The $q\bar{q}$ semirelativistic interaction in the Wilson loop approach

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The complete $q\bar{q}$ semirelativistic interaction is obtained as a gauge-invariant function of the Wilson loop and its functional derivatives. The approach is suitable for analytic evaluations as well as for lattice calculations. Here we consider three different models for the Wilson loop non-perturbative behaviour and discuss the related semirelativistic dynamics.

1. INTRODUCTION

The derivation of the $q\bar{q}$ potential is a longstanding problem; in Ref. [2] the static and the spin-dependent potentials were first obtained. Recently, working with a path integral representation for the quark propagator in the external field, the complete $1/m^2$ $q\bar{q}$ potential was obtained as a gauge-invariant function of the Wilson loop and its functional derivatives [3]:

$$\int_{t_i}^{t_f} dt V_{q\bar{q}} = i \log \langle W(\Gamma) \rangle$$

$$- \sum_{j=1}^{2} \frac{g}{m_j} \int_{\Gamma_j} dx^\mu \left( S^l_j \langle \hat{F}_{l\mu}(x) \rangle - \frac{1}{8m_j} \langle D^{l'} F_{\nu l\mu}(x) \rangle \right)$$

$$- \frac{1}{2} \sum_{j,j'=1}^{m} \frac{ig^2}{m_j m_{j'}} T_s \int_{\Gamma_j} dx^\mu \int_{\Gamma_{j'}} dx^\nu S^l_j S^l_{j'}$$

$$\times \left( \langle \hat{F}_{l\mu}(x) \hat{F}_{l\nu}(x') \rangle - \langle \hat{F}_{l\mu}(x) \rangle \langle \hat{F}_{l\nu}(x') \rangle \right)$$

where

$$W(\Gamma) \equiv P \exp \left[ ig \oint_{\Gamma} dx^\mu A_\mu(x) \right].$$

$$\langle \rangle$$ is the normalized average over the gauge fields $A_\mu$ and

$$\langle f(A) \rangle \equiv \langle f(A) W(\Gamma) \rangle / \langle f(A) W(\Gamma) \rangle.$$ 

The closed loop $\Gamma$ is defined by the quark (antiquark) trajectories $z_1(t)$ ($z_2(t)$) running from $y_1$ to $x_1$ ($x_2$ to $y_2$) along with two straight-lines at fixed time connecting $y_1$ to $y_2$ and $x_1$ to $x_2$. Moreover we have

$$g \langle F_{\mu\nu}(z_j) \rangle = (-1)^{j+1} \frac{\delta \log \langle W(\Gamma) \rangle}{\delta S^{\mu\nu}(z_j)}, \quad (3)$$

$$g^2 \left( \langle F_{\mu\nu}(z_1) F_{\nu\rho}(z_2) \rangle - \langle F_{\mu\nu}(z_1) \rangle \langle F_{\nu\rho}(z_2) \rangle \right)$$

$$= -ig \frac{\delta}{\delta S^{\mu\nu}(z_2)} \langle F_{\mu\nu}(z_1) \rangle, \quad (4)$$

where $\delta S^{\mu\nu}(z_j) = dz^\mu_j dz^\nu_j - dz^\nu_j dz^\mu_j$. As a consequence we have that the actual form of the potential in [3] is completely known once the Wilson loop behaviour is given.

The semirelativistic $q\bar{q}$ potential can be written as

$$V_{q\bar{q}} = V_0 + V_{VD} + V_{SD}, \quad (5)$$

with

$$i \log \langle W(\Gamma) \rangle = \int_{t_i}^{t_f} dt V_0(r(t)) + V_{VD}(r(t)),$$

$$V_{VD}(r(t)) = \frac{1}{m_1 m_2} \left\{ \frac{1}{3} \mathbf{p}_1 \cdot \mathbf{p}_2 \left( \frac{\mathbf{p}_1 \cdot \mathbf{r}}{r^2} \cdot \mathbf{r} \right) V_e(r) \right\}_w$$

