Possible retardation effects of quark confinement on the meson spectrum. II

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We present results of a study of heavy-light-quark bound states in the context of the reduced Bethe-Salpeter equation, with relativistic vector and scalar interactions. We find that satisfactory fits may also be obtained when the retarded effect of the quark-antiquark interaction is concerned. [S0556-2821(99)02919-7]

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Because of the limited understanding of confinement at present, more theoretical effort made regarding this issue is worthwhile. In an earlier paper [1], we presented results of a relativistic analysis of the spectrum of light- and heavy-quark–antiquark bound states based on the reduced Bethe-Salpeter (BS) equation, while retardation effects in the quark-antiquark interaction kernel were taken into consideration. The results are stimulating and appear to have clarified the problem pointed out by Gara *et al.* [2] for the static scalar confinement in the reduced Salpeter equation. The ''intrinsic flaw'' of the Salpeter equation with static scalar confinement could be remedied to some extent by taking the retardation effect into confinement. In the on-shell approximation for the retardation term of linear confinement, the notorious trend of narrow level spacings for quarkonium states, especially for light quarkonium states, is found to be removed. A good fit for mass spectrum of S-wave heavy and light quarkonium states (except for the light pseudoscalar mesons) is obtained using the one-gluon exchange potential and the scalar linear confinement potential, with retardation taken into account.

In this paper we extend our previous study to the heavy-light-quark systems ($Q\bar{q}$ or $q\bar{Q}$) in order to obtain a complete understanding of the retardation effects on meson spectra. Based on the same procedure taken in Ref. [1], by solving the reduced BS equation numerically, the previous conclusion is further substantiated through this study.

We assume the confinement kernel in momentum space to take the form

$$G(p)\propto\frac{1}{(-p^2)^2}=\frac{1}{(\vec{p}^2-p_0^2)^2},\tag{1}$$

as suggested by some authors to be the dressed gluon propagator to implement quark confinement [3]. Here p is the four-momentum exchanged between the quark and antiquark in a meson. If the system is not highly relativistic, we may make the approximation

$$G(p)\propto\frac{1}{(\vec{p}^2-p_0^2)^2}\approx\frac{1}{(\vec{p}^2)^2}\left(1+\frac{2p_0^2}{\vec{p}^2}\right),\tag{2}$$

and may further express p_0 in terms of its on-shell values which are obtained by assuming that quarks are on their mass shells. This should be a good approximation for $c\bar{c}$ and $b\bar{b}$ systems. However, to obtain a qualitative feeling about the retardation effect considered here, we will also use Eq. (2) for heavy-light-quark mesons, though the approximations are not as good as for heavy-heavy-quark mesons. With the above procedure, the scalar confinement kernel again becomes instantaneous, but with some retardation effects taken into the kernel. In the static limit, the retardation term vanishes and the kernel returns to $G(q)\propto 1/(\vec{p}^2)^2$, which is just the Fourier transformation of the linear confining potential. In this paper, we will use this modified scalar confining potential in which the retardation effect is incorporated and the one-gluon-exchange potential in the framework of the reduced Salpeter equation is used to study the mass spectrum of $q\bar{Q}$ mesons. The structure of the hyperfine splittings of the heavy-light-quark systems will also be investigated in this paper.

In quantum field theory, one of the basic descriptions for the bound states is the Bethe-Salpeter equation [4]. We can define the BS wave function of the bound state $|P\rangle$ of a quark $\psi(x_1)$ and an antiquark $\bar{\psi}(x_2)$ as

$$\chi(x_1,x_2)=\langle 0|T\psi(x_1)\bar{\psi}(x_2)|P\rangle.\tag{3}$$

Here T represents the time-order product, and the wave function can be transformed into the momentum space:

