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Relativistic properties of the quark-antiquark potential

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Abstract. The relativistic transformation properties of the heavy quark-antiquark interaction potential are considered in the framework of the relativistic quark model. A special attention is paid to the long-range (confining) contribution to the spin-independent part of $q\bar{q}$ interaction. The retardation effects are consistently taken into account.

The relativistic properties of the quark-antiquark interaction potential play an important role in analysing different static and dynamical characteristics of heavy mesons. The Lorentz-structure of the confining quark-antiquark interaction is of particular interest. In the literature there is no consent on this item. For a long time the scalar confining kernel has been considered to be the most appropriate one [1,2]. The main argument in favour of this choice is based on the nature of the heavy quark spin-orbit potential. The scalar potential gives a vanishing long-range magnetic contribution, which is in agreement with the flux tube picture of quark confinement of [3], and allows to get the fine structure for heavy quarkonia in accord with experimental data. However, the calculations of electroweak decay rates of heavy mesons with a scalar confining potential alone yield results which are in worse agreement with data than for a vector potential [4,5]. The radiative $M1$-transitions in quarkonia such as e. g. $J/\psi \to \eta_c\gamma$ are the most sensitive for the Lorentz-structure of the confining potential. The relativistic corrections for these decays arising from vector and scalar potentials have different signs [4,5]. In particular, as it has been shown in [5], agreement with experiments for these decays can be achieved only for a mixture of vector and scalar potentials. In this context, it is worth remarking, that the recent study of the $q\bar{q}$ interaction in the Wilson loop approach [6] indicates that it cannot be considered as simply a scalar. Moreover, the found structure of spin-independent relativistic corrections is not compatible with a scalar potential. A similar conclusion has been obtained in [7] on the basis of a Foldy-Wouthuysen reduction of the full Coulomb gauge Hamiltonian of QCD. There, the Lorentz-structure of confinement has been found to be of vector nature. The scalar character of spin splittings in heavy quarkonia in this approach is dynamically generated through the interaction with collective gluonic degrees of freedom. Thus we see that the spin-dependent structure of $(q\bar{q})$ interaction is well established now, while the spin-independent part is still controversial in the literature.

In preceding papers [8,9] we have developed the relativistic quark model with the $(q\bar{q})$ potential consisting of the perturbative one-gluon exchange part and a nonperturbative one which is a mixture of the Lorentz scalar and vector confining potentials:

$$V(\mathbf{p},\mathbf{q};M) = \bar{u}_a(\mathbf{p})\bar{u}_b(-\mathbf{p})\left\{\frac{4}{3}\alpha_s D_{\mu\nu}(k)\gamma_a^\mu\gamma_b^\nu \right. \left. + V^V_{\rm conf}(k)\Gamma_a^\mu\Gamma_{b;\mu} + V^S_{\rm conf}(k)\right\}u_a(\mathbf{q})u_b(-\mathbf{q}), \quad (1)$$

where $k = p - q$, $D_{\mu\nu}$ is the gluon propagator in the Coulomb gauge and Γ_μ is the effective vector long-range vertex, containing both the Dirac and Pauli terms

$$\Gamma_\mu = \gamma_\mu + \frac{i\kappa}{2m}\sigma_{\mu\nu}k^\nu, \quad (2)$$

$u_{a,b}(\mathbf{p})$ are the Dirac bispinors. Using the identity

$$\bar{u}(\mathbf{p})\Gamma_\mu u(\mathbf{q}) = \bar{u}(\mathbf{p})\left\{\frac{p_\mu + q_\mu}{2m} + \frac{i(1+\kappa)}{2m}\sigma_{\mu\nu}k^\nu\right\}u(\mathbf{q})$$

we can treat the parameter $(1+\kappa)$ as the nonperturbative (long-range) chromomagnetic moment of the quark and κ as its anomalous part (flavour independent).

In the nonrelativistic limit the Fourier transform of (1) gives the static potential

$$V_0(r) = V_{\rm Coul}(r) + V^S_{\rm conf}(r) + V^V_{\rm conf}(r), \quad (3)$$

where

$$V_{\rm Coul}(r) = -\frac{4}{3}\frac{\alpha_s}{r}.$$

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In order to reproduce the linear confining potential

$$V_{\rm conf}(r) = Ar + B$$

in this limit one should put

$$\begin{aligned} V^V_{\rm conf}(r) &= (1-\varepsilon)(Ar+B) \\ V^S_{\rm conf}(r) &= \varepsilon(Ar+B), \end{aligned} \qquad (4)$$

where ε is the mixing parameter.

