QUARKONIUM IN THE WILSON LOOP FORMALISM: NEW PERSPECTIVES

N. Brambilla \textsuperscript{a} and A. Vairo \textsuperscript{b}

\textsuperscript{a} Institut für Theoretische Physik, Universität Wien
Boltzmanngasse 5, A-1090 Vienna, Austria

\textsuperscript{b} Institut für Hochenergiephysik, Österreichische Akademie der Wissenschaften
Nikolsdorfgasse 18, A-1050 Vienna, Austria

Abstract

We review the present knowledge on the heavy quark interaction. The framework is the NRQCD effective theory and the interaction is expressed in terms of Wilson loop chromoelectric and chromomagnetic insertions.

1 Introduction

The experimental evidence that for heavy quark bound states like $b\bar{b}$, $c\bar{c}$, ... all the splittings are considerably less than the masses suggests that all the dynamical energy scales of these systems are small with respect to the quark masses. As a consequence the quark velocities $v$ are small and these systems can be considered as non relativistic. The hierarchy of the scales is then the typical one of a non relativistic system. Called $m$ the mass of the heavy quark, the quark momenta scale like $mv$ and the quark energies like $mv^2$. Trickier is the situation in the gluon sector, but the binding interaction is essentially characterized by the same energy scale distribution. Therefore the dominant gluon interaction among heavy quarks appears “instantaneous”. A potential picture should hold, at least in first approximation, and the energy levels can be obtained by solving the corresponding Schrödinger equation. In particular for infinitely heavy quarks the spin splittings vanish and the system can be described only by means of a single static central potential.

Indeed, potential models have been quite successful \cite{1} in explaining the hadron spectrum. However, being these phenomenological models, their connection with the QCD parameters is hidden, the scale at which they are defined is not clear and they cannot be systematically improved. In this paper we outline the kind of rigorous description of quarkonium that can be obtained from QCD at the present.

Several perturbative evaluations of the quarkonium potential have been performed in the last twenty years \cite{2} and recently the complete $\alpha_s^2$ corrections to the levels have been calculated \cite{3}. The main difficulty in these calculations is connected with the inclusion of nonperturbative contributions which not only set

\footnote{1 Plenary talk presented at “Nuclear and Particle physics with CEBAF at Jefferson Lab”, Dubrovnik, November 3-10, 1998.}
a bound on the reachable precision, but turn out to be relevant for quantitative
predictions even on the bottomonium spectrum. A typical approach is to consider
all the nonperturbative physics to be encoded into few local condensates. This
corresponds to assume that the typical length associated, for instance, with the
nonperturbative gluodynamics is larger than any other length scale of the system
and can be put equal to infinity. In a pure analytic calculation this statement
is also dictated by the necessity to have a small number of nonperturbative pa-
rameters and maintain predictability. Even for the ground state this assumption
is doubtful, but surely it does not hold for large radii quarkonia (i.e. excited
heavy mesons) where the nonperturbative gluonic length cannot be considered
large with respect to the size of the bound state. Therefore a proper treatment of
nonperturbative effects includes non-local condensates. Another way to say the
same thing is to notice that the contribution to the levels associated with the lo-
cal gluon condensate is proportional to \( n^6 \langle \alpha_s F^2(0) \rangle \). By increasing the principal
quantum number \( n \) beyond the ground state, it grows very soon out of control.

A way of including nonperturbative effects, as they are, in the evaluation of the
quarkonium spectrum is lattice QCD. This technique is becoming more and more
successful and in the near future it is expected to be the only competitive one (see
for instance [4] and references therein). There are different ways in which lattice
QCD calculations can be performed. We mention lattice NRQCD [5] where the
quarkonium spectrum is directly evaluated on the lattice by means of an effective
action derived from QCD by an expansion in the quark velocities. Very close is the
approach we will discuss in the following where the QCD Lagrangian is replaced
by its effective non relativistic formulation before doing any lattice evaluation.
All the non perturbative physics is encoded in this way in the Wilson loop made
up by the quark trajectories and in field strength insertions on it. These (non
local) objects are then evaluated on the lattice. We call this the Wilson loop
approach. The advantage is that in this way our expressions are safer to handle
for lattice purposes, since less affected by finite size effects, and easier to treat also
for analytic purposes like the implementation of vacuum models. In particular
we get an expression for the heavy quark-antiquark potential. In the limit where
the insertions of two field strengths on the Wilson loop can be approximated by a
local condensate, one gets back the “improved” perturbative expression discussed
previously.

