mu\) dependent. Furthermore, the wave function is not defined over all values of \(p\). We will work this out in detail in sect. 4. In order to simplify the notation we will not displayed the cut-off dependence explicitly in the rest of the section, but it must be understood throughout.

For \(H^\mu\), the hamiltonian associated to \(L^\mu\), we denote the eigenstates and eigenvalues by

$$\left|(ab, g, s; \vec{v})\right|, \quad E_g \quad (3.12)$$

where \(g\) labels the low momentum state. We cannot give explicit expressions since their dynamics is governed by low momentum. (3.12) has the non-relativistic normalization

$$\langle (ab, g, s; \vec{v}' | (a'b', g', s'); \vec{v}' \rangle = (2\pi)^3 \delta^{(3)}(m_{ab}(\vec{v} - \vec{v}')) \delta_{s,s'} \delta_{(ab),(a'b')} \delta_{g,g'} \quad (3.13)$$

Of course the states (3.1) and (3.12) are orthogonal since they belong to different momentum regimes.

Our Hilbert space is (before switching on \(L^\mu\)) \(\{N\} = \{(n, A), g\}\) and the identity reads in this base

$$1 \approx |0\rangle\langle 0| + 1 + 1 = |0\rangle\langle 0| + \sum_{ab, N, s} \int \frac{d^3\vec{p}}{(2\pi)^3} \left|(ab, N, s; \vec{v})\right\langle (ab, N, s; \vec{v}') \right| \quad (3.14)$$

Let us now calculate the matrix elements of \(H^\mu\), the hamiltonian associated to \(L^\mu\). We note that the only matrix element different from zero is

$$\langle (ab, g, s; \vec{v})| H^\mu| (a'b', n, s', A); \vec{v}' \rangle$$
\[
(2\pi)^3 \delta^{(3)}(m_{ab}(\vec{v} - \vec{v}')) E^A_{ab,rr'} \tilde{\Phi}^A_{ab,rr'}(0) \frac{f_{ab,g}^A N_A}{\sqrt{2}} \delta_{s',s}(a',(a''b')) ,
\]
where
\[
\langle (ab, g, s); \vec{v} | \bar{h}^A_a T^A \Gamma^s h^A_b(0) | 0 \rangle =: f_{ab,g}^A \delta_{s',s}
\]

In the calculations above we have not made any explicit assumption about the relative size of \( L_\mu \) and \( L^\mu \). We are mainly interested in very heavy quark-antiquark bound states where small momenta can be considered as corrections, at least for the lower energy levels. Clearly, these bound states should be singlets since the octet potential is repulsive. In fact, at the level we are working, the octet states are not going to give contributions to the physical observables so we will neglect them in the following. Hence from now on colour singlets are understood and colour indeces dropped. We also remark that we are always working in the CM frame, even though sometimes we keep \( \vec{v} \neq 0 \) in some intermediate steps for convenience. Following standard Quantum Mechanics perturbation theory \cite{17} we can obtain the corrected bound state energy\footnote{The correction to the bound state energy was found to be zero in \cite{6} because the existence of \( L^{\mu}_\mu \) was not known.} and wave functions (states) for the lower levels. They read

\[
\delta E_{ab,n} = \frac{\Pi_{ab}(E_{ab,n})}{2N_c} |E_{ab,n} \tilde{\Phi}_{ab,n}(0)|^2
\]

\[
|\langle (ab, n, s); \vec{v} | F = Z_n^{1/2} |(ab, n, s); \vec{v} \rangle_F
\]

\[
|\langle (ab, n, s); \vec{v} | F = |\langle (ab, n, s); \vec{v} \rangle + |\langle (ab, n, s); \vec{v} \rangle^{(1)} + |\langle (ab, n, s); \vec{v} \rangle^{(2)} + ... \]

\[
|\langle (ab, n, s); \vec{v} \rangle^{(1)} = \frac{\tilde{\Phi}_{ab,n}(0)^* E_{ab,n} \hat{\Gamma}^\mu(E_{ab,n}) J^{ab}_{\mu}(0)|0\rangle}{\sqrt{2N_c}}
\]

\[
|\langle (ab, n, s); \vec{v} \rangle^{(2)} = \sum_{m \neq n} |\langle (ab, m, s); \vec{v} \rangle | E_{ab,n} \tilde{\Phi}^*_{ab,m}(0) \bar{\Phi}_{ab,n}(0) \frac{E_{ab,m}}{E_{ab,n} - E_{ab,m}}
\]

\[
Z_n \approx 1 + \frac{1}{2N_c} \frac{d\Pi_{ab}(E_{ab,n})}{dE_{ab,n}} |E_{ab,n} \tilde{\Phi}_{ab,n}(0)|^2
\]

\[
\hat{G}^\mu(z) := \frac{1}{z - H^\mu + i\epsilon}
\]
We should stress that in the last two equations there is only small momentum dynamics. High energies may come from the external bound state energy insertion.