$$+ \frac{2}{m_j^3} \left\{ \left( \mathbf{p}_j \cdot \mathbf{r} \right) V_d(r) + \left( \frac{1}{3} \mathbf{p}_j \cdot \mathbf{r} \right) V_e(r) \right\}_w, \quad (6)$$

where $r(t) \equiv z_1(t) - z_2(t)$ and the symbol $\{ \}$ \textit{w} stands for the Weyl ordering prescription. The
spin dependent interaction appearing in \( \text{[1]} \) has a transparent physical meaning the first term being the magnetic interaction, the second the Thomas precession, the third the Darwin term and the last one the spin-spin interaction. However, it can be rewritten in the usual Eichten–Feinberg form \( \text{[2]} \):

\[
V_{SD} = \frac{1}{8} \left( \frac{1}{m_1^2} + \frac{1}{m_2^2} \right) \Delta [V_0(r) + V_a(r)] + \sum_{j=1,2} \left( \frac{(-1)^{j+1}}{2m_j^2} \mathbf{L}_j \cdot \mathbf{S}_j \right) \frac{1}{r} \frac{d}{dr} [V_0(r) + 2V_1(r)]
\]

\[
+ \frac{1}{m_1 m_2} \left( \mathbf{L}_1 \cdot \mathbf{S}_2 - \mathbf{L}_2 \cdot \mathbf{S}_1 \right) \frac{1}{r} \frac{d}{dr} V_2(r)
\]

\[
+ \frac{1}{3m_1 m_2} \mathbf{S}_1 \cdot \mathbf{S}_2 V_3(r).
\]

It is clear that all the dynamics is contained in the \( V_i(r) \) of Eqs. \( \text{[1]} \)–\( \text{[3]} \) and these are unambiguous functions of the Wilson loop (cf. \( \text{[4]} \)). Similar expressions for the \( V_i \) in terms of insertions on the static Wilson loop were obtained in \( \text{[3]} \). These expressions are suitable for lattice evaluation and were used in \( \text{[4]} \) to obtain lattice predictions for the semirelativistic interaction. However, this is also the ideal framework to formulate hypothesis on the Wilson loop behaviour (and so on the confinement mechanism) to be checked on the lattice and on the experimental data. In the following we use three different models for the QCD vacuum to obtain an analytic behaviour of the Wilson loop in the non-perturbative region and to predict the semirelativistic quark dynamics. Comparison among the results allows us to get some insight in the mechanism of confinement \( \text{[5]} \).

2. THE CONFINING SEMIRELATIVISTIC DYNAMICS

2.1. Minimal area law model (MAL)

In Ref. \( \text{[3,5]} \) \( \langle W(\Gamma) \rangle \) was approximated by the sum of a perturbative part given at the leading order by the gluon propagator \( D_{\mu \nu} \) and a non-perturbative part given by the value of the minimal area of the deformed Wilson loop of fixed contour \( \Gamma \) plus a perimeter contribution \( \mathcal{P} \):

\[
i \log \langle W(\Gamma) \rangle = -\frac{4}{3} g^2 \oint_{\Gamma} dx_1^x \oint_{\Gamma} dx_2^x 
\times i D_{\mu \nu}(x_1 - x_2) + \alpha S_{\min} + \frac{C}{2} \mathcal{P}.
\]

Denoting by \( u^\mu = u^\mu(s, t) \) the equation of any surface with contour \( \Gamma (s \in [0,1], t \in [t_1, t_2]) \), \( u^0(s, t) = t, \mathbf{u}(1, t) = \mathbf{z}_1(t), \mathbf{u}(0, t) = \mathbf{z}_2(t) \) we have:

\[
S_{\min} = \min \int_{t_1}^{t_2} dt \int_0^1 ds \times \left[ \left( \frac{\partial u^\mu}{\partial t} \right)^2 - \left( \frac{\partial u^\mu}{\partial s} \right)^2 \right]^{1/2}
\]

which coincides with the Nambu–Goto action. Up to the order \( 1/m^2 \) the minimal surface can be identified exactly with the surface spanned by the straight-line joining \( (t, \mathbf{z}_1(t)) \) to \( (t, \mathbf{z}_2(t)) \). From Eq. \( \text{[3]} \) \( V_{q\bar{q}} \) is obtained \( \text{[3,4]} \). In particular for the static potential we have

\[
V_0 = -\frac{4 \alpha_s}{3} \frac{r}{\pi} + \sigma r + C.
\]

See Tab. 1 for the complete results.