$$\chi_P(q)=e^{-iP\cdot X}\int d^4x e^{-iq\cdot x}\chi(x_1,x_2),\tag{4}$$

where P is the four-momentum of the meson and q is the relative momentum between the quark and antiquark. As the standard measure in solving the BS equation, we choose center-of-mass and relative coordinates as variables,

$$X=\eta_1 x_1+\eta_2 x_2, \quad x=x_1-x_2, \tag{5}$$

where $\eta_i=m_i/m_1+m_2$ ($i=1$ and 2). After taking the Fourier transformation, in momentum space the BS equation reads as

$$(\not{p}_1-m_1)\chi_P(q)(\not{p}_2+m_2)=\frac{i}{2\pi}\int d^4k G(P,q-k)\chi_P(k), \tag{6}$$

where p_1 and p_2 represent the momenta of quark and antiquark, respectively,

$$p_1=\eta_1 P+q, \quad p_2=\eta_2 P-q, \tag{7}$$

and $G(P,q-k)$ is the interaction kernel which acts on χ and is determined by the interquark dynamics. Note that in Eq. (6) m_1 and m_2 represent the effective constituent quark masses, so that we could use the effective free propagators of quarks instead of the full propagators. This is an important approximation, and a simplification for light quarks. Furthermore, because of the lack of a fundamental description for the nonperturbative QCD dynamics, we have to make some approximations for the interaction kernel of quarks. In solving Eq. (6), the kernel is taken to be instantaneous, but with some retardation effect taken into it; the negative-energy projectors in the quark propagators are neglected, because in general the negative-energy projectors only contribute in higher orders due to $M-E_1-E_2 \ll M+E_1+E_2$, where M, E_1, and E_2 are the meson mass, the quark kinetic energy, and the antiquark kinetic energy, respectively. Based on the above assumptions the BS equation can be reduced to a three-dimensional equation: i.e., the reduced Salpeter equation

$$(P^0-E_1-E_2)\Phi_{\vec{P}}(\vec{q})=\Lambda_+^1\gamma^0\int d^3k G(\vec{P},\vec{q},\vec{k})\Phi_{\vec{P}}(\vec{k})\gamma^0\Lambda_-^2. \tag{8}$$

Here

$$\Phi_{\vec{P}}(\vec{q})=\int dq^0\chi_P(q^0,\vec{q}) \tag{9}$$

is the three dimensional BS wave function, and

$$\Lambda_+^1=\frac{1}{2E_1}(E_1+\gamma^0\vec{\gamma}\cdot\vec{p}_1+m_1\gamma^0), \tag{10}$$

$$\Lambda_-^2=\frac{1}{2E_2}(E_2-\gamma^0\vec{\gamma}\cdot\vec{p}_2-m_2\gamma^0) \tag{11}$$

are the remaining positive energy projectors of the quark and antiquark, respectively, with $E_1=\sqrt{m_1^2+\vec{p}_1^2}$ and $E_2=\sqrt{m_2^2+\vec{p}_2^2}$. The formal products of $G\Phi$ in Eq. (8) take the form

$$G\Phi=\sum_i G_i O_i \Phi O_i = G_s\Phi+\gamma_\mu\otimes\gamma^\mu G_v\Phi, \tag{12}$$

where $O_i=\gamma_\mu$ corresponds to the perturbative one-gluon-exchange interaction and $O_i=1$ to the scalar confinement potential.

From Eq. (8) it is easy to see that

$$\Lambda_+^1\Phi_{\vec{P}}(\vec{q})=\Phi_{\vec{P}}(\vec{q}), \quad \Phi_{\vec{P}}(\vec{q})\Lambda_-^2=\Phi_{\vec{P}}(\vec{q}). \tag{13}$$