Some comments are in order. Notice first that for \( l \neq 0 \) (angular momentum) the wave function (state) and the energy remain unchanged. Notice also that the s-wave state does not receive contributions from \( l \neq 0 \) states either. The previous statement is true due to the fact that the momentum wave function at zero momentum for \( l \neq 0 \) is zero. This means that the new interaction does not couple \( l = 0 \) states with \( l \neq 0 \) states. This result would change if we kept further terms in the effective lagrangian (see (2.16)) but, of course, these contributions would be subleading.

Let us next calculate the decay constant. In order to do it we split the current as in the last section. The soft current only gives a contribution with the low momentum states \( g \) in the same way as the hard current only gives a contribution with the modified Coulomb bound states. Thus, we obtain

\[
\langle 0 | J^a_{\Gamma}(0) \rangle (ab, g, s); \bar{v} \rangle = \langle 0 | J^a_{\Gamma}(0) \rangle (ab, g, s); \bar{v} \rangle = -\frac{tr \left( \Gamma^* \Gamma \right)}{2} f_{ab,g} \tag{3.25}
\]

\[
\langle 0 | J^a_{\Gamma}(0) \rangle (ab, n, s); \bar{v} \rangle = \langle 0 | J^a_{\Gamma}(0) \rangle (ab, n, s); \bar{v} \rangle = -tr \left( \Gamma^* \Gamma \right) \sqrt{\frac{N_c}{2}} \Phi_{ab,n}(0) \tag{3.26}
\]

Finally the decay constant reads (changing to relativistic normalization)

\[
\langle 0 | J^a_{\Gamma}(0) \rangle (ab, n, s); \bar{v} \rangle_F = -tr \left( \Gamma^* \Gamma \right) \sqrt{\frac{m_{ab,n} N_c \Phi_{ab,n}(0)}} \times \left\{ 1 + \frac{1}{4N_c} \frac{d \Pi_{ab} (E_{ab,n})}{dE_{ab,n}} E_{ab,n} \Phi_{ab,n}(0) \right\}^2 + \frac{\Pi_{ab} (E_{ab,n})}{2N_c} \frac{\Phi_{ab,n}(0)}{E_{ab,n} - E_{ab,m}} \times \left( 1 + \sum_{m \neq n} \Phi_{ab,m}(0) \Phi_{ab,m}(0) \frac{E_{ab,m}}{E_{ab,n} - E_{ab,m}} \right) \right\} \tag{3.27}
\]

Finally let us obtain the bilinear currents at zero recoil. For that we need to know

\[
J^a_{\Gamma}(0) = \tilde{Q}^a \Gamma Q'(0) \tag{3.28}
\]

\[
\langle (ac, N', s'); \bar{v} | J^a_{\Gamma}(0) \rangle (ab, N, s); \bar{v} \rangle \tag{3.29}
\]

In order to deal with them we need to perform the splitting between large and small momentum. However this current cannot be in general splitted in two terms. We have
mixing between large and small momentum. Fortunately, the mixing terms disappear
if both initial and final states have the same velocity. This will not longer be true for
non-zero recoil matrix elements. Thus, we obtain

\[ \langle (ac, n', s'); \vec{v}|J_{t'}^{bc}(0)|(ab, n, s); \vec{v}'\rangle = -\frac{tr(\Gamma' \Gamma^*)}{2} \int d^3k (2\pi)^3 \hat{\Phi}^*_{ac, n'}(\vec{k}) \hat{\Phi}_{ab, n}(\vec{k}) \]  

(3.30)

\[ \langle (ac, g', s'); \vec{v}|J_{t'}^{bc}(0)|(ab, g, s); \vec{v}'\rangle = -\frac{tr(\Gamma' \Gamma^*)}{2} f_{ac, ab}^{d'g} \]  

(3.31)

where we have used

\[ \sum_s (\Gamma^*_a)_{\alpha_2\alpha_4} (\Gamma^*_b)_{\alpha_1\alpha_3} = -2(p_+\gamma_{\alpha_2\alpha_4})(p_-\gamma_{\alpha_1\alpha_3}). \]  

(3.32)

The remaining possible matrix elements are zero. Notice that \( f_{ac, ab}^{d'g} = \delta_{d'g} \) because of baryonic charge conservation.