2.2. Stochastic vacuum model

Using the non-Abelian Stokes theorem and the cumulant expansion it is possible to write

\[
\langle W(\Gamma) \rangle = \exp \left( ig \int_S dS_{\mu \nu}(u) F_{\mu \nu}(u, x_0) \right)
\]

\[
= \exp \sum_{j=1}^{\infty} \frac{(ig)^j}{j!} \int_S dS_{\mu_1 \nu_1}(u_1) \cdots \int_S dS_{\mu_j \nu_j}(u_j)
\times \langle F_{\mu_j \nu_j}(u_1, x_0) \cdots F_{\mu_1 \nu_1}(u_j, x_0) \rangle_{\text{cum}}
\]

with \( \langle \cdots \rangle_{\text{cum}} \) defined in terms of average values over the gauge fields \( \langle \cdots \rangle \) and \( F_{\mu \nu}(u, x_0) \) being the path-ordered product of the field strength tensor \( F_{\mu \nu}(u) \) times two Schwinger strings connecting the point \( u \) with an arbitrary reference point \( x_0 \) on the surface \( S \) appearing in the non-Abelian Stokes theorem.

Equation \( \text{[10]} \) is exact. The first cumulant vanishes trivially. The second cumulant gives the
The first non-zero contribution to the cluster expansion. In the SVM, one assumes that in the context of heavy quark bound states higher cumulants can be neglected and the second cumulant dominates the cluster expansion, i.e., that the vacuum fluctuations are of a Gaussian type:

\[
\log\langle W(\Gamma) \rangle = -\frac{g^2}{2} \int_S dS_{\mu}(u) \int_S dS_{\lambda}(v) \times (F_{\mu \nu}(u, x_0) F_{\lambda \rho}(v, x_0))_{\text{cum}},
\]

(11)

with

\[
(F_{\mu \nu}(u, x_0) F_{\lambda \rho}(v, x_0))_{\text{cum}} = \frac{\beta}{g^2} \left\{ (\delta_{\mu \lambda} \delta_{\nu \rho} - \delta_{\mu \rho} \delta_{\nu \lambda}) D\left((u - v)^2\right) + \frac{1}{2} \left[ \frac{\partial}{\partial u_\mu} ((u - v) \lambda \delta_{\nu \rho} - (u - v) \rho \delta_{\nu \lambda}) + \frac{\partial}{\partial u_\nu} ((u - v) \mu \delta_{\lambda \rho} - (u - v) \rho \delta_{\lambda \mu}) \right] \times D_1((u - v)^2) \right\}
\]

(12)

\[
\beta \equiv \frac{g^2 \langle \text{Tr} F_{\mu \nu}(0) F_{\mu \nu}(0) \rangle}{36 D(0) + D_1(0)}.
\]

Eqs. (11) and (12) define the SVM for heavy quarks. The correlator functions \(D\) and \(D_1\) are unknown. The perturbative part of \(D_1\), which is expected to be dominant in the short-range behaviour, can be obtained by means of the standard perturbation theory:

\[
D_1^{\text{pert}}(x^2) = \frac{160 \pi}{3\pi} \frac{1}{x^4} + \text{higher orders}.
\]

(13)

One of the main features of SVM is to insert the information coming from the lattice inside the analytic model. Indeed, the non-perturbative behaviour of \(D\) and \(D_1\) was evaluated on the lattice (see [7])

\[
\beta D_1^{LR}(x^2) = d e^{-\delta|x|}, \beta D_1^{LR}(x^2) = d_1 e^{-\delta_1|x|}
\]