The (Wegner)-Wilson loop formalism has a long story. It was first suggested
by K. Wilson [6] that the object called after him Wilson loop would be the relevant
one in order to describe confinement. The strong coupling expansion suggested
an area law behaviour further confirmed by lattice simulations, which were born
with this pioneering work. In [7] it was shown how to relate rigorously the static
Wilson loop with the static quark-antiquark potential. On the same line few
years later also spin dependent corrections to the potentials where expressed in
terms of Wilson loop and chromoelectric and chromomagnetic field insertions [8].
In particular it was proven that a nonperturbative behaviour in the static poten-
tial must give rise to nonperturbative spin dependent corrections as well. Non
spin dependent corrections were treated in the same framework some time later
[9]. This was the situation pictured at the beginning of the '90 in [1]. Several
problems were still open. No relativistic formulation in terms of the Wilson loop
was available. There was an apparent mismatch between the Eichten–Feinberg–
Gromes expression for the spin dependent sector of the potential (analytic in
the quark masses) and the perturbative one loop estimate (containing logarithms
of the quark masses). The inclusion of non potential terms was not understood.
While not all these problems have been solved, some remarkable progress has been
achieved in the last years. Attempts in the direction of a relativistic formulation
have been done in [10]. In [11] it was shown that, by performing properly the
matching between the effective theory and QCD, Wilson coefficients carrying log-
arithms of the quark masses appear in the Eichten–Feinberg–Gromes expression
for the potential. The inclusion of non potential contributions in this frame-
work is now in progress [12] and the way to precision calculations in quarkonium
seems finally to be open. Here we only mention that this goal is met by a new
effective theory built from NRQCD, where explicitly potential and non potential
terms have been separated [13]. Finally, all the Wilson loop averages relevant
for the potential have been calculated on the lattice and the bottomonium and
charmonium spectra have been calculated with good agreement with the data
[14, 15].

In the meantime there have been several progress in the building of a non
relativistic effective theory from QCD, mainly due to the success of Heavy Quark
Effective Theory in describing heavy-light systems [16]. Wilson coefficients have
been calculated at higher order and the role played by reparameterization invar-
iance has been better understood.

Here, we summarize the present level of understanding of the heavy quark
potential which contributes to the energy levels of quarkonium at the order $\nu^4$.
The framework is NRQCD, the tool the Wilson approach.

## 2 NRQCD

In order to define an effective theory we typically need three ingredients: an
effective Lagrangian, a regularization scheme and therefore a matching scale and
a power counting set of rules. The effective theory we will use is NRQCD [1].
Let us discuss its key ingredients.

The NRQCD Lagrangian is obtained from QCD by expanding with respect
to the heavy quark masses. The matching with QCD is performed like in Heavy
Quark Effective Theory (HQET) [17, 18]. We emphasize that in order to build
up the effective Lagrangian from the QCD Lagrangian, we have to ignore the
specific dynamical problem we are dealing with and expand with respect to the
heavy quark masses which are explicit parameters of QCD. Typically the effective
Lagrangian turns out to be the sum of a pure gauge part $L_g$ plus two, four, ... fermion terms ($L_{2f}$, $L_{4f}$, ...). Since we are interested in two-body bound states we will take into account only two and four fermion terms (terms involving more fermions will contribute only in intermediate states). Therefore the NRQCD Lagrangian obtained from QCD by expanding with respect of the mass $m_1$ of a heavy quark and the mass $m_2$ of a heavy antiquark is given by \[17, 18\]

$$L = L_{2f} + L_{4f} + L_g,$$  \(1\)