Now assuming that both quarks are heavy enough we evaluate the (v^2/c^2) relativistic corrections to the static potential (3), (4). For the one-gluon exchange part the retardation effect is taken into account by the contribution of transverse gluon exchange and thus it is sufficient in the adopted approximation to set $k^0 = 0$. For the confining part one should utilize a different procedure (see [10–12]). The Fourier transform of the linear potential Ar in the momentum space looks like:

$$A\int {\rm d}^3 r r e^{-i\mathbf{k}\cdot\mathbf{r}} = -A\frac{8\pi}{|\mathbf{k}|^4}, \quad \mathbf{k} = \mathbf{p} - \mathbf{q}. \qquad (5)$$

The natural (though not unique) relativistic extension (dependent only on the four-momentum transfer) of expression (5) is to substitute $(-\mathbf{k}^2) \to (k_0^2 - \mathbf{k}^2)$ and thus

$$\frac{1}{|\mathbf{k}|^4} \to \frac{1}{(k_0^2 - \mathbf{k}^2)^2}. \qquad (6)$$

Now as mentioned above we should choose the procedure of fixing k_0. On the mass shell due to energy conservation we have $k_0 = 0$. So k_0 may be considered as the measure of deviation either from the mass shell or from the energy shell. We choose the second possibility and set k_0 equal to $\epsilon_a(\mathbf{p}) - \epsilon_a(\mathbf{q})$ or to $\epsilon_b(\mathbf{q}) - \epsilon_b(\mathbf{p})$. Then in the symmetrized form [10,12] we get

$$\begin{aligned} k_0^2 &= -(\epsilon_a(\mathbf{p}) - \epsilon_a(\mathbf{q}))(\epsilon_b(\mathbf{p}) - \epsilon_b(\mathbf{q})), \qquad (7) \\ \epsilon_{a,b}(\mathbf{p}) &= \sqrt{\mathbf{p}^2 + m_{a,b}^2}. \end{aligned}$$

This form is not unique and other possible expressions for k_0^2 are discussed in [12,13]. In favour of choice (7) we mention the following arguments. It is well-known [10,14] that for the one-photon exchange contribution in QED only choice (7) in the Feynman (diagonal) gauge leads to the same correct result (the Breit-Fermi Hamiltonian) as the prescription $k_0 = 0$ in the Coulomb (or transverse Landau) gauge. The same is naturally true for the one-gluon exchange contribution in QCD. Moreover as shown in [14] for any effective vector potential generated by a vector exchange and its couplings to conserved vector currents (vertices) there is the so-called instantaneous gauge which plays the role of the Coulomb gauge. In the instantaneous gauge the prescription $k_0 = 0$ reproduces the same result as the expansion in k_0^2 fixed by (7) in the diagonal gauge used here. The other reason to utilize prescription (7) is the reproduction of the correct Dirac limit in this case [13].

The (p^2/m^2) expansion of (6) with the account of (7) yields:

$$\begin{aligned} &\frac{1}{[-(\epsilon_a(\mathbf{p}) - \epsilon_a(\mathbf{q}))(\epsilon_b(\mathbf{p}) - \epsilon_b(\mathbf{q})) - \mathbf{k}^2]^2} \\ &\quad \cong \frac{1}{|\mathbf{k}|^4}\left[1 - \frac{(\mathbf{p}^2 - \mathbf{q}^2)^2}{2m_a m_b |\mathbf{k}|^2}\right] \\ &\quad = \frac{1}{|\mathbf{k}|^4} - \frac{1}{2m_a m_b |\mathbf{k}|^6}\big[(\mathbf{k}\cdot\mathbf{p})^2 \\ &\qquad + 2(\mathbf{k}\cdot\mathbf{p})(\mathbf{k}\cdot\mathbf{q}) + (\mathbf{k}\cdot\mathbf{q})^2\big]. \qquad (8) \end{aligned}$$