The physical matrix element reads (again with relativistic normalization)

\[ F\langle (ac, n', s'); \vec{v}|J_{t'}^{bc}(0)|(ab, n, s); \vec{v}\rangle_F = -\sqrt{m_{ab,n} m_{ac,n'}} tr(\Gamma' \Gamma^*) \]

\[ \times \left\{ \int d^3k (2\pi)^3 \hat{\Phi}^*_{ac, n'}(\vec{k}) \hat{\Phi}_{ab, n}(\vec{k}) \right\} \]

\[ \times \left\{ 1 + \frac{1}{4N_c} \frac{d\Pi_{ab}(E_{ab,n})}{dE_{ab,n}} |E_{ab,n} \hat{\Phi}_{ab, n}(\vec{0})|^2 + \frac{1}{4N_c} \frac{d\Pi_{ac}(E_{ac,n'})}{dE_{ac,n'}} |E_{ac,n'} \hat{\Phi}_{ac, n'}(\vec{0})|^2 \right\} \]

\[ + \frac{\Pi_{ac}(E_{ac,n'})}{2N_c} E_{ac, n'} \sum_{m \neq n'} \Phi^*_{ac, m}(\vec{0}) \Phi_{ab, n}(\vec{0}) \frac{E_{ac, m}}{E_{ac, n'}} \frac{E_{ab, n}}{E_{ab, m}} \int \frac{d^3k}{(2\pi)^3} \hat{\Phi}_{ac, m}(\vec{k}) \hat{\Phi}_{ab, n}(\vec{k}) \]

\[ + \frac{\Pi_{ab}(E_{ab,n})}{2N_c} E_{ab, n} \sum_{m \neq n} \Phi^*_{ab, m}(\vec{0}) \Phi_{ab, n}(\vec{0}) \frac{E_{ab, m}}{E_{ab, n}} \frac{E_{ab, n}}{E_{ab, m}} \int \frac{d^3k}{(2\pi)^3} \hat{\Phi}_{ac, n'}(\vec{k}) \hat{\Phi}_{ab, m}(\vec{k}) \]  

(3.33)

where

\[ \int d^4x_1 d^4x_2 e^{iP_{n'} x_1} e^{-iP_{n} x_2} \langle 0|T \left\{ J_{t'}^{ca}(x_1) J_{t'}^{bc}(0) J_{t'}^{ab}(x_2) \right\} |0\rangle \]

\[ =: \Pi_{ac, ab}(E_{ac, n'}, E_{ab, n}) tr(\Gamma' \Gamma^*) \]  

(3.34)

We can easily check that orthonormality is fulfilled when \( b = c \). We expect the last
statement to be true since spin symmetry relates the matrix element with the baryonic
charge when \( b = c \).

\(^6\)This was not always the case for the result given in [6].
Before finishing this section let us make some remarks. Both correlators (3.24) and (3.34) should be small quantities for perturbation theory to hold. This is the case if \( \mu \) is small against the typical momentum in the Coulomb interaction (i.e. \( \mu a_{ab,n} \ll 1 \), where \( a_{ab,n} = n/\mu_{ab} \alpha \) is the Bohr radius). It constrains the possible applications to the lower energy levels. On the other hand \( \Lambda_{QCD} \ll \mu \) should hold so that the low momentum dynamics is not strongly affected by the cut-off.

4 Cut-off independence

Our results in the last section may look like strongly cut-off dependent. We have two sources of cut-off dependence. On the one hand we have a cut-off separating small momentum gluons from large momentum gluons. This cut-off is the responsible for the absence of Coulomb interaction in \( L^{\mu} \). It has been mentioned at several instances but it has never been written down explicitly in the formulas. This cut-off dependence has been analysed before [15] so we shall ignore it in the following. On the other hand we have the cut-off separating large and small relative momenta. It plays the role of an infrared cut-off in the perturbative Coulomb wave function (large momentum) and the role of an ultraviolet cut-off for the small momentum contributions. We prove in this section the cut-off independence to the desired order (\( \mu^2, \Lambda_{QCD}^3 \)) of this last cut-off. This is crucial to ensure that our approach respects colour \( SU(3) \) gauge symmetry. It is important to use the same cut-off procedure in both large and small momentum regions in order to neatly cancel the cut-off dependence. We use a hard three momentum cut-off for convenience, as we have done in the previous sections.

First of all, let us study the cut-off dependence in the low momentum correlators we found in the last section. Although they are non-perturbative objects we can always perform a perturbative calculation in order to see how they depend on the cut-off.

Let us start by (2.13) (which is formally equal to NRQCD). For (3.24) we obtain at the lowest order (in the CM frame, \( \bar{\nu}' = 0 \))

\[
\Pi_{ab}(k^0) = -\frac{N_c \mu_{ab} \mu}{\pi^2} \left[ 1 - \frac{1}{2x} \text{tr} \left( \frac{1 + x}{1 - x} \right) \right], \quad x = \frac{\mu}{\sqrt{2 \mu_{ab} (k^0 + i\epsilon)}}. \tag{4.1}
\]