Considering the constraint of Eq. (13) and the requirements of space reflection of bound states, in the meson rest frame ($\vec{P}=0$) the wave function $\Phi_{\vec{P}}(\vec{q})$ for the 0^- and 1^- mesons can be expressed as

$$\Phi_{\vec{P}}^{0^-}(\vec{q})=\Lambda_+^1\gamma^0(1+\gamma^0)\gamma_5\gamma^0\Lambda_-^2\varphi(\vec{q}), \tag{14}$$

$$\Phi_{\vec{P}}^{1^-}(\vec{q})=\Lambda_+^1\gamma^0(1+\gamma^0)\not{e}\gamma^0\Lambda_-^2 f(\vec{q}), \tag{15}$$

where $\not{e}=\gamma_\mu e^\mu$ is the polarization vector of the 1^- meson, and $\varphi(\vec{q})$ and $f(\vec{q})$ are scalar functions of $\vec{q}^2$. It is easy to show that Eqs. (14) and (15) are the most general forms of *S*-wave wave functions for 0^- and 1^- mesons in the rest frame.

Substituting Eqs. (12), (14), and (15) into Eq. (8), one can derives the equations for $\varphi(\vec{q})$ and $f(\vec{q})$ in the meson rest frame [5]:

$$\begin{aligned} M\varphi_1(\vec{q}) &= (E_1+E_2)\varphi_1(\vec{q})-\frac{E_1E_2+m_1m_2+\vec{q}^2}{4E_1E_2} \\ &\quad\times\int d^3k[G_S(\vec{q},\vec{k})-4G_V(\vec{q},\vec{k})]\varphi_1(\vec{k}) \\ &\quad-\frac{(E_1m_2+E_2m_1)}{4E_1E_2}\int d^3k[G_S(\vec{q},\vec{k}) \\ &\quad+2G_V(\vec{q},\vec{k})]\frac{m_1+m_2}{E_1+E_2}\varphi_1(\vec{k}) \\ &\quad+\frac{E_1+E_2}{4E_1E_2}\int d^3k G_S(\vec{q},\vec{k}) \\ &\quad\times(\vec{q}\cdot\vec{k})\frac{m_1+m_2}{E_1m_2+E_2m_1}\varphi_1(\vec{k}) \\ &\quad+\frac{m_1-m_2}{4E_1E_2}\int d^3k[G_S(\vec{q},\vec{k})+2G_V(\vec{q},\vec{k})] \\ &\quad\times(\vec{q}\cdot\vec{k})\frac{E_1-E_2}{E_1m_2+E_2m_1}\varphi_1(\vec{k}), \end{aligned} \tag{16}$$

with

$$\varphi_1(\vec{q})=\frac{(m_1+m_2+E_1+E_2)(E_1m_2+E_2m_1)}{4E_1E_2(m_1+m_2)}\varphi(\vec{q}) \tag{17}$$

and

$$
\begin{aligned}
Mf_1(\vec{q})=&(E_1+E_2)f_1(\vec{q})-\frac{1}{4E_1E_2}\int d^3k[G_S(\vec{q},\vec{k})-2G_V(\vec{q},\vec{k})](E_1m_2+E_2m_1)f_1(\vec{k})\\
&-\frac{E_1+E_2}{4E_1E_2}\int d^3kG_S(\vec{q},\vec{k})\frac{E_1m_2+E_2m_1}{E_1+E_2}f_1(\vec{k})+\frac{E_1E_2-m_1m_2+\vec{q}^2}{4E_1E_2\vec{q}^2}\int d^3k[G_S(\vec{q},\vec{k})+4G_V(\vec{q},\vec{k})](\vec{q}\cdot\vec{k})f_1(\vec{k})\\
&-\frac{E_1m_2-E_2m_1}{4E_1E_2\vec{q}^2}\int d^3k[G_S(\vec{q},\vec{k})-2G_V(\vec{q},\vec{k})](\vec{q}\cdot\vec{k})\frac{E_1-E_2}{m_2+m_1}f_1(\vec{k})-\frac{E_1+E_2-m_2-m_1}{2E_1E_2\vec{q}^2}\int d^3kG_S(\vec{q},\vec{k})\\
&\times(\vec{q}\cdot\vec{k})^2\frac{1}{E_1+E_2+m_1+m_2}f_1(\vec{k})-\frac{m_2+m_1}{E_1E_2\vec{q}^2}\int d^3kG_V(\vec{q},\vec{k})(\vec{q}\cdot\vec{k})^2\frac{1}{E_1+E_2+m_1+m_2}f_1(\vec{k}),
\end{aligned}
\tag{18}
$$