(14)

with \(\delta = (1 \pm 0.1) \text{ GeV}, d = 0.073 \text{ GeV}^4, \delta_1 = (1 \pm 0.1) \text{ GeV}, d_1 = 0.0254 \text{ GeV}^4\). From Eqs. (14) all the potentials follow directly. In particular the static potential is given by

\[
V_0(r) = \beta \int_{-\infty}^{+\infty} d\tau \left\{ \int_0^\tau d\lambda (r - \lambda) \times D(r^2 + \lambda^2) + \int_0^\tau d\lambda \frac{\lambda}{2} D_1(r^2 + \lambda^2) \right\},
\]

(15)

In the \(r \to 0\) limit (13) reproduces one gluon exchange while in the limit \(r \to \infty\)

\[
V_0(r) = \sigma_2 r + \frac{1}{2} C_2^{(1)} - C_2
\]

\[
\sigma_2 = \beta \int_{-\infty}^{+\infty} d\tau \int_0^\infty d\lambda \ D_2(r^2 + \lambda^2),
\]

\[
C_2 = \beta \int_{-\infty}^{+\infty} d\tau \int_0^\infty d\lambda \ D(r^2 + \lambda^2),
\]

\[
C_2^{(1)} = \beta \int_{-\infty}^{+\infty} d\tau \int_0^\infty d\lambda \ D_1(r^2 + \lambda^2).
\]

It is clear that (13) contains the MAL result for the static potential upon identification of the string tension and of the constants. For the other potentials see Ref. [3] (complete form) and Tab. 1 (\(r \to \infty\) limit).

2.3. Dual QCD

The duality assumption that the long distance physics of a Yang–MILLS theory depending upon strong coupled gauge potentials \(A_\mu\) is the same as the long distance physics of the dual theory describing the interactions of weakly coupled dual potentials \(C_\mu \equiv \sum_{a=1}^{s} C^a_\mu \lambda_a/2\) and monopole fields \(B_i \equiv \sum_{a=1}^{s} B_i^a \lambda_a/2\), forms the basis of DQCD [8]. The model is constructed as a concrete realization of the Mandelstam–’tHooft dual superconductor mechanism of confinement. Since the main interest is solving such a theory in the long-distance regime, the Lagrangian \(\mathcal{L}_{\text{eff}}\) is explicitly constructed as the minimal dual gauge invariant extension of a quadratic Lagrangian with the further requisite to give a mass to the dual gluons (and to the monopole fields) via a spontaneous symmetry breaking of the dual gauge group.

We denote by \(\langle W_{\text{eff}}(\Gamma) \rangle\) the average over the fields of the Wilson loop of the dual theory [8]:

\[
\langle W_{\text{eff}}(\Gamma) \rangle = \frac{\int DC DB e^{i \int dz [\mathcal{L}_{\text{eff}}(G_{\mu \nu}^s) + \mathcal{L}_{\text{GP}}]} \int DC DB e^{i \int dz [\mathcal{L}_{\text{eff}}(G_{\mu \nu}^{\text{new}} = 0) \mathcal{L}_{\text{GP}}]}}{\int DC DB e^{i \int dz [\mathcal{L}_{\text{eff}}(G_{\mu \nu}^s) + \mathcal{L}_{\text{GP}}]}}.
\]
$\mathcal{L}_{\text{GF}}$ is a gauge fixing term and the effective Lagrangian in presence of quarks is given by

$$\mathcal{L}_{\text{eff}} = 2 \text{Tr} \left\{ -\frac{1}{4} G^{\mu\nu} G_{\mu\nu} + \frac{1}{2} (D_\mu B_i)^2 \right\} - U(B_i).$$

$U(B_i)$ is the Higgs potential with a minimum at a non-zero value $B_{01} = B_0 \lambda_7$, $B_{02} = -B_0 \lambda_5$ and $B_{03} = B_0 \lambda_2$. Moreover we have taken $B_1 = B_2 = B$, the dual potential proportional to the hypercharge matrix $C_i = C^\mu Y$ and

$$D_\mu B_i = \partial_\mu B_i + i e [C_\mu, B_i], \quad e \equiv \frac{2\pi}{g},$$

$$G_{\mu\nu} = (\partial_\mu C_\nu - \partial_\nu C_\mu + G^{S}_{\mu\nu}) Y,$$

$$G^S_{\mu\nu}(x) = g \varepsilon_{\mu\nu\alpha\beta} \int ds \int d\tau \left( \frac{\partial y^\alpha}{\partial s} \frac{\partial y^\beta}{\partial \tau} \right) \delta(x - y(s, \tau)),$$

where $y(s, \tau)$ is a world sheet with boundary $\Gamma$ swept out by the Dirac string.