After the Fourier transform of (8) we obtain in the configuration space:

$$\begin{aligned} &-8\pi A\int \frac{{\rm d}^3 k}{(2\pi)^3} e^{i\mathbf{k}\cdot\mathbf{r}} \frac{1}{|\mathbf{k}|^4}\left[1 - \frac{(\mathbf{p}^2 - \mathbf{q}^2)^2}{2m_a m_b |\mathbf{k}|^2}\right] \\ &\quad = Ar - \left\{\frac{Ar}{2m_a m_b}\left[\mathbf{p}^2 + \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}\right]\right\}_W, \qquad (9) \end{aligned}$$

where the notation $\{\ldots\}_W$ means the Weyl ordering prescription for $\mathbf{p}$ and $\mathbf{r}$ variables.

Now we turn to the constant B term in (4). The Fourier transform of it gives in the momentum space

$$B\int {\rm d}^3 r e^{-i\mathbf{k}\cdot\mathbf{r}} = B(2\pi)^3\delta(\mathbf{k}), \quad \mathbf{k} = \mathbf{p} - \mathbf{q}. \qquad (10)$$

The simplest relativistic version of (10) is multiplying it by the energy factor $\epsilon(\mathbf{p})/m$, which in the symmetric form looks like

$$\sqrt{\frac{\epsilon_a(\mathbf{p})\epsilon_b(\mathbf{p})}{m_a m_b}}\,\delta(\mathbf{p} - \mathbf{q}). \qquad (11)$$

Expanding (11) in (p^2/m^2) we get

$$\delta(\mathbf{p} - \mathbf{q})\left[1 + \frac{\mathbf{p}^2}{4}\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)\right]. \qquad (12)$$

So in the configuration space the constant term acquires the form

$$B\left[1 + \frac{1}{4}\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)\mathbf{p}^2\right]. \qquad (13)$$

All other relativistic corrections of order (p^2/m^2) have been considered in [8]. After using the Weyl ordering notations, expressions (9) of [8] take on the form:

$$\begin{aligned} &V_0(r) + \frac{1}{8}\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)\Delta V_{\rm Coul}(r) \\ &+ \frac{1}{8}\left(\frac{1+2\kappa_a}{m_a^2} + \frac{1+2\kappa_b}{m_b^2}\right)\Delta V^V_{\rm conf}(r) \\ &+ \frac{1}{2m_a m_b}\left\{V_{\rm Coul}\left[\mathbf{p}^2 + \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}\right]\right\}_W \\ &+ \frac{1}{m_a m_b}\left\{V^V_{\rm conf}(r)\mathbf{p}^2\right\}_W \\ &- \frac{1}{2}\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)\left\{V^S_{\rm conf}(r)\mathbf{p}^2\right\}_W. \end{aligned}$$

Adding contributions (9) and (13) to the above expressions we obtain the complete spin-independent part of the $(q\bar{q})$ potential ($\kappa_a = \kappa_b = \kappa$):

$$V_{\rm SI}(r) = V_0(r) + V_{\rm VD}(r) + \frac{1}{8}\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right) \times \Delta\big[V_{\rm Coul}(r) + (1+2\kappa)V^V_{\rm conf}(r)\big], \qquad (14)$$

where $V_0(r)$ is given by (3), (4). For the velocity-dependent part $V_{\rm VD}(r)$ we have

$$\begin{aligned} V_{\rm VD}(r) = {}& \frac{1}{2m_am_b}\left\{\left(-\frac{4}{3}\frac{\alpha_s}{r}\right)\left[\mathbf{p}^2 + \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}\right]\right\}_W \\ & + \frac{(1-\varepsilon)}{2m_am_b}\left\{Ar\left[\mathbf{p}^2 - \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}\right]\right\}_W \\ & - \frac{\varepsilon}{2}\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)\{Ar\mathbf{p}^2\}_W \\ & - \frac{\varepsilon}{2m_am_b}\left\{Ar\left[\mathbf{p}^2 + \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}\right]\right\}_W \\ & - \left(\frac{\varepsilon}{2} - \frac{1}{4}\right)\left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)B\mathbf{p}^2 \\ & + \frac{1-\varepsilon}{m_am_b}B\mathbf{p}^2. \end{aligned} \qquad (15)$$