Let us consider two limits.
In the limit $x >> 1$ (i.e. near threshold) it reduces to
\[ \Pi_{ab}(k^0) \simeq - \frac{N_c \mu_{ab} \mu}{\pi^2} \left[ 1 + \frac{i \pi}{2x} \theta(k^0) \right] \] (4.2)
and no pole appears. In the limit $x << 1$ (4.1) reduces to
\[ \Pi_{ab}(k^0) \simeq \frac{N_c \mu^3}{6\pi^2} \frac{1}{(k^0 + i\epsilon)^2}. \] (4.3)

This expression is going to be important in the following. We stress that (4.3) is $\mu_{ab}$ independent, and amounts to drop the $1/m$ terms in (2.13) which is nothing but the HQET for quarks and antiquarks. Let us now look for the physical situation we are interested in. Thus, we take $k^0 = E_{ab,n}$ and we obtain $|x| = \frac{\mu}{\mu_{ab} C_{F} \alpha_{s}}$, but this is nothing but the parameter we need to keep small so that the small relative momentum contributions are subleading, and hence our expansion makes sense. In the following, we always consider that we are in the limit $|x| << 1$.

For (3.34) we obtain ($x << 1$)
\[ \Pi_{ac,ab}(k'^0, k^0) \simeq \frac{N_c \mu^3}{6\pi^2} \frac{1}{(k^0 + i\epsilon)^2}. \] (4.4)

At this point we would like to stress that in the limit $x << 1$ the same perturbative results are found using HQET. This is going to be determinant in the next section.

(4.3) and (4.4) make explicit the UV cut-off dependence coming from the small relative momentum region. Let us next go on with the IR cut-off dependence coming from the large relative momentum region.

Let us then study the cut-off dependence of the wave function (for simplicity we omit the flavour indices). In order to do it we solve the wave equation (3.11) perturbatively in $\mu$. $n$ labels continuum or discrete spectrum. Because of the radial symmetry, we can write (we follow [18])
\[ \tilde{\Phi}_{n,l,m}(\vec{p}; \mu) = F_{n,l}(p; \mu) Y_{l,m}(\hat{p}) \] (4.5)
where $F_{n,l}(p; \mu)$ fulfills
\[ \frac{p^2}{2\mu_{ab}} F_{n,l}(p; \mu) - \frac{C_F \alpha_s}{\pi \mu} \int_{\mu} q dq F_{n,l}(q; \mu) Q_{d}(\frac{p^2 + q^2}{2pq}) = E_{n}(\mu) F_{n,l}(p; \mu) \quad p > \mu \] (4.6)
\[ Q_l(z) = \frac{1}{2} \int_{-1}^{1} dx \frac{P_l(x)}{z - x} \]  

(4.7)

and \( P_l \) is the Legendre function of the first kind. We stress that we are interested in \( F_{n,l}(p; \mu) \) for \( p > \mu \) only, although in the intermediate steps it is going to be defined over all \( p > 0 \) values. Now we perform a cut-off parameter expansion and we work as in usual quantum mechanics perturbation theory where we demand the corrections to be orthogonal to the leading result.

\[
F_{n,l}(p; \mu) = \sum_{r=0}^{\infty} F_{n,l}^r(p) \frac{\mu^r}{r!} \quad E_n = \sum_{r=0}^{\infty} E_n^r \frac{\mu^r}{r!}. 
\]

(4.8)

We also expand the cut-off in the integral.

\[
\int dq \frac{F_{n,l}^r(q)}{2pq} \Phi_l(q) \approx \hbar_{n,l}(p, \mu) \equiv \sum_{i=0}^{\infty} h_{n,l}^i(p) \frac{\mu^i}{i!}.
\]

(4.9)

On general grounds we can see that the corrections to the Coulomb wave function and energy go like \( O(\mu^{2+3}) \), therefore, as expected, we can neglect the \( l \neq 0 \) states since their contributions are subleading.

At leading order we obtain the standard Schrödinger equation with a Coulomb potential with no \( \mu \) dependence. Furthermore, for the following terms in perturbation theory we obtain

\[
E_n^1 = E_n^2 = 0, \quad \Phi_n^1(p) = \Phi_n^2(p) = 0.
\]

(4.10)

Finally to third order we obtain

\[
E_n^3 = -E_n \frac{\lvert \Phi_n(\vec{0}) \rvert^2}{\pi^2},
\]

(4.11)

\[
\Phi_n^3(p) = \sum_{m \neq n} \Phi_m(p) \frac{\Phi_m^*(\vec{0}) \Phi_n^*(\vec{0})}{\pi^2} \frac{E_m}{E_m - E_n}.
\]

(4.12)

We have not yet normalized the cut-off dependent wave function, as we can see from

\[
\int_{\mu} \frac{d^3q}{(2\pi)^3} \Phi_n^*(q) \Phi_n(q) \simeq 1 - \frac{\lvert \Phi_n(\vec{0}) \rvert^2}{6\pi^2} \frac{\mu^3}{12\pi^2}.
\]

(4.13)

Therefore, we must change

\[
\Phi_n(p; \mu) \rightarrow \Phi_n(p; \mu) \left( 1 + \frac{\lvert \Phi_n(\vec{0}) \rvert^2}{12\pi^2} \right) \mu^3.
\]

(4.14)
(4.11)-(4.14) provide the explicit IR cut-off dependence from the large relative momentum region.