The assumption that the dual theory describes the long distance $q\bar{q}$ interaction in QCD then takes the form:

$$\langle W(\Gamma) \rangle = \langle W_{\text{eff}}(\Gamma) \rangle, \quad \text{for large loops } \Gamma. \quad (16)$$

Large loop means that the size $R$ of the loop is large compared to the inverse mass ($M_R^{-1} \simeq (600 \text{ MeV})^{-1}$) of the Higgs particle (monopole field). Furthermore, since the dual theory is weakly coupled at large distances, we can evaluate $\langle W_{\text{eff}}(\Gamma) \rangle$ via a semiclassical expansion to which the classical configuration of the dual potentials and monopoles gives the leading contribution. In the leading classical approximation

$$i \log \langle W_{\text{eff}}(\Gamma) \rangle = - \int dx \mathcal{L}_{\text{eff}}(G^S_{\mu\nu}), \quad (17)$$

with $\mathcal{L}_{\text{eff}}(G^S_{\mu\nu})$ evaluated at the solution of the classical equations of motion. For the static potential we have

$$V_0 = -\frac{4}{3} \frac{\alpha_s}{r} \exp \left( -0.511 \sqrt{\frac{\sigma}{\alpha_s}} r \right) + \sigma r - 0.646 \sqrt{\sigma \alpha_s}.$$ 

This results directly from the classical solution to the nonlinear equations obtained from $\mathcal{L}_{\text{eff}}$ and reproduces the one gluon exchange in the limit $r \to 0$ and the string term in the limit $r \to \infty$. The coefficient of the exponent can be actually given in terms of the dual gluon mass $M = 6 g^2 B_0^2 \simeq \frac{4}{g^2}$. An interpolation of the numerical results for the potentials can be found in [33]. In Tab.1 we report the potentials in $r \to \infty$ limit.

3. FLUX TUBE STRUCTURE AND CONFINEMENT MECHANISM

To discuss the semirelativistic non-perturbative dynamics is convenient to study the large $r$ limit in the three above models as reported in Tab.1. The MAL result is the realization of the intuitive Buchm"uller’s picture of zero magnetic field in the flux tube comoving system: $dV_2/dr$ which is given by the magnetic interaction is zero in the non-perturbative region; similarly the velocity dependent potentials $V_3$–$V_6$ come from the consideration of the flux–tube energy. From Tab.1 it is apparent that at the leading order in the long–range limit, neglecting exponentially falling off terms, the SVM contains exactly the MAL model results. In the spin dependent sector of the potential, both the SVM and DQCD not only reproduce the long-range behaviour given by the area law, but also give $1/r$ corrections to $dV_1/dr$ and $dV_2/dr$. These corrections are equal in both models and very near to the absolute value of the constant term in the static potential (the SVM also supplies for the explanation of this fact). This perfect agreement is absolutely not trivial and seems to be very meaningful, since it arises from two very different models in a region of distances in which the physics cannot be described by the area law alone. It is now clear that the vanishing of the magnetic part in the non-perturbative region takes place only at the leading order in the long-range limit. Therefore, working in a Bethe–Salpeter context, there is no need to assume an effective pure convolution kernel which is a Lorentz scalar. Velocity dependent contributions to the quark-antiquark potential are important. In fact the string behaviour of the non-perturbative interaction shows up when we consider the velocity dependent part of the potential and this is also what the data
Table 1
Complete MAL potential and long distance SVM and DQCD potentials; from parameterization \( \frac{d}{a} \delta \) we have: \( \sigma_2 = \frac{\pi d}{3r} \), \( C_2 = \frac{4d}{3r} \); \( C_2^{(1)} = \frac{4d}{3r} \); \( E_2 = \frac{3\pi d}{3r} \); \( E_2 = \frac{3\pi d}{3r} \); see the values of \( d \) and \( \delta \) given in the text. The dual gluon mass \( M^2 = \frac{\pi \sigma}{4\alpha_s} \); typical values are \( \sigma = 0.18; \alpha_s = 0.35 \).