where at order $1/m^2$ and up to field redefinitions

$$L_{2f} = Q_1 \left( iD_0 + c_2^{(1)} \frac{D^2}{2m_1} + c_4^{(1)} \frac{D^4}{8m_1^3} + c_F^{(1)} g \frac{\sigma \cdot B}{2m_1} + c_D^{(1)} g \frac{D \cdot E - E \cdot D}{8m_1^2} \right) Q_1 + O \left( \frac{1}{m_1^3} \right)$$

$$+ Q_2 \left( -iD_0 + c_2^{(2)} \frac{D^2}{2m_2} + c_4^{(2)} \frac{D^4}{8m_2^3} + c_F^{(2)} g \frac{\sigma \cdot B}{2m_2} - c_D^{(2)} g \frac{D \cdot E - E \cdot D}{8m_2^2} \right) Q_2 + O \left( \frac{1}{m_2^3} \right),$$  \(2\)

$Q_j$ are the heavy quark fields and the covariant derivative is defined as $D_\mu = \partial_\mu + igA_\mu^a T^a$. The explicit form of $L_{4f}$ is given in \[18\]. $L_g$ is the SU(3) Yang–Mills Lagrangian modified in order to give rise to an effective $\alpha_s$ running with 2 (heavy) flavors less.

The effective Lagrangian \([1]\) is not renormalizable. Therefore, it is necessary to regularize it. In a given regularization scheme the reproduction of the correct ultraviolet regime of QCD is obtained by means of the Wilson coefficients. The effective Lagrangian is complete only once these coefficients are given. The Wilson coefficients are evaluated at a matching scale where perturbation theory still holds. They encode the ultraviolet regime of QCD up to a given scale $\mu$ order by order in the coupling constant $\alpha_s$. Renormalization group (RG) transformations should be used in order to resum leading log contributions ($\sim \ln m/\mu$). The matching coefficients are known in the literature at different level of precision (i.e. at different order in the coupling constant), see \[17, 18, 19\]. Here we remember only that reparameterization invariance \[17\] (i.e. the invariance of the effective Lagrangian with respect a variation of $v$) fixes $c_2^{(j)} = c_4^{(j)} = 1$.

Heavy quark bound states are characterized by a dynamical dimensionless parameter, the quark velocity $v$. As explained in the introduction, this parameter is small. From phenomenological potential models \[20\] and from lattice studies \[14\] we get the usually accepted values of $\langle v_b^2 \rangle \sim 0.07$ for the bottomonium system and $\langle v_c^2 \rangle \sim 0.24$ for the charmonium system. This allows the ordering of the energy scales of the problem, $m$, $mv$ and $mv^2$. The first scale has been explicitly integrated out in the Lagrangian \([1]\) and, as discussed above, the contributions
coming from it are carried order by order in $\alpha_s$ by the Wilson coefficients. The last two scales, sometimes known with the name “soft” and “ultrasoft” respectively, are still mixed up. The relation between $v$ and the QCD parameters is in general unknown. It is expected to be the result of perturbative and nonperturbative effects. For infinitely heavy quarks $v$ coincides with $\alpha_s$ like in QED where, for instance in the hydrogen atom, the hyperfine structure constant $\alpha$ is equal to the electron velocity. The main point is that the existence of this small dynamical parameter enables us to set up power counting rules for the operators in the effective Lagrangian. For the sake of simplicity we reproduce here the rough power counting argument of [5], keeping in mind that an exact power counting cannot be done until soft and ultrasoft degrees of freedom have not been completely disentangled. Noticing that the number operator for heavy quarks, $\int d^3 x Q^\dagger(x)Q(x)$, does not depend on $v$, one gets $Q \sim (mv)^{3/2}$. The kinetic energy $\int d^3 x Q^\dagger(x)\frac{D^2}{2m}Q(x)$ is by definition of order $mv^2$ and therefore $D \sim mv$ (i.e. the contribution of $D$ to the levels is of order $mv$). Finally, using the equation of motion one gets the other counting rules $D_0 \sim mv^2$, $gA_0 \sim mv^2$ $gA_0 \sim mv^3$, $gE \sim m^2v^3$ and $gB \sim m^2v^4$. With respect to these rules the Lagrangian of Eq. (2) is accurate up to order $O(v^4)$ of the levels (or up to $O(v^2)$ with respect to the leading contribution). This should guarantee a rough 10% of accuracy on the absolute value of the levels.