We have obtained the explicit cut-off dependence to the desired order $\mu^3$ in both large and small momentum regions. Now we will see they match properly, that is, the observables are cut-off independent. In fact what we will see is that the physical states (3.18) themselves are already cut-off independent. In this way we prove the cut-off independence for any observable.

Consider first the bound state energy $E_n^F$

$$E_n^F = E_n + \delta E_n.$$  \hspace{1cm} (4.15)

The cut-off dependence of $E_n$ is given by (4.8), (4.10) and (4.11), whereas the cut-off dependence of $\delta E_n$ is given by (3.17) and (4.3). One can then easily check that $E_n^F$ is cut-off independent.

Consider next the state $|((ab, n, s); \vec{v})_F \rangle$ in (3.19). Recall that the first and last term on the rhs belong to the large relative momentum region whereas the term in the middle belongs to small relative momentum region. Let us keep apart for a moment the explicit cut-off separating these two regions in the relative momentum integrals. The remaining cut-off dependences of the first term are given by (4.5), (4.8), (4.10) and (4.12), while for the last term are given by (3.21) and (4.3), which cancel each other.

It remains the UV cut-off dependence coming from (3.20) (which has been already studied in [6]) and the explicit IR cut-off dependence coming from the integral over relative momentum in the first term of (3.19) (see (3.1)), which we kept apart for a while. Recall that the wave function in the first term of (3.19) is, except for the normalization factor (4.14), the Coulomb wave function since we have already cancelled the cut-off dependences coming from (4.12). Let us next calculate (3.20) perturbatively at lowest order. It reads

$$|((ab, n, s); \vec{v})^{(1)} \rangle = \frac{\tilde{\Phi}_{ab, n}(0)}{\sqrt{2N_c}} \int \mu \frac{d^3 \vec{K}}{(2\pi)^3} a^\alpha(p_1) \Gamma_{\alpha \beta}(p_2) b^\dagger_{\alpha, i}(p_1) d_{\beta, i}(p_2) |0\rangle$$

$$= \frac{1}{\sqrt{2N_c}} \int \mu \frac{d^3 \vec{K}}{(2\pi)^3} \tilde{\Phi}_{ab, n}(\vec{K}) a^\alpha(p_1) \Gamma_{\alpha \beta}(p_2) b^\dagger_{\alpha, i}(p_1) d_{\beta, i}(p_2) |0\rangle.$$  \hspace{1cm} (4.16)

The second equality holds at the order we are working at. Notice finally that this is nothing but the piece we need to add to the first term of (3.19) in order to obtain a relative momentum integral independent of the cut-off. Finally, the cut-off dependence
of the normalization in (4.14) and of (3.22) also cancel each other in (3.18) (again taking into account (4.3)).

We have thus seen that at the level of physical states we are able to prove the cut-off independence. The cut-off independence can also be checked explicitly in the observables (3.17), (3.27) and (3.33). This demonstrates that the HQET ultraviolet behavior cancels the NRQCD infrared behavior in Coulomb type bound states, which guarantees that we have performed a proper matching between large and small relative momentum. This issue has also been pursued in [6, 15].

5 Evaluation of the low momentum correlators

In section 3 we learnt how to parametrize the possible non-perturbative contributions in the small relative momentum region in terms of two low momentum correlators ((3.24) and (3.34)) with external Coulomb bound state energy insertions. It is remarkable that these contributions only exist for s-states. At the beginning of section 4 we also saw that the kinetic term, which is suppressed by a mass invers power, can be safely neglected in the correlators we are interested in, and hence we can use HQET for quarks and antiquarks to discuss their properties.

The HQET for quarks and antiquarks enjoys a $U(4N_{hf})$ symmetry which breaks spontaneously down to the $U(2N_{hf}) \otimes U(2N_{hf})$ Isgur-Wise symmetry [12].

Let us first analyse the consequences of the unbroken $U(2N_{hf}) \otimes U(2N_{hf})$ symmetry. In fact the spin symmetry which is included in it has already been used in (3.24) and (3.34). The flavour symmetry implies the following

$$f_{ab,g} = f_g, \quad f_{ac,ab}^{d,g} = f^{d,g}.$$

Therefore we get

$$\Pi_{ab}(E_{ab,n}) = \Pi_1(E_{ab,n}), \quad \Pi_{ac,ab}(E_{ac,n'}, E_{ab,n}) = \Pi_2(E_{ac,n'}, E_{ab,n}).$$

The correlators (3.24) and (3.34) are thus given in terms of two unknown universal (flavour independent) functions $\Pi_1$ and $\Pi_2$. But if we go further, using flavour number conservation together with flavour symmetry, we obtain $f^{d,g} = \delta_{d,g}$ for any flavour. From that it follows that if $\Pi_1$ is known for any energy insertion we can obtain $\Pi_2$. 