with

$$f_1(\vec{q})=-\frac{m_1+m_2+E_1+E_2}{4E_1E_2}f(\vec{q}). \tag{19}$$

Equation (16) and (18) can also be formally expressed as more compact forms:

$$(M-E_1-E_2)\varphi_1(\vec{q})=\int d^3k\sum_{i=S,V}F_i^{0^-}(\vec{q},\vec{k})G_i(\vec{q},\vec{k})\varphi_1(\vec{k}), \tag{20}$$

$$(M-E_1-E_2)f_1(\vec{q})=\int d^3k\sum_{i=S,V}F_i^{1^-}(\vec{q},\vec{k})G_i(\vec{q},\vec{k})f_1(\vec{k}). \tag{21}$$

When taking the nonrelativistic limit for both the quark and antiquark, and expanding in terms of $\vec{q}^2/m_1^2$ and $\vec{q}^2/m_2^2$, it can be proved that Eqs. (16) and (18) are identical to the Schrödinger equation to zeroth order, and to the Breit equation to first order.

To solve Eq. (6), one must have a good command of the potential between two quarks. At present, the reliable information about the potential only comes from the lattice QCD result, which shows that the potential for a heavy quark-antiquark pair $Q\bar{Q}$ in the static limit is well described by a long-ranged linear confining potential (Lorentz scalar V_S) and a short-ranged one gluon exchange potential (Lorentz vector V_V), i.e., [6,7]

$$V(r)=V_S(r)+\gamma_\mu\otimes\gamma^\mu V_V(r), \tag{22}$$

with

$$V_S(r)=\lambda r\frac{(1-e^{-\alpha r})}{\alpha r}, \tag{23}$$

$$V_V(r)=-\frac{4}{3}\frac{\alpha_s(r)}{r}e^{-\alpha r}, \tag{24}$$

where the introduction of the factor $e^{-\alpha r}$ is not only for the sake of avoiding the infrared divergence but also incorporating the color screening effects of the dynamical light quark pairs on the ''quenched'' $Q\bar{Q}$ potential [8]. Although the lattice QCD result for the $Q\bar{Q}$ potential is supported by the heavy quarkonium spectroscopy including both spin-independent and -dependent effects [9–11], we will employ this static potential below to the heavy-light-quark systems as an assumption. The interaction potentials in Eq. (22) can be transformed straightforwardly into momentum space, where the strong coupling constant

$$\alpha_s(\vec{p})=\frac{12\pi}{27}\frac{1}{\ln\left(a+\dfrac{\vec{p}^2}{\Lambda_{\rm QCD}^2}\right)}. \tag{25}$$

is assumed to be a constant of $O(1)$ as $\vec{p}^2\to 0$. The constants λ, α, a, and $\Lambda_{\rm QCD}$ are the parameters that characterize the potential.

In taking the retardation effect of the scalar confinement into consideration, as discussed in Ref. [1], the confinement will be approximately introduced by adding a retardation term $2p_0^2/\vec{p}^6$ to the instantaneous part $1/(\vec{p}^2)^2$ as given in Eq. (2), and p_0^2 will be treated to take its on-shell values which are obtained by assuming that the quarks are on their mass shells, which means that the retardation term will become instantaneous rather than convoluted. By this procedure the modified scalar confinement potential will include the retardation effect, and become

$$
\begin{aligned}
G_S(\vec{p})\to G_S(\vec{p},\vec{k})=&-\frac{\lambda}{\alpha}\delta^3(\vec{p})+\frac{\lambda}{\pi^2}\frac{1}{(\vec{p}^2+\alpha^2)^2}\\
&+\frac{2\lambda}{\pi^2}\frac{1}{(\vec{p}^2+\alpha^2)^3}(\sqrt{(\vec{p}+\vec{k})^2+m^2}\\
&-\sqrt{\vec{k}^2+m^2})^2,
\end{aligned}
\tag{26}
$$

which is related not only to the interquark momentum exchange $\vec{p}$ but also the quark momentum $\vec{k}$ itself.