Now representing (15) in the form

$$\begin{aligned} V_{\rm VD}(r) = {}& \frac{1}{m_am_b}\left\{\mathbf{p}^2V_{bc}(r) + \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}V_c(r)\right\}_W \qquad (16) \\ & + \left(\frac{1}{m_a^2} + \frac{1}{m_b^2}\right)\left\{\mathbf{p}^2V_{de}(r) - \frac{(\mathbf{p}\cdot\mathbf{r})^2}{r^2}V_e(r)\right\}_W \end{aligned}$$

with

$$\begin{aligned} V_{bc}(r) &= -\frac{2\alpha_s}{3r} + \left(\frac{1}{2} - \varepsilon\right)Ar + (1-\varepsilon)B; \\ V_{de}(r) &= -\frac{\varepsilon}{2}Ar + \left(\frac{1}{4} - \frac{\varepsilon}{2}\right)B; \\ V_c(r) &= -\frac{2\alpha_s}{3r} - \frac{1}{2}Ar; \quad V_e(r) = 0, \end{aligned} \qquad (17)$$

we are able to test the fulfilment of the exact Barchielli, Brambilla, Prosperi (BBP) relations [15], which follow from the Lorentz invariance of the Wilson loop. In our notations these relations look like

$$\begin{aligned} V_{de} - \frac{1}{2}V_{bc} + \frac{1}{4}V_0 &= 0 \\ V_e + \frac{1}{2}V_c + \frac{r}{4}\frac{{\rm d}V_0}{{\rm d}r} &= 0 \end{aligned} \qquad (18)$$

(in the original version $V_{bc} \equiv -V_b - \frac{1}{3}V_c$ and $V_{de} \equiv V_d + \frac{1}{3}V_e$). One can easily find that the functions (17) identically satisfy relations (18) independently of values of the parameters ε and κ. This is a highly nontrivial result. For the perturbative one-gluon-exchange part of $V_{\rm VD}$ our expressions for $V_b, \ldots, V_e$ are the same as in [15,16], but for the confining (long-range) part they are different, namely the result of [15,16] (from the minimal area law) is as follows:

$$\begin{aligned} V_{bc}(r) &= -\frac{2\alpha_s}{3r} + \frac{1}{6}Ar; \quad V_c(r) = -\frac{2\alpha_s}{3r} - \frac{1}{6}Ar; \\ V_{de}(r) &= -\frac{1}{6}Ar - \frac{1}{4}B; \quad V_e = -\frac{1}{6}Ar. \end{aligned} \qquad (19)$$

No value of ε in (17) can reproduce the above result. The terms with the Laplacian in (14) coincide only for $\kappa = 0$ and $\varepsilon = 0$, i. e. for purely vector confining interaction without the Pauli term in the vertex (2). Our expressions (14) and (15) for purely vector ($\varepsilon = 0$) and purely scalar ($\varepsilon = 1$) interactions and for $\kappa = 0$ coincide with those of [13] except for the constant B term. Our B term for $\varepsilon = 1$ (scalar potential) is the same as in [15]. The B term from [13] does not satisfy the BBP relations (it gives contribution $-B/2$ only to V_{de}). Our result (15) for the scalar ($\varepsilon = 1$) confining potential also differs from the one obtained in [17], where the prescription $k_0 = 0$ was used and as a result the contribution of retardation was lost. The differences between our results and the results presented in [18] originate from the use of specific models such as minimal area law, flux tube, dual superconductivity and stochastic vacuum.

The spin-dependent part of the $(q\bar{q})$ potential is given in [8] ($\kappa_a = \kappa_b = \kappa$):

