COMMENTS

Comments are short papers which criticize or correct papers of other authors previously published in the **Physical Review**. *Each Comment should state clearly to which paper it refers and must be accompanied by a brief abstract. The same publication schedule as for regular articles is followed, and page proofs are sent to authors.*

Comment on "Relativistic description of quark-antiquark bound states"

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By means of an effective-Hamiltonian method we reconsider the derivation of the effective interaction in a fermion-antifermion system. Furthermore, we point out some errors of Gara and co-workers in their treatment of fermion-antifermion bound states by solving the reduced Salpeter equation in configuration space.

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In two recent papers [1,2] Gara and co-workers investigated the relativistic description of bound states of quark-antiquark pairs. Their analysis was based on the reduced Salpeter equation. The Salpeter equation [3] is derived from the Bethe-Salpeter equation [4] upon eliminating any dependence on timelike variables in a suitable manner. Some standard and plausible approximations, such as the restriction to positive-energy solutions, then lead to the reduced Salpeter equation.

In this Comment we would like to present an alternative way of treating fermion-antifermion bound states relativistically, namely, by construction of an effective Hamiltonian for the two-particle system under consideration. The procedure advocated for consists of two main steps [5,6].

(1) Compute the effective interaction potential between the bound-state constituents (at least to the extent you trust in perturbation theory) from the elastic scattering of the involved particles, more precisely, from the Fourier transform of the corresponding transition amplitude [7,8].

(2) Use this potential in a multiparticle Schrödinger equation with relativistically correct kinetic Hamiltonian in order to determine the energy eigenvalues and corresponding eigenstate vectors of the bound state under consideration.

Obviously, this effective-Hamiltonian method might be regarded as the relativistic generalization of the description of fermion-antifermion bound states in terms of nonrelativistic potential models. As far as the incorporation of relativistic kinematics is concerned, it provides a description of bound states which is of equal quality as the reduced Salpeter equation adopted in Refs. [1,2]. The obvious advantage of our approach is its physical transparency.

The basic idea of the proposed effective-Hamiltonian method [5,6] is to approximate by a potential the (perturbatively accessible part of the) interaction between particles which in fact are described by some quantum field theory. To this end consider the elastic scattering

$$f(\mathbf{p}_1,\tau_1)+\bar{f}(\mathbf{p}_2,\tau_2)\rightarrow f(\mathbf{q}_1,\tau_3)+\bar{f}(\mathbf{q}_2,\tau_4) \tag{1}$$

of the involved fermion f and antifermion $\bar{f}$ (with masses m_1 and m_2, respectively). Expressed in terms of Dirac spinors $u(\mathbf{p},\tau)$ and $v(\mathbf{p},\tau)$, the general form of the corresponding transition amplitude T is

$$T=\frac{1}{(2\pi)^6}\frac{m_1 m_2}{\sqrt{E_1E_2E_3E_4}} \times\bar{u}(\mathbf{q}_1,\tau_3)\Gamma_1 u(\mathbf{p}_1,\tau_1)\,\bar{v}(\mathbf{p}_2,\tau_2)\Gamma_2 v(\mathbf{q}_2,\tau_4)K\ , \tag{2}$$

where Γ_i, $i=1,2$, represent some Dirac matrices. The (unspecified) interaction kernel, which depends on the dynamics of the theory responsible for the interaction of the particles under consideration, is denoted by K. It is usually assumed to be a function of the (modulus of the) momentum transfer $\mathbf{k}\equiv\mathbf{p}_1-\mathbf{q}_1$ only. The Fourier transform of just this kernel yields the static interaction potential in the Salpeter equation, but with the fermion and antifermion restricted from the very beginning to have positive energies, that is, with the "small-small" term in the Salpeter equation dropped as in Refs. [1,2].

There are many indications that the dominant spin structure for the quark-antiquark interaction, originating from quantum chromodynamics, is a vector, i.e., $\Gamma_1\otimes\Gamma_2=\gamma_\mu\otimes\gamma^\mu$, plus scalar, i.e., $\Gamma_1\otimes\Gamma_2=1\otimes 1$. (For a very recent review on the phenomenological aspects of the forces acting within bound states of quarks, see, e.g., Ref. [8].) Accordingly, in the following we shall focus

our interest to precisely this spin structure of the fermion-antifermion interaction.