TABLE I. Calculated mass spectrum of $q\bar{Q}$ states using a reduced Salpeter equation with retardation for scalar confinement. The experimental data are taken from Ref. [13].

State	Quarks	Data (MeV)	Fit I[a] (Mev)	Error (MeV)	Fit II[b] (Mev)	Error (MeV)
$\bar{B}^0$	$b\bar{d}$	5279	5381	$+102$	5258	-21
D_s^+	$c\bar{s}$	1969	2097	$+128$	1946	-23
D_s^{*+}	$c\bar{s}$	2112	2148	$+35$	2094	-19
D^+	$c\bar{d}$1869	1983	$+114$	1862	-7	
D^{*+}	$c\bar{d}$	2010[c]	0	2003	-7	
$\bar{K}^0$	$s\bar{d}$	498	743	$+245$	652	$+154$
$\bar{K}^{*0}$	$s\bar{d}$	892	870	-22	898	$+6$
D_s^{*+}-D_s^+		144	51	-93	148	$+4$
D^{*+}-D^+	141	27	-114	141	0	
$\bar{K}^{*0}$-$\bar{K}^0$	394	127	-267	244	-150	

[a]Results without retardation.
[b]Results with retardation.
[c]Used to fix the parameters.

Based on the formalism and discussions above, we can now embark on numerical calculations, in which, for input parameters, we take the values

$$\lambda = 0.21\,\text{GeV}^2, \quad \alpha = 0.06\,\text{GeV}, \quad a = e = 2.7183,$$

$$\Lambda_{\text{QCD}} = 0.19\,\text{GeV}, \quad C = -0.05 \tag{27}$$

and

$$m_u = m_d = 0.35\,\text{GeV}, \quad m_s = 0.5\,\text{GeV},$$

$$m_c = 1.68\,\text{GeV}, \quad m_b = 4.925\,\text{GeV}, \tag{28}$$

which fall in the scopes of customarily usage. The numerical results with retardation are listed in Table I. For the convenience of comparison, results obtained without retardation are also presented.

From Table I one can immediately see that the calculated masses with the retardation effect are generally well fitted compared with those without retardation considered, and the spin splittings $1^3S_1 - 1^1S_0$ are significantly improved by adding the retardation term to the scalar confinement potential. These conclusions obviously shed light on the usefulness of the Bethe-Salpeter equation in describing systems containing light quarks; however, at the same time they also indicate the need for further investigations on the interaction kernels.

As noted in Ref. [2], the smallness of the hyperfine splitting obtained for $q\bar{Q}$ mesons is due to the weakness of the binding potential in these systems. However, similarly to what was shown in Ref. [1], this situation may also be changed after including the retardation effect into the interaction kernel. For example, in the equal-quark-mass special case the coefficients for the scalar potential G_s, which plays the main role in setting the spin splittings in Eqs. (20) and (21), will reduce to

$$F_S^{0^-}(\vec{q},\vec{k}) = -\frac{1}{2} + \frac{\vec{q}\cdot\vec{k} - m^2}{2E_qE_k}, \tag{29}$$

$$F_S^{1^-}(\vec{q},\vec{k}) = -\frac{m(E_q+E_k) - \vec{q}\cdot\vec{k}}{2E_q^2} - \frac{(\vec{q}\cdot\vec{k})^2}{2E_q^2\vec{q}^2}\frac{(E_q+m)}{(E_k+m)} \tag{30}$$