$$\begin{aligned} V_{\rm SD}(r) = {}& \frac{1}{m_am_b}\frac{1}{r}\frac{\rm d}{{\rm d}r}V_{\rm Coul}(r)\mathbf{L}\cdot(\mathbf{s}_a + \mathbf{s}_b) \\ & + \frac{1}{2m_a^2}\frac{1}{r}\frac{\rm d}{{\rm d}r}\Big\{[V_{\rm Coul}(r) - V_{\rm conf}(r)] \\ & + 2(1+\kappa)\left(1 + \frac{m_a}{m_b}\right)V^V_{\rm conf}(r)\Big\}\mathbf{L}\cdot\mathbf{s}_a \\ & + \frac{1}{2m_b^2}\frac{1}{r}\frac{\rm d}{{\rm d}r}\Big\{[V_{\rm Coul}(r) - V_{\rm conf}(r)] \\ & + 2(1+\kappa)\left(1 + \frac{m_b}{m_a}\right)V^V_{\rm conf}(r)\Big\}\mathbf{L}\cdot\mathbf{s}_b \\ & + \frac{1}{3m_am_b}\Big\{\frac{1}{r}\frac{\rm d}{{\rm d}r}[V_{\rm Coul}(r) + (1+\kappa)^2V^V_{\rm conf}(r)] \\ & - \frac{{\rm d}^2}{{\rm d}r^2}[V_{\rm Coul} + (1+\kappa)^2V^V_{\rm conf}(r)]\Big\} \\ & \times\left[\frac{3}{r^2}(\mathbf{s}_a\cdot\mathbf{r})(\mathbf{s}_a\cdot\mathbf{r}) - \mathbf{s}_a\cdot\mathbf{s}_b\right] + \frac{2}{3m_am_b} \\ & \times\Delta[V_{\rm Coul}(r) + (1+\kappa)^2V^V_{\rm conf}(r)]\mathbf{s}_a\cdot\mathbf{s}_b, \end{aligned} \qquad (20)$$

where $\mathbf{L}$ is the orbital momentum and $\mathbf{s}_{a,b}$ are the spin momenta.

The correct description of the fine structure of the heavy quarkonium mass spectrum requires vanishing of the vector confinement contribution. This can be achieved by putting $1 + \kappa = 0$, i.e. the total long-range quark chromomagnetic moment equals to zero, which is in accord with the flux tube [3] and minimal area [16,18] models. One can see from (20) that for the spin-dependent part of the potential this conjecture is equivalet to the assumption about the scalar structure of confinement interaction

[1]. The specific value of vector-scalar mixing parameter $\varepsilon = -1$ provides the correct description of radiative decays of heavy quarkonia [5].

In this way setting $\kappa = -1$, we obtain:

$$\begin{aligned} V_{\mathrm{SD}}(r) = &\frac{1}{2}\left(\frac{4\alpha_s}{3r^3} - \frac{A}{r}\right)\left[\frac{1}{m_a^2}\mathbf{L}\cdot\mathbf{s}_a + \frac{1}{m_b^2}\mathbf{L}\cdot\mathbf{s}_b\right] \\ &+\frac{1}{m_a m_b}\left(\frac{4\alpha_s}{3r^3}\right)\mathbf{L}\cdot(\mathbf{s}_a+\mathbf{s}_b) \\ &+\frac{8\pi}{3m_a m_b}\left(\frac{4\alpha_s}{3}\right)\delta(\mathbf{r})\mathbf{s}_a\cdot\mathbf{s}_b + \frac{1}{m_a m_b}\left(\frac{4\alpha_s}{3r^3}\right) \\ &\times\left[\frac{3}{r^2}(\mathbf{s}_a\cdot\mathbf{r})(\mathbf{s}_b\cdot\mathbf{r}) - \mathbf{s}_a\cdot\mathbf{s}_b\right]. \end{aligned} \quad (21)$$

Expression (21) for V_{SD} completely coincides with the one found in [16,3]. The Gromes relation is identically fulfilled. Our result supports the conjecture that the long-range confining forces are dominated by chromoelectric interaction and that the chromomagnetic interaction vanishes. It is also in accord with the dual superconductivity picture [19,18]. It is important to mention, that our relativistic quark model is in complete agreement with the heavy quark effective theory (HQET). The model correctly reproduces HQET results for heavy-to-heavy and heavy-to-light weak transition matrix elements with the specific choice of the parameters $\varepsilon = -1$, $\kappa = -1$ (see [9] for details) in accord with the ones found previously [5, 8,20]. It includes the proper description of invariant form factors within the inverse heavy-quark-mass expansion up to the terms of order of $1/m_Q^2$. The mass spectra of D and B mesons have been also calculated in our model in complete agreement with the HQET predictions and available experimental data [21]. It is interesting to note that the relations which are equivalent to the BBP relations (18) can be obtained by use of the reparametrization invariance (in four-velocity) within HQET [22]. The phenomenological implications of retardation corrections will be considered elsewhere.

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