Explicitly, in the Dirac representation the Dirac spinors are given by

$$u(\mathbf{p}_i,\tau_i)=\left(\frac{S_i}{2m_i}\right)^{1/2}\begin{pmatrix}1\\ \dfrac{\boldsymbol{\sigma}\cdot\mathbf{p}_i}{S_i}\end{pmatrix}\chi_{\tau_i}\,,$$
$$v(\mathbf{p}_i,\tau_i)=\left(\frac{S_i}{2m_i}\right)^{1/2}\begin{pmatrix}\dfrac{\boldsymbol{\sigma}\cdot\mathbf{p}_i}{S_i}\\ 1\end{pmatrix}\chi^c_{\tau_i},\quad \chi^c_{\tau_i}\equiv -i\sigma_2\chi^*_{\tau_i}\,, \tag{3}$$

where χ_τ is the two-component spinor corresponding to spin projection τ, and we defined

$$E_1\equiv\sqrt{\mathbf{p}_1^2+m_1^2},\quad E_2\equiv\sqrt{\mathbf{p}_2^2+m_2^2}\,,$$
$$E_3\equiv\sqrt{\mathbf{q}_1^2+m_1^2},\quad E_4\equiv\sqrt{\mathbf{q}_2^2+m_2^2} \tag{4}$$

and

$$S_1\equiv E_1+m_1,\quad S_2\equiv E_2+m_2\,,$$
$$S_3\equiv E_3+m_1,\quad S_4\equiv E_4+m_2\,. \tag{5}$$

Inserting this into (2), we immediately obtain, in the center-of-momentum system,

$$\mathbf{p}\equiv\mathbf{p}_1=-\mathbf{p}_2,\quad \mathbf{q}\equiv\mathbf{q}_1=-\mathbf{q}_2\,, \tag{6}$$

for the scattering amplitude of vectorial spin structure [5],

$$\begin{aligned}T_V=N_1N_2N_3N_4\Big[&\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}+\frac{1}{S_1S_3}\delta_{\tau_2\tau_4}(\mathbf{p}\cdot\mathbf{q}\delta_{\tau_1\tau_3}-i\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_1)+\frac{1}{S_2S_4}\delta_{\tau_1\tau_3}(\mathbf{p}\cdot\mathbf{q}\delta_{\tau_2\tau_4}-i\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_2)\\
&+\frac{1}{S_1S_2}[\mathbf{p}^2\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-\mathbf{p}^2\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2+(\mathbf{p}\cdot\boldsymbol{\sigma}_1)(\mathbf{p}\cdot\boldsymbol{\sigma}_2)]\\
&+\frac{1}{S_3S_4}[\mathbf{q}^2\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-\mathbf{q}^2\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2+(\mathbf{q}\cdot\boldsymbol{\sigma}_1)(\mathbf{q}\cdot\boldsymbol{\sigma}_2)]\\
&+\frac{1}{S_1S_4}[\mathbf{p}\cdot\mathbf{q}\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-i\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_++(\mathbf{p}\cdot\mathbf{q})(\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2)-(\mathbf{q}\cdot\boldsymbol{\sigma}_1)(\mathbf{p}\cdot\boldsymbol{\sigma}_2)]\\
&+\frac{1}{S_2S_3}[\mathbf{p}\cdot\mathbf{q}\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-i\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_++(\mathbf{p}\cdot\mathbf{q})(\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2)-(\mathbf{p}\cdot\boldsymbol{\sigma}_1)(\mathbf{q}\cdot\boldsymbol{\sigma}_2)]\\
&+\frac{1}{S_1S_2S_3S_4}[(\mathbf{p}\cdot\mathbf{q})^2\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-i(\mathbf{p}\cdot\mathbf{q})(\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_+)-(\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_1)(\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_2)]\Big]K_V\,,\end{aligned} \tag{7}$$

and, for the scattering amplitude of scalar spin structure [5],

$$\begin{aligned}T_S=-N_1N_2N_3N_4\Big[&\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-\frac{1}{S_1S_3}\delta_{\tau_2\tau_4}(\mathbf{p}\cdot\mathbf{q}\delta_{\tau_1\tau_3}-i\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_1)-\frac{1}{S_2S_4}\delta_{\tau_1\tau_3}(\mathbf{p}\cdot\mathbf{q}\delta_{\tau_2\tau_4}-i\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_2)\\
&+\frac{1}{S_1S_2S_3S_4}[(\mathbf{p}\cdot\mathbf{q})^2\delta_{\tau_1\tau_3}\delta_{\tau_2\tau_4}-i(\mathbf{p}\cdot\mathbf{q})(\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_+)-(\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_1)(\mathbf{p}\times\mathbf{q}\cdot\boldsymbol{\sigma}_2)]\Big]K_S\,.\end{aligned} \tag{8}$$