Concluding this section, we stress that the NRQCD power counting defined above is not the same as that one used in HQET. In particular, in HQET the kinetic energy $D^2/2m$ is suppressed by a factor $\Lambda_{\text{QCD}}/m$ with respect to the operator $D_0$ while in NRQCD the two operators are of the same order. As a consequence the heavy quark propagator contains in NRQCD a kinetic part which is absent in HQET where the heavy quark propagator is static. In a very general way one can say that these differences are due to the fact that, even if the effective Lagrangian is essentially the same, the physical problem is different and it is the physical problem which defines the counting rules.

3 The Wilson loop formalism

The next step is to derive the heavy quark interaction in the so-called Wilson loop formalism. In this context the use of the effective Lagrangian (2) (with tree level matching) was first suggested by L. S. Brown and W. I. Weisberger, later by E. Eichten and F. Feinberg in [7, 8]. A one loop RG improved calculation was done in [11]. In those papers an expansion is performed around the static solution. Here, we adopt the approach of [9] where the kinetic energy was kept during all the calculations.

The 4-point gauge invariant Green function $G$ associated with the Lagrangian
is defined as
\[ G(x_1, y_1, x_2, y_2) = \langle 0 | Q_2^\dagger(x_2) \phi(x_2, x_1) Q_1(x_1) Q_1^\dagger(y_1) \phi(y_1, y_2) Q_2(y_2) | 0 \rangle, \]
where \( \phi(x_2, x_1) \equiv \exp \left\{ -ig \int_0^1 ds (x_2 - x_1)^\mu A_\mu(x_1 + s(x_2 - x_1)) \right\} \) is a Schwinger line added to select the singlet state contribution. For large time separations the string vanishes. After integrating out the heavy quark fields \( Q_j \) and \( Q_j^\dagger \), \( G \) can be expressed as a quantum-mechanical path integral over the quark trajectories

\[
G = \int_{y_1}^{x_2} \mathcal{D}z_1 \mathcal{D}p_1 \int_{y_2}^{x_2} \mathcal{D}z_2 \mathcal{D}p_2 \exp \left\{ i \int_{-T/2}^{T/2} dt \sum_{j=1}^{2} p_j^i \cdot z_j - \frac{p_j^2}{2m_j} + \frac{p_j^4}{8m_j^3} - i \int_{-T/2}^{T/2} dt \mathcal{U} \right\}
\]
where \( z_j = z_j(t) \) and \( y_2^0 = y_1^0 \equiv -T/2, x_2^0 = x_1^0 \equiv T/2 \). The function \( \mathcal{U} \) describes the heavy quark interaction. Since the kinetic energy has been properly isolated, it is consistent with the counting rules given in the previous section to expand the interaction around the static quark trajectories \( z_1 = (t, \mathbf{r}) \) and \( z_2 = (t, 0) \). Moreover, we define the heavy quark-antiquark potential as \( V = \lim_{T \to \infty} \frac{1}{T} \int_{-T/2}^{T/2} dt \mathcal{U} \), with the warning that, having soft and ultrasoft degrees of freedom not been disentangled in NRQCD, non potential terms could still contribute to some extent to \( V \).

Working out this expression \( \mathfrak{a} \) \( \mathfrak{b} \) we get

\[
V = V_0 + \text{spin independent terms} - \frac{1}{m_1 m_2} (d_{sv} + C_F d_{vw}) \delta^3(r)
\]

\[
+ \frac{1}{8} \left( \frac{c_{D}^{(1)}}{m_1^2} + \frac{c_{D}^{(2)}}{m_2^2} \right) (\Delta V_0(r) + \Delta V_a^E(r)) + \frac{1}{8} \left( \frac{c_F^{(1)}}{m_1^2} + \frac{c_F^{(2)}}{m_2^2} \right) \Delta V_a^B(r)
\]