coresponding to 0^- and 1^- mesons, respectively. It is clear that these coefficients will be limited to m/E or of higher order, while $\vec{q}\to\vec{k}$. On the other hand, however, the static linear confining potential in momentum space, which behaves as $G_s(\vec{q}-\vec{k}) \propto (\vec{q}-\vec{k})^{-4}$, is strongly weighted as $\vec{q}\to\vec{k}$ in Eqs. (20) and (21). Because the coefficients will diminish in the relativistic limit as $\vec{q}\to\vec{k}$, the strength of the confining potential would be reduced in turn. This is the reason which leads to the small spin splittings obtained for heavy-light-quark mesons. However, this depressing situation would change if retardation is taken into account. In fact, the covariant form of the confinement interaction may take the form $G_S(q,k) \propto [(\vec{q}-\vec{k})^2 - (q_0-k_0)^2]^{-2}$, and in the on-shell approximation $q_0^2 = m^2 + \vec{q}^2$, $k_0^2 = m^2 + \vec{k}^2$ it becomes

$$G_S(q,k) \propto (-2m_q^2 + qk - \vec{q}\cdot\vec{k})^{-2} \tag{31}$$

in the high energy limit, i.e., $p,k \gg m$. We can see that with the retardation effect the scalar interaction $G_S(q,k)$ is heavily weighted as $\vec{q}$ and $\vec{k}$ are collinear ($\vec{q}\|\vec{k}$), whereas the static linear potential only peaks at $\vec{q}=\vec{k}$. This indicates that the former is weighted in a much wider kinematic region than the latter, especially for systems including light components. As a consequence, the wave functions in coordinate space would tend to be short ranged, and hence the magnitudes of the wave function at the origin, $\psi(0)$, would increase. Because, to leading order in the $1/m^2$ expansion, the hyperfine splitting is proportional to $|\psi(0)|^2$, the Hermitan square of the wave function of the meson at the origin, the splittings would enlarge as well. The above analysis indicates that in the equal-quark-mass situation the modified effective scalar interaction will not be weakened too much as $\vec{q}$ and $\vec{k}$ approach a parallel configuration and this is just due to the retardation effect. In our opinion, the difficulty in the reduced Salpeter equation with static scalar confinement is probably due to the improper treatment, in which the confining interaction is purely instantaneous [1,12].

In practice, for the constituent quark model, which is essentially used in the present work, the equal-mass and on-shell approximations may not be good simplifications for the $q\bar{Q}$ systems; however, in any case the analysis given above is qualitatively correct, which is supported by the numerical results listed in Table I. It is obvious that the new procedure gives a much better fit to data than the previous one, especially in the spin splittings. The results obtained for the mass spectrum, particularly the sd system, still seem not fully satisfactory, but it is noted that our results are just schematic ones. A fine tuning may lead to improvement.

In conclusion, in this paper we have extended our previous study in clarifying the problems pointed out by Gara *et al.* for static scalar confinement in the reduced Salpeter equation to heavy-light-quark systems. The conclusion remains the same in that the ''intrinsic flaw'' of the Salpeter equation with static scalar confinement could be remedied to some extent by taking the retardation effect of the confinement into consideration. A fit for the mass spectrum of *S*-wave heavy-light-quark systems is obtained by using a scalar linear confinement potential with retardation, and a one-gluon exchange potential. Results show that a great improvement is achieved on fitting the data with retardation taken into account. Although the on-shell approximation may not be a rigorous treatment here, the qualitative feature of the retardation effect is still manifest. Nevertheless, it is still premature to assess whether or not quark confinement is really represented by a scalar exchange of the form of $(\vec{p}^2 - p_0^2)^{-2}$, as suggested by some authors, and as used here as a dressed gluon propagator to implement quark confinement. Therefore, further investigations on this subject are necessary.

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