Here we introduced the shorthand notation

$$\boldsymbol{\sigma}_1\equiv\chi^\dagger_{\tau_3}\boldsymbol{\sigma}\chi_{\tau_1}\,,$$
$$\boldsymbol{\sigma}_2\equiv\chi^\dagger_{\tau_4}\boldsymbol{\sigma}\chi_{\tau_2}\,, \tag{9}$$
$$\boldsymbol{\sigma}_+\equiv\boldsymbol{\sigma}_1\delta_{\tau_2\tau_4}+\boldsymbol{\sigma}_2\delta_{\tau_1\tau_3}\,.$$

The normalization factors N_i are given by

$$N_i\equiv\frac{1}{(2\pi)^{3/2}}\left(\frac{S_i}{2E_i}\right)^{1/2}. \tag{10}$$

From this, the interaction Hamiltonian may be found from the Fourier transform with respect to the momentum transfer $\mathbf{k}\equiv\mathbf{p}-\mathbf{q}$. In the nonrelativistic expansion up to next-to-lowest order one recovers in this way, of course, the well-known Breit-Fermi Hamiltonian.

We observe, however, the following discrepancies between our result and the one given in Eqs. (6a) and (6b) of Ref. [2]: (1) an overall sign in the third line of Eq. (6b); (2) the last term in the fourth line of Eq. (6b) does not exhibit the symmetry under interchange of $\mathbf{p}$ and $\mathbf{q}$ one would expect for this term; (3) an overall sign on the right-hand side of Eq. (6a). (The authors of Refs. [1,2] concur with the results obtained here; see Ref. [9].)

Furthermore, we feel obliged to make some additional remarks to the two papers by Gara and co-workers [1,2].

First of all, an eventually confusing point might be the somewhat asymmetric behavior of the interaction functions as defined in Refs. [1,2], F_V and F_S, under interchange of the relative momenta of initial and final state. In contrast with that, the approach to relativistic bound states via the effective-Hamiltonian method introduces these momenta in a symmetric way, as can immediately be seen from the transition amplitudes (7) and (8). For the case of equal fermion masses, $m_1=m_2=m$, the relevant T-matrix elements read explicitly, for vectorial spin structure,

$$T_V=\frac{1}{(2\pi)^6}\frac{1}{4E_pE_q}\left[(E_p+m)(E_q+m)+2\mathbf{p}\cdot\mathbf{q}(2+\sigma_1\cdot\sigma_2)-3i\mathbf{p}\times\mathbf{q}\cdot\sigma_+-(\mathbf{q}\cdot\sigma_1)(\mathbf{p}\cdot\sigma_2)-(\mathbf{p}\cdot\sigma_1)(\mathbf{q}\cdot\sigma_2)\right.$$
$$+\frac{E_q+m}{E_p+m}[\mathbf{p}^2(1-\sigma_1\cdot\sigma_2)+(\mathbf{p}\cdot\sigma_1)(\mathbf{p}\cdot\sigma_2)]+\frac{E_p+m}{E_q+m}[\mathbf{q}^2(1-\sigma_1\cdot\sigma_2)+(\mathbf{q}\cdot\sigma_1)(\mathbf{q}\cdot\sigma_2)]$$
$$\left.+\frac{1}{(E_p+m)(E_q+m)}[(\mathbf{p}\cdot\mathbf{q})^2-i(\mathbf{p}\cdot\mathbf{q})(\mathbf{p}\times\mathbf{q}\cdot\sigma_+)-(\mathbf{p}\times\mathbf{q}\cdot\sigma_1)(\mathbf{p}\times\mathbf{q}\cdot\sigma_2)]\right]K_V \tag{11}$$

and, for scalar spin structure,

$$T_S=-\frac{1}{(2\pi)^6}\frac{1}{4E_pE_q}\left[(E_p+m)(E_q+m)-2\mathbf{p}\cdot\mathbf{q}+i\mathbf{p}\times\mathbf{q}\cdot\sigma_+\right.$$
$$\left.+\frac{1}{(E_p+m)(E_q+m)}[(\mathbf{p}\cdot\mathbf{q})^2-i(\mathbf{p}\cdot\mathbf{q})(\mathbf{p}\times\mathbf{q}\cdot\sigma_+)-(\mathbf{p}\times\mathbf{q}\cdot\sigma_1)(\mathbf{p}\times\mathbf{q}\cdot\sigma_2)]\right]K_S\ , \tag{12}$$

where now

$$E_p\equiv\sqrt{\mathbf{p}^2+m^2},\quad E_q\equiv\sqrt{\mathbf{q}^2+m^2},\quad \sigma_+\equiv\sigma_1+\sigma_2\ , \tag{13}$$

and any reference to the spin degrees of freedom has been suppressed, so that now the Pauli matrices σ_i have to be understood simply to act on particle i, $i=1,2$, respectively.