\[
+ \left( \frac{S^{(1)} \cdot L^{(1)}}{m_1^2} + \frac{S^{(2)} \cdot L^{(2)}}{m_2^2} \right) 2c_F V_1(r) + \frac{c_S V_0(r)}{2} + \frac{S^{(1)} \cdot L^{(2)} + S^{(2)} \cdot L^{(1)}}{m_1 m_2} \frac{c_F V_2(r)}{r}
\]

\[
+ \frac{c_F^{(1)} c_F^{(2)}}{m_1 m_2} \left( \frac{S^{(1)} \cdot r S^{(2)} \cdot r}{r^2} - \frac{S^{(1)} \cdot S^{(2)}}{3} \right) V_3(r)
\]

\[
+ \frac{S^{(1)} \cdot S^{(2)}}{3 m_1 m_2} \left( c_F^{(1)} c_F^{(2)} V_4(r) - 12 (d_{sv} + C_F d_{vw}) \delta^3(r) \right).
\]

\[ S^{(j)} = \mathbf{\sigma}^{(j)}/2 \text{ and } L^{(j)} = \mathbf{r} \times \mathbf{p}_j \text{ are the spin and orbital angular momentum operators respectively.} \]

The matching coefficients \( d_\ell \) come from the 4-fermion sector \( L_{4f} \) \( \mathfrak{c} \) and \( 2c_{F,S}^{\pm} \equiv c_{F,S}^{(1)} \pm c_{F,S}^{(2)} \). The static potential \( V_0 \) is

\[
V_0(r) = \lim_{T \to \infty} \frac{i}{T} \ln \langle W(\Gamma_0) \rangle, \quad W(\Gamma) = \mathbb{P} \exp \left\{ -ig \oint dz^\mu A_\mu(z) \right\}.
\]
where $\Gamma$ is the loop made up by the quark trajectories $z_1$ and $z_2$ and the endpoints Schwinger strings, and the static loop $\Gamma_0$ is a $r \times T$ rectangle. The bracket means the color trace of the average over the gauge fields weighted by the gluon Lagrangian $L_g$. The “potentials” $V_1, V_2, \ldots$ are scale dependent gauge field averages of electric and magnetic field strength insertions in the static Wilson loop, see [14, 15].

Since the spin independent corrections come from the terms $D_0$ and $D^2/2m$ of the Lagrangian (3) whose matching coefficients are protected by reparameterization invariance, they are scale independent. An evaluation of the so-called momentum dependent corrections can be found in [3]. Terms involving logarithms of the quark masses are present. As a consequence of the matching procedure they are all encoded in the matching coefficients. Therefore, as argued in [14], the correct handling of the matching allows finally the agreement between the potential derived here evaluated in the perturbative regime with the traditional QCD one loop perturbative calculation [3, 21]. For a careful discussion of this point see [19].

The “potentials” are known exactly (up to a given order in $\alpha_s$) only in the perturbative regime, i.e. in the short range behaviour. Nevertheless there exist some exact relations between them which hold at any range. For instance, from Lorentz invariance it follows that

\[ V'_0(r) + V'_1(r) - V'_2(r) = 0. \]

Since for reparameterization invariance $c_S^{(j)} = 2c_F^{(j)} - 1$ [17], equation (5) holds at any scale $\mu$. Similar relations exist for the momentum dependent “potentials” of Ref. [4]. Moreover the scale independence of the potential (3) (i.e. $\frac{dV}{d\mu} = 0$) establishes several relations between quantities at different renormalization scales [14].