Explicitly they read

\[ \Pi_1(E_{ab,n}) = \sum_g \left| \frac{f_g}{2} \right|^2 \frac{1}{E_{ab,n} - E_g}, \quad \Pi_2(E_{ac,n'}, E_{ab,n}) = \sum_g \left| \frac{f_g}{2} \right|^2 \frac{1}{E_{ac,n'} - E_g} \frac{1}{E_{ab,n} - E_g}. \]  

(5.3)

These low momentum correlators can be further specified at least in two situations.

i) \( E_{ab,n} \gg \Lambda_{QCD}, \quad (m_Q \to \infty, \quad \alpha \text{ small}), \)

ii) \( E_{ab,n} \ll \Lambda_{QCD}, \quad (m_Q \text{ large}, \quad \alpha \to 0). \)

Notice that situation ii) is conceivable if \( \alpha \) is very small since so far we have only assumed that the inverse Bohr radius is much bigger than \( \Lambda_{QCD} \) and the energy is suppressed by a factor \( \alpha \) with respect to the former.\(^7\) \(^8\)

In the situation i) the operator product expansion holds. If we carry it out for the low momentum correlators we just obtain (4.3) and (4.4). Their cut-off dependence just cancels the cut-off dependence from the large relative momentum region, as we saw in section 4. Hence, we conclude that there are no new non-perturbative contributions in this situation, thus confirming the fact that the VL contributions from the condensate are indeed the leading non-perturbative effects in the \( m_Q \to \infty \) limit.\(^9\) This result follows from the observation that there is no local gauge invariant object that can be built out of \( D_0 \) alone. We have explicitly checked it for lower order terms.

In the situation ii) we are in the low energy regime of the HQET. In this regime it is important that the HQET with quarks and antiquarks with the same velocity undergoes a spontaneous symmetry breaking of a \( U(4N_{hf}) \) symmetry down to the Isgur-Wise symmetry \( U(2N_{hf}) \otimes U(2N_{hf}) \), since the Goldstone modes associated to the broken generators dominate the dynamics. The Heavy Quark Hadronic effective lagrangian describing the Goldstone modes was worked out in \(^6\), where the correlators (3.24) and (3.34) were also calculated. Using those results we obtain

\[ \Pi_{ab}(k^0) = \frac{f_h^2}{2} \frac{1}{(k^0 + i\epsilon)}, \]  

(5.4)

\(^7\)In practise we must remember that \( \alpha \) should better be substituted by the running coupling constant at the quarkonium scale, which is in fact an implicit function of \( m_Q \) and \( \Lambda_{QCD} \).

\(^8\)In [6], the bound state energy \( E_{ab,n} \) was understood as giving rise to a residual mass for the heavy quark and antiquark in the Heavy Quark Effective lagrangian, which was later on subtracted. That definitely obscures its actual role, which eventually led to some confusion: in [6] the situation ii) was not allowed whereas the Heavy Quark Hadronic lagrangian was used for situation i), which is not correct.

\(^9\)This point was not properly specified in [6].
\[ \Pi_{ac,ab}(k^0, k') = \frac{\pi_H^2}{2} \frac{1}{(k^0 + i\epsilon)} \frac{1}{(k'^0 + i\epsilon)}, \quad (5.5) \]

\[ \frac{\pi_H^2}{2} = \frac{\pi_H^2}{2} + \frac{\mu^3 N_c}{6\pi^2}, \quad (5.6) \]

where \( \pi_H^2 \) is cutoff independent. Notice that in this situation all non-perturbative effects in the small relative momentum region are parametrized by a single nonperturbative constant which is spin and flavour independent. This is a non-trivial consequence of the \( U(4N_{hf}) \) symmetry. The fact that the latter is spontaneously broken down to \( U(2N_{hf}) \otimes U(2N_{hf}) \) allows us to know the Green function behaviour at low energy insertion with a single nonperturbative constant since no mass term appears in the pole. All the spin and flavour dependence is explicitly known in the observables.

However, caution must be taken in the situation ii). This is due to the fact that, in this situation the standard evaluation of non-perturbative contributions in the large relative three momentum region coming from (2.15) becomes unreliable. Let us briefly recall the two approximations involved, namely the multipole expansion and the adiabatic approximation. The first one is an expansion in \( \Lambda_{QCD} \) over the inverse Bohr radius, which has also been assumed to hold throughout. The second one requires the time evolution of the soft gluon fields to be slow in comparison with the energies involved in the Coulomb spectrum. This requirement is in fact the opposite of situation ii). Thus we are in the unfortunate even that when we have an excellent parametrization of the non-perturbative effects in the small relative momentum region ((5.4) and (5.5)) we loose control of them in the large relative momentum region.