The spin-independent parts of the above scattering amplitudes, $T_V^{\rm SI}$ and $T_S^{\rm SI}$, may be extracted by taking the trace over the spin degrees of freedom of the fermions, with the result

$$T_V^{\rm SI}\equiv\frac{1}{4}\,{\rm Tr}_1\,{\rm Tr}_2 T_V$$
$$=\frac{1}{(2\pi)^6}\frac{1}{4E_pE_q}\left[(E_p+m)(E_q+m)+4\mathbf{p}\cdot\mathbf{q}\right.$$
$$+\frac{E_q+m}{E_p+m}\mathbf{p}^2+\frac{E_p+m}{E_q+m}\mathbf{q}^2$$
$$\left.+\frac{(\mathbf{p}\cdot\mathbf{q})^2}{(E_p+m)(E_q+m)}\right]K_V \tag{14}$$

and

$$T_S^{\rm SI}\equiv\frac{1}{4}\,{\rm Tr}_1\,{\rm Tr}_2 T_S$$
$$=-\frac{1}{(2\pi)^6}\frac{1}{4E_pE_q}\left[(E_p+m)(E_q+m)-2\mathbf{p}\cdot\mathbf{q}\right.$$
$$\left.+\frac{(\mathbf{p}\cdot\mathbf{q})^2}{(E_p+m)(E_q+m)}\right]K_S\ . \tag{15}$$

These expressions have to be compared to (the equal-mass case of) Eqs. (15a) and (15b) of Ref. [1]. (In doing this, we detected a further sign error in Eq. (A7) of Ref. [1].) The interaction functions in the effective-Hamiltonian approach, $T_V^{\rm SI}/K_V$ and $T_S^{\rm SI}/K_S$, differ from those entering in the reduced Salpeter equation, $F_V^{\rm SI}$ and $F_S^{\rm SI}$, by an overall factor $(E_p+m)(E_q+m)^{-1}E_qE_p^{-1}$.

The reason for this discrepancy may be traced back to the choice of normalization of the employed wave functions. Our momentum-space wave function, denoted by $\psi(\mathbf{p})$, differs from the (unnormalized) one adopted in Refs. [1,2], $\phi(\mathbf{p})$, by a factor $2E_p(E_p+m)^{-1}$ [10]:

$$\psi(\mathbf{p})=\frac{2E_p}{E_p+m}\phi(\mathbf{p})\ . \tag{16}$$

Of course, any change of normalization, like the above one, has no influence on the resulting spectrum of eigenvalues. In the coordinate-space representation of the equation of motion, however, the free energy E has to be regarded as the nonlocal differential operator $E\equiv\sqrt{-\Delta+m^2}$. As a consequence of this, any rescaling in the definition of the wave function by an energy-dependent factor entails a change in the ordering of the various factors which involve this free-energy expression.

Second, one angular-momentum-dependent term is missing in one of the identities used by the authors of Refs. [1,2] in order to obtain from the configuration-space representation of the reduced Salpeter equation, by pushing the spherical harmonics $\mathcal{Y}_{lm}(\theta,\phi)$ through all the differential operators, a differential equation for the radial part $R(r)$ of the bound-state wave function $\psi(\mathbf{r})=R(r)\mathcal{Y}_{lm}(\theta,\phi)$. In general, the identity in question is given by the relation

$$\nabla_i\nabla_j V(r)\nabla_i\nabla_j R(r)\mathcal{Y}_{lm}(\theta,\phi)=\mathcal{Y}_{lm}(\theta,\phi)\left[\frac{d^2V(r)}{dr^2}-\frac{1}{r}\frac{dV(r)}{dr}\right]\left[\frac{2}{3}\left[\frac{d^2R(r)}{dr^2}-\frac{1}{r}\frac{dR(r)}{dr}\right]+\frac{1}{3}\Delta R(r)\right]$$
$$+\frac{1}{3}\left[\frac{1}{r}\frac{dV(r)}{dr}-\frac{d^2V(r)}{dr^2}+\Delta V(r)\right]\Delta R(r)\mathcal{Y}_{lm}(\theta,\phi)$$
$$+\frac{2}{r}\frac{dV(r)}{dr}\mathbf{r}\cdot\nabla\Delta R(r)\mathcal{Y}_{lm}(\theta,\phi)+V(r)\Delta\Delta R(r)\mathcal{Y}_{lm}(\theta,\phi)\ . \tag{17}$$