| \( V_0 \) | \( \frac{-4}{3} \frac{\alpha_s}{r} + \sigma r + C \) | \( \sigma_2 r + \frac{1}{2} C_2^{(1)} - C_2 \) | \( \sigma r - 0.646 \sqrt{\sigma \alpha_s} \) |
| \( \Delta V_a \) | 0 | const. self-energy terms | const. terms \(-2 \frac{\sigma}{r}\) |
| \( dV_1/dr \) | \(-\sigma\) | \(-\sigma_2 + \frac{C_2}{r}\) | \(-\sigma + \frac{0.681 \sqrt{\sigma \alpha_s}}{r}\) |
| \( dV_2/dr \) | \( \frac{4}{3} \frac{\alpha_s}{r^2} \) | \( \frac{4}{3} \frac{\alpha_s}{r^2} \) | \( \frac{4}{3} \alpha_s (M^2 + \frac{3}{3} M + \frac{3}{3} \epsilon) e^{M_r} \) |
| \( V_3 \) | \( 32 \pi \alpha_s \delta^3 (r) \) | exp. fall off | \( \frac{4}{3} \alpha_s M^2 e^{M_r} \) |
| \( V_6 \) | \( \frac{8}{9} \frac{\alpha_s}{r} - \frac{1}{9} \sigma r \) | \( \frac{1}{9} \sigma_2 r - \frac{2}{3} \frac{D_2}{r} + \frac{8}{9} \frac{E_2}{r} \) | \(-0.097 \sigma r - 0.226 \sqrt{\sigma \alpha_s} \) |
| \( V_c \) | \( \frac{-2}{3} \frac{\alpha_s}{r} - \frac{1}{6} \sigma r \) | \( \frac{-1}{6} \sigma_2 r - \frac{2}{3} \frac{D_2}{r} + \frac{3}{3} \frac{E_2}{r} \) | \(-0.146 \sigma r - 0.516 \sqrt{\sigma \alpha_s} \) |
| \( V_d \) | \( \frac{-1}{6} \sigma r - \frac{1}{4} C \) | \( \frac{-1}{3} \sigma_2 r + \frac{1}{4} C_2 - \frac{1}{8} C_2^{(1)} + \frac{1}{4} \frac{D_2}{r} - \frac{2}{3} \frac{E_2}{r} \) | \(-0.118 \sigma r + 0.275 \sqrt{\sigma \alpha_s} \) |
| \( V_c \) | \( \frac{-1}{6} \sigma r \) | \( \frac{-1}{6} \sigma_2 r + \frac{1}{4} C_2 - \frac{1}{8} C_2^{(1)} + \frac{1}{4} \frac{D_2}{r} - \frac{2}{3} \frac{E_2}{r} \) | \(-0.177 \sigma r + 0.258 \sqrt{\sigma \alpha_s} \) |

require. The velocity dependent structure which arises from the DQCD model differs slightly in the coefficients with respect to the area law behaviour. In conclusion SVM and DQCD reproduce the flux tube distribution measured on the lattice and the general features coming from the area law. Both give analytical expressions for the Wilson loop which describe the evolving behaviour of \( \langle W (\Gamma) \rangle \) from the short to the long distances but not all predictions are equal in the two models in the intermediate distances region, in particular in the velocity dependent sector of the potential, and also in the spin-spin interaction as well as in \( \Delta V_a \). The up to now available lattice data \( \delta \) confirm the MAL results but are not sufficiently accurate to discriminate between SVM and DQCD models.

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