Nevertheless, we envisage a situation where the parametrization (5.4) and (5.5) may be useful. Recall that although the parameter controlling the adiabatic approximation and the parameter controlling the expansion in the hadronic effective lagrangian are both of order \( \Lambda_{QCD} \), they need not be exactly the same. The former was shown to be \( <DFDF>/ <FF> \) in [7] and let \( 2\pi \tilde{\pi}_H^2 \) be the latter. Suppose then that

\[ E_{ab,n} > \left( \frac{<DFDF>}{<FF>} \right)^{\frac{1}{2}}, \]

\[ \text{Notice also that although at first sight the contributions obtained by substituting (5.4) and (5.5) in (3.17), (3.27) and (3.33) look like more important than those from the condensate when } m_Q \rightarrow \infty, \text{ they are actually not so since the smallness of } \alpha \text{ required in situation ii) maintains the condensate contribution dominant. Some statements made in [6] implying the opposite must be corrected.} \]
In such a situation it would be reasonable to use both the adiabatic approximation in the large relative momentum region and the hadronic effective lagrangian in the small relative momentum region. Some bottomonium, charmonium, and presumably $B_c$ states may well be considered in the situation (5.7). However, the mass of the $b$ quark and mainly the mass of the $c$ quark are not large enough to allow for a straightforward application of our formalism to phenomenology. Relativistic and radiative corrections are in general important and this is also so for the non-perturbative corrections due to the gluon condensate [3]. All them must be taken into account.

Let us next discuss the expected size of our contributions. It is not our aim to present a full-fledged phenomenological analysis in order to extract $f_H^0$ from the data, which would definitely be premature as it should be clear in the following discussion, but just give reasonable estimates of the expected magnitude of its contributions. For simplicity, we will concentrate on the mass corrections.

We start with the bottomonium system where our formalism is expected to apply for the lowest lying states [3, 19]. We proceed as follows. First of all, we fix $m_b$ and $a_{b_{0,1}}$ using the experimental data and the available theoretical results while ignoring the contribution $\delta E_{ab,n}$ in (3.17). Then we estimate the size of $\delta E_{ab,n}$ by letting $f_H^0$ run within values of the order of $\Lambda_{QCD}$. We should keep in mind that although we will take $f_H^0$ positive for definiteness it can also be negative. We extract $m_b$ and $a_{b_{0,1}}$ from the selfconsistency equation $a_{b_{0,1}}(\alpha_s(a_{b_{0,1}})) = a_{b_{0,1}}$, and the $\Upsilon(1s)$ mass. We use the following equation to fit the latter

$$m_{\Upsilon(1s)} = 2m_b + A_2 + A_3 + A_{VL}$$

where

$$a_{b_{0,1}}^{-1} = \frac{m_b C_f \hat{\alpha}(a_{b_{0,1}})}{2}$$

$$A_2 = -2m_b C_f^2 \hat{\alpha}^2(a_{b_{0,1}})$$

$$A_3 = -2m_b C_f^2 \beta_0 \hat{\alpha}^2(a_{b_{0,1}}) \hat{\alpha}(a_{b_{0,1}})^{-1} \left( \ln \left[ \frac{(a_{b_{0,1}})^{-1}}{m_b C_f \hat{\alpha}(a_{b_{0,1}})} \right] + 1 - \gamma_E \right)$$

$$A_{VL} = m_b e^{10\pi(\alpha_s G^2)} \left( \frac{m_b C_f \hat{\alpha}(a_{b_{0,1}})}{4} \right)^4$$

(5.10)
We have taken the formulas above, which include relativistic, radiative and the VL non-perturbative corrections, from [3]. We allow for different values of $\Lambda_{QCD}$ and give the relative weight of each contribution in the table.

Let us next assume that we are in the situation (5.7). As mentioned before, this may well be the case for the $\Upsilon(1S)$, $\Upsilon(2S)$, $\chi_b(1P)$, $J/\psi$ (and $\eta_c$) and $B_c$ (and $B_c^*$). If we let $f_H^{2/3}$ run between the values

$$E_{ab,n} < 2 \pi f_H^{2/3} < a_{ab,n}^{-1}$$  \hspace{1cm} (5.11)

we can give an estimate of $\delta E_{ab,n}$. If we allow $f_H^{2/3}$ between $100 - 150 MeV$, our results turn out to be quite stable under values of $\Lambda_{QCD}^{nf=3}$ between $200 - 300$ MeV. We obtain

$$-9 MeV < \delta E_{bb,1} < -2 MeV$$  \hspace{1cm} (5.12)

where the explicit expression used for calculating $\delta E_{ab,n}$ reads

$$\delta E_{ab,n} = -4 \mu_{ab} \frac{16 \pi f_H^n}{N_c C_f \hat{\alpha}_s(a_{ab,n}^{-1})} \left( \frac{n}{2 \mu_{ab}} \right)^3$$  \hspace{1cm} (5.13)

Although the smallness of the result above is discouraging at first sight, it justifies the procedure used and makes it selfconsistent.