If $R(r)$ is regular, i.e., nonsingular, the action of the Laplacian on the wave function $\psi(\mathbf{r})$ may be written in the form

$$\Delta R(r)\mathcal{Y}_{lm}(\theta,\phi)=\mathcal{Y}_{lm}(\theta,\phi)\Delta_l R(r)\ , \tag{18}$$

where, in the notation of Gara and co-workers,

$$\Delta_l\equiv\frac{d^2}{dr^2}+\frac{2}{r}\frac{d}{dr}-\frac{l(l+1)}{r^2}\ . \tag{19}$$

Similarly, under the assumption that $R(r)$, $\Delta_l R(r)$, as well as the potential $V(r)$ are regular functions of the radial coordinate r, the above identity reduces to

$$\nabla_i\nabla_j V(r)\nabla_i\nabla_j R(r)\mathcal{Y}_{lm}(\theta,\phi)=\mathcal{Y}_{lm}(\theta,\phi)\left[\frac{d^2V(r)}{dr^2}\frac{d^2}{dr^2}+\frac{dV(r)}{dr}\left[\frac{d}{dr}\Delta_l+\Delta_l\frac{d}{dr}+\frac{l(l+1)}{r^3}\right]+V(r)\Delta_l\Delta_l\right]R(r)\ . \tag{20}$$

Comparing this result to the relations given in Refs. [1,2], we find that the term

$$\frac{dV(r)}{dr}\frac{l(l+1)}{r^3}$$

in the square brackets on the right-hand side of (20) is missing in the fourth of Eqs. (22) in Ref. [1] (which is identical to Eq. (A2d) of Ref. [2]), and consequently also in the relativistic wave equation which provides the basis for the whole analysis of Gara and co-workers (Eq. (24) of Ref. [1]). (The authors of Refs. [1,2] agree with the form (20) of the identity in question [10].)

Taking into account the above remarks, that is, the rearrangement of the energy-dependent factors brought about by the difference in the overall normalization of the interaction functions as well as the inclusion of the missing term in the expressions involving the fourth derivative (20), the spin-independent equation of motion for the radial part $R(r)$ of the wave function becomes

$$\begin{aligned}(M-2E_l)R(r)=\frac{1}{4E_l}\Bigg\{&(E_l+m)[V_V(r)+V_S(r)](E_l+m)-2\left[2\frac{dV_V(r)}{dr}-\frac{dV_S(r)}{dr}\right]\frac{d}{dr}-2[2V_V(r)-V_S(r)]\Delta_l\\&-(E_l+m)V_V(r)\frac{\Delta_l}{E_l+m}-\frac{\Delta_l}{E_l+m}V_V(r)(E_l+m)\\&+\frac{1}{E_l+m}\Bigg[\left[\frac{d^2V_V(r)}{dr^2}+\frac{d^2V_S(r)}{dr^2}\right]\frac{d^2}{dr^2}+\left[\frac{dV_V(r)}{dr}+\frac{dV_S(r)}{dr}\right]\left[\frac{d}{dr}\Delta_l+\Delta_l\frac{d}{dr}+\frac{l(l+1)}{r^3}\right]\\&+[V_V(r)+V_S(r)]\Delta_l\Delta_l\Bigg]\frac{1}{E_l+m}\Bigg\}\frac{1}{E_l}R(r)\ .\end{aligned} \tag{21}$$

Here M denotes the mass of the bound state, E_l labels the free-energy differential operator given by

$$E_l\equiv\sqrt{-\Delta_l+m^2}\ , \tag{22}$$

and $V_V(r)$ and $V_S(r)$ are the static potentials obtained by Fourier transformation from the interaction kernels K_V and K_S, respectively, which are usually assumed to depend only on the modulus of the momentum transfer $\mathbf{k}\equiv\mathbf{p}-\mathbf{q}$. Similar modifications apply, of course, also to the spin-dependent discussion in Ref. [2].

In their rather comprehensive and detailed investigation of quark-antiquark bound states Gara and co-workers arrive at the somewhat surprising conclusion that, contrary to one's physical intuition, a relativistic treatment yields no improvement in the description of the meson mass spectra compared to the corresponding nonrelativistic one. In view of the importance of these findings, a reanalysis of the significance of the (spin-dependent counterpart of the) wave equation (21) for the description of relativistic bound states has been performed [11]. We find, in accordance with [10], that the effects of the $l(l+1)/r^3$ term on quarkonia spectra are quite small [11].

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