The potentials $V_0, V_1, \ldots$ are suitable for direct lattice computation and for analytic evaluation inside a QCD vacuum model. The first possibility relies on the fact that all the dynamical quantities are expressed in terms of field strength insertions on a static Wilson loop which is an object traditionally measured on the lattice. Such an analysis has been performed in [14]. All the “potentials” have been measured on the lattice and the previous mentioned exact relations have been used to check the accuracy of the results. Using the parameterized form obtained in this way for the “potentials” the bound state equation has been solved and the bottomonium and charmonium spectra evaluated. The agreement with the experimental data is found to be quite good. On the other hand, an analytic evaluation of the heavy quark potential inside a QCD vacuum model turns out to be very convenient in this framework since by means of functional derivatives all the averages of field strength insertions on the Wilson loop can be expressed in terms of the average of the non-static Wilson loop alone. This is usually a quantity provided by QCD vacuum models. A study of different models
in this framework has been done in \cite{15, 22} where a comparison with the existing lattice data has also been provided. Up to now these are not accurate enough to really discriminate between different models. More precise lattice measurements will be performed in the near future providing in this way strong constraints on all the infrared QCD models with predictability on the long-range quarkonium interaction.

Finally we emphasize that until soft and ultrasoft degrees of freedom have not been disentangled an exact value in $v$ cannot be assigned to each term of the effective Lagrangian and the power counting has to be interpreted at the leading order. For instance the $O(1)$ NRQCD Lagrangian, $L = Q_1^\dagger (iD_0 + \partial^2 / 2m_1) Q_1 + Q_2^\dagger (-iD_0 + \partial^2 / 2m_2) Q_2$, does not contribute to Eq. (3) only with the static potential $V_0$. Since the corresponding Wilson loop $P \exp \left\{-ig \oint_\Gamma dz^0 A_0 (z) \right\}$ is a function of the non-static loop $\Gamma$, its expansion produces spin independent terms of order $O(v)$ as well.

4 Comments and outlook

In the framework of NRQCD we have shown how to get the complete order $O(v^4)$ expression of the heavy quark-antiquark potential in terms of field strength insertions on a static Wilson loop. This has several advantages. Nonperturbative contributions can be evaluated by means either of traditional lattice calculations or of different QCD vacuum models. Having worked out the matching procedure, we find that the potential of Eq. (3) is consistent in the short range with the existing perturbative calculations and with the lattice data \cite{13}. These are sensitive to one loop and in some cases to the next to leading correction too. We stress here that terms proportional to the static potential in Eq. (3) have to be protected from the running. For instance, using Eq. (3) and reparameterization invariance, the spin-orbit interaction term $\sim 2C_F^+ V_1' + C_F^+ V_0'$ can be written as $2C_F^+ V_2' - V_0'$ where it is apparent that no Wilson coefficient multiplies the static potential. Moreover we can draw the following consequence about $V_2'$: either $V_2'$ does not contain any nonperturbative contribution at all (as present lattice data seem to suggest) or, if $V_2'$ contains some nonperturbative contributions, they have to satisfy the RG equation $\frac{d}{d\mu} C_F^+ V_2' = 0$. The same argument suggests for the Darwin term that $\Delta V_a^E = -\Delta V_0 +$ perturbative contributions. This gives some constraints on the QCD vacuum models which, introducing at some point some approximations, lose scale invariance.

All the corrections to the effective Lagrangian discussed here are relevant in order to obtain the quarkonium spectrum with an accuracy of $O(v^4)$. For some Wilson coefficients only the leading $\alpha_s$ contribution needs to be taken into account. This is no more the case if we aim to reach a 10% accuracy in the quarkonium spin splittings. Being these quantities an order $O(v^4)$ effect, a 10%
accuracy is achievable only if $O(v^6)$ and $O(\alpha_s v^4)$ effects are calculated as well. Therefore operators of order $O(v^6)$ should be added to the effective Lagrangian \[5\]. The inclusion of such operators in Eq. \(3\) in terms of (two and three) field strength insertions on a static Wilson loop, is only a technical problem, but has not been done so far. The main reason is that non potential contributions are expected to become even more important. Ultrasoft gluonic degrees of freedom, characterized by a time scale $1/mv^2$, exist. In perturbative QCD they are responsible, for instance, of the Lamb-shift (which is a $\alpha_s v^4$ effect). In the language of the Wilson loop, this means that the interaction $U$ (see Sec. 3) can be affected by non potential contributions which have to be subtracted from $V$ in order to define properly a heavy quark potential. Since $\alpha_s$ is not a small parameter at the ultrasoft scale, these non potential terms could be an effect of order $O(v^4)$. Therefore, precision studies of quarkonium have to take into account it. An approach was recently proposed in \[12, 13\]. The ultrasoft degrees of freedom are integrated out directly from the NRQCD Lagrangian giving rise to another effective theory, called potential NRQCD, where all the energy scales of the bound state are disentangled explicitly. The advantages are enormous, since in the new theory potential and non potential contributions are clearly separated. The novel feature is that the matching this time takes place in a energy region dominated by the nonperturbative physics.