For $n = 2$ we obtain

$$-55 MeV < \delta E_{bb,2} < -15 MeV$$  \hspace{1cm} (5.14)

Recall that only the s-wave states receive this correction. If the sign of $f_H^n$ was negative, the signs above would be reversed. This would help to understand the mass difference between the $\chi_b(1P)$ and the $\Upsilon(2S)$.

Let us finally give some estimates for $\delta E_{cc,1}$ and $\delta E_{bc,1}$ corresponding to the $J/\Psi$ (and $\eta_c$) and the $B_c$ (and $B_c^*$) ground states. We have taken the mass of the charm $m_c = 1570 MeV$ as given in reference [3]. For the $J/\Psi$ we find

$$\delta E_{cc,1} \sim -42 MeV$$

taking $\Lambda_{QCD}^{nf=3} = 300 MeV$, $a_{cc,1}^{-1} = 848 MeV$ and $f_H^{2/3} = 150 MeV$. For the $B_c$ we find

$$\delta E_{bc,1} \sim -23 MeV$$

\footnote{However, we have not taken into account the contributions of order $O(\alpha^4, \alpha^5)$ given in [3] since the complete calculation at this order is still lacking.}
taking $\Lambda_{QCD}^{n_f=3} = 300 \, \text{MeV}$, $a_{b,c}^{-1} = 1013 \, \text{MeV}$ and $f_H^{2/3} = 150 \, \text{MeV}$.

The above contributions for the energy shifts are, on the one hand, small enough to make us confident that our results are under control and, on the other hand, large enough to hope for its eventual observation. However, it is important to realize that the VL contributions are exceedingly big for $Y(2S)$, $\chi_b(1P)$, $J/\psi$ (and $\eta_c$) and $B_c$ (and $B_c^*$). We suspect that the framework used so far to calculate the VL contributions in the large relative momentum region is not appropriated for these states. We believe that in order to make realistic QCD-based predictions for these states one should devise a reliable approximation in the large relative momentum region to deal with the situation (ii) above, namely invers Born radius and energy larger and smaller than $\Lambda_{QCD}$ respectively. Work in this direction is in progress [20].

6 Conclusions

We are confident that the theoretical framework above is going to be useful for an eventual QCD-based formalism attempting to encompass situations where the Coulomb energy is large (small $n$) and situations where it is small (large $n$) with respect to $\Lambda_{QCD}$ in heavy quarkonium. Even more, this formalism could also be useful in order to obtain explicitly the perturbative Coulomb corrections to the non-perturbative heavy quarks bound states (large $n$).

Our formalism is clearly inspired by the Wilson renormalization group approach. We separate the fields into large and small momentum components by an explicit cut-off, and work out what the effective action for the latter is. However, there is an important point which makes our formalism rather peculiar: integrating out the large momentum components does not give rise to local counterterms only. There is non-trivial physics in the ultraviolet, namely Coulomb type bound states. As far as we know, this is the first example of a Wilsonian approach where effects due to bound states have been taken into account.

Let us finally summarize the main contributions of this paper. Elaborating on the ideas first presented in [6], we have produced a detailed derivation of the effective theory governing the small relative momentum degrees of freedom in heavy quarkonium. In particular this includes an interaction term, which had been overlooked before, that leads to a few corrections in the observables. We have proven the cut-off independence
of the formalism. We have also discussed in detail when non-perturbative contributions which cannot be expressed in terms of local condensates arise, namely when a description in terms of a Heavy Quark Hadronic Theory is adequate. Our preliminary estimations suggests that these contributions lead to energy shifts of a few tens of MeV. Unfortunately, more theoretical work is necessary to establish them from the data. This is mainly due to the lack of control on the non-perturbative effects in the large relative momentum region of most of the systems where our approach should apply, namely $\Upsilon(2S), \chi_b(1P), J/\psi$ (and $\eta_c$) and $B_c$ (and $B'_c$). Work in this direction is in progress [20].

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References


Table 1: We display $A_2$, $A_3$ and $A_{VL}$ defined in (5.10). The last two columns give our results for $m_b$ and $a_{bb,1}^{-1}$.

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<th>$\Lambda^{n_{3q}}_{QCD}(MeV)$</th>
<th>$A_2(MeV)$</th>
<th>$A_3(MeV)$</th>
<th>$A_{VL}(MeV)$</th>
<th>$m_b(MeV)$</th>
<th>$a_{bb,1}^{-1}(MeV)$</th>
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