**Acknowledgements**

It is a pleasure for N. B. to thank Andrei Afanesev, Nathan Isgur, Dubravko Klabucar, Elio Soldi and Branko Vlahovic for the organization of this nice Conference. N. B. acknowledges financial support from the European Community, Marie Curie fellowship-TMR Contract n. ERBFMBICT961714, and from the theory section of Milano.

**References**

[1] W. Lucha, F. F. Schöberl and D. Gromes, Phys. Rep. **200**, 127 (1991).

[2] Quarkonia ed. by W. Buchmüller (North-Holland, Amsterdam, 1992).

[3] S. Titard and F. J. Yndurain, Phys. Rev. **D 49**, 6007 (1994); A. Pineda and F. J. Yndurain, Phys. Rev. **D 58**: 094022 (1998); F. J. Yndurain, Nucl. Phys. Proc. Suppl. **B 64**, 433 (1998).

[4] C. Davies, in “Schladming 1997, Computing Particle Properties”, (Springer, Berlin, 1997).
[5] G. P. Lepage, L. Magnea, C. Nakhleh, U. Magnea and K. Hornbostel, Phys. Rev. D 46, 4052 (1992).

[6] K. G. Wilson, Phys. Rev. D 10, 2445 (1974).

[7] L. S. Brown and W. I. Weisberger, Phys. Rev. D 20, 3239 (1979).

[8] E. Eichten and F. Feinberg, Phys. Rev. D 23, 2724 (1981); M. E. Peskin in Proceeding of the 11th SLAC Institute, SLAC Report No. 207, 151, edited by P. Mc Donough (1983); D. Gromes, Zeit. Phys. C 22, 265 (1984).

[9] N. Brambilla, P. Consoli, G. Prosperi, Phys. Rev. D 50, 5878 (1994); A. Barchielli, N. Brambilla and G. Prosperi, Nuovo Cimento 103 A, 59 (1990); A. Barchielli, E. Montaldi and G. Prosperi, Nucl. Phys. B 296, 625 (1988).

[10] N. Brambilla and A. Vairo, hep-ph/9809460, Phys. Rev. D 56, 1445 (1997).

[11] Y. Chen, Y. Kuang and R. J. Oakes, Phys. Rev. D 52, 264 (1995).

[12] N. Brambilla, A. Pineda, J. Soto and A. Vairo, in preparation.

[13] A. Pineda and J. Soto, Nucl. Phys. Proc. Suppl. B 64, 428 (1998); J. Soto in the Proceedings of Confinement III, Jefferson Lab (1998).

[14] G. S. Bali, K. Schilling and A. Wachter, Phys. Rev. D 56, 2566 (1997); G. S. Bali and P. Boyle, hep-ph/9809180.

[15] N. Brambilla, hep-ph/9809263, in the Proceedings of Confinement III, Jefferson Lab (1998).

[16] M. Neubert, Phys. Rep. 245, 259 (1994).

[17] A. V. Manohar, Phys. Rev. D 56, 230 (1997).

[18] A. Pineda and J. Soto, Phys. Rev. D 58: 114011 (1998).

[19] N. Brambilla and A. Vairo, in preparation: hep-ph/9809230; A. Vairo, hep-ph/9809229.

[20] C. Quigg and J. L. Rosner, Phys. Rep. 56 C, 167 (1979).

[21] Y. J. Ng, J. Pantaleone and S. H. Tye, Phys. Rev. Lett. 55, 916 (1985); Phys. Rev D 33, 777 (1986).

[22] N. Brambilla and A. Vairo, Phys. Rev. D 55, 3974 (1997).