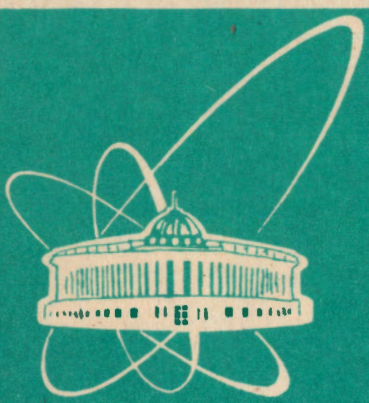

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«DRESSING» AND BOUND STATES IN QUANTUM FIELD THEORY

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I. Introduction

The interaction of charged particles with photons or a quark-gluon interaction give rise to bound states such as the hydrogen atom (pe^-), the positronium (e^+e^-), and hadrons. The problem of determining the states is of importance in field theories.

Different approaches to this problem are known, e.g., nonrelativistic Breit's approach [1,2], Bethe-Salpeter equation (see the book [3] and bibliography given in [4]), three-dimensional quasi-potential approaches (see [5,6] and references therein).

I shall show here that the "dressing" formalism given by L.Faddeev [7] (see also [8]) provides one more approach. The idea of the approach is illustrated by using a simple nonrelativistic model, see sections 2 and 3. The peculiarities of the relativistic case (the spinor QED) are briefly discussed in sect.3. The simple model used in sect.2 allows one to consider an instructive application of the formalism to the so-called multipolar form of QED. Some consequences of the "dressed" description of particles are discussed in sect.5. In the summarizing sect.6 I list several qualitative peculiarities of the suggested approach to the bound-state problem.

2. Introduction of the "dressed" particle creation- destruction operators

The definition of "dressed" particles and their introduction proposed in [7,8] need modifications when applied to QED or QCD. The "dressing" according to [7,8] requires an infinite series of specific unitary transformations. A simple "partial" dressing presented here is realised by the simplest of these transformations applied moreover to a nonrelativistic field model.

2.1. The model

Consider the second-quantized theory of nonrelativistic spinless charged particles (named electrons below) interacting with photons. The total Hamiltonian is

$$H=\int d^3x\,\psi^+(\bar{x})\Big\{\big(-i\bar{\nabla}\big)^2/2\mu_e+V(\bar{x})+\frac{e}{2\mu_e}\big[\bar{A}(\bar{x})(-i\bar{\nabla})+(-i\bar{\nabla})\bar{A}(\bar{x})\big]\Big\}\psi(x)+$$
$$+\;H_{ph}\;,\qquad H_{ph}=\int d^3x\,\big[\bar{E}^2(\bar{x})+\bar{H}^2(\bar{x})\big]. \tag{1}$$

Here $\psi(\bar{x})$ is the particle field satisfying the anticommutation relation $\{\psi(\bar{x}), \psi^+(\bar{y})\}_+ = \delta(\bar{x}-\bar{y})$. The absence of antiparticles is the main simplification as compared to the spinor QED.

Let φ_n be the complete set of solutions of the equation

$$[-\Delta/2\mu_0 + V_0(\bar{x})]\,\varphi_n(\bar{x}) = E_n \varphi_n(\bar{x}). \qquad (2)$$

Here V_0 is a part of the potential V : $V = V_0 + P$ (such that eq.(2) can be solved exactly). In the expansion

$$\psi(\bar{x}) = \sum_n a_n \varphi_n(\bar{x}) \qquad (3)$$

the coefficients a_n satisfy $\{a_m, a_n^+\}_+ = \delta_{mn}$; $\sum_n$ denotes the sum over discrete values of n , n=0,1,2,3,... and the integral over continuous values. Inserting (3) in the rhs of eq. (I) along with the usual expansion of the electromagnetic potential $\bar{A}(\bar{x})$ in photon operators one obtains

$$H = \sum_n E_n a_n^+ a_n + \sum_m \sum_n P_{mn} a_m^+ a_n + T + \sum_{\bar{k}} \kappa c_{\bar{k}}^+ c_{\bar{k}}, \qquad (4)$$

$$T = \sum_m \sum_n \sum_{\bar{k}} [T_{mn}(\bar{k}) a_m^+ a_n c_{\bar{k}} + T_{mn}^+(\bar{k}) a_m^+ a_n c_{\bar{k}}^+]. \qquad (5)$$

Here $T_{mn}^+(\bar{k}) = T_{nm}^*(\bar{k})$ and $T_{mn} \sim e$. The index $\bar{k}$ denotes the photon momentum along with the photon polarization λ and $\sum_{\bar{k}} \equiv \int d^3k \sum_\lambda$.

The hamiltonian H commutes with $N = \sum_n a_n^+ a_n$, the operator of the number of electrons. So the model can be considered in the sectors of the total Fock space which correspond to fixed N values. In each such sector the model is equivalent to the first-quantized theory of N electrons, interacting with photons (cf. [9], ch.6; [10], ch. XIII, §46).

Let us represent $\sum_m \sum_n T_{mn} a_m^+ a_n$ as the product $a^+ T a$ of the row $a^+ = (a_0^+, a_1^+, \ldots)$, matrix T with elements T_{mn} and the column a. Then, H can be written as

$$H = a^+ E a + a^+ P a + \sum_{\bar{k}} \kappa c_{\bar{k}}^+ c_{\bar{k}} + \sum_{\bar{k}} [a^+ T_{\bar{k}} a c_{\bar{k}} + a^+ T_{\bar{k}}^+ a c_{\bar{k}}^+]. \qquad (6)$$

Here E is the diagonal matrix: $E_{mn} = E_n \delta_{mn}$.

2.2. "Dressed" creation- destruction operators

Consider the transition m $\longrightarrow$ nk from the state "electron is in the state $|m\rangle = a_m^+ \Omega_0$ with the energy E_m, no photons"

into the state" electron is in (n) , one photon with the energy K". The corresponding amplitude $T^{*}_{nm}(\bar{k})$ is nonzero even if the energy is not conserved: $E_m \neq E_n + K$, the case $E_m < E_n$ being the example. Here "energy" means the proper value of the free part $H_0 = a^{+} E a + S_{\bar{k}} K c^{+}_{\bar{k}} c_{\bar{k}}$ of the total Hamiltonian. The corresponding probability in unit time interval (the probability of the transition for a large time interval t divided by t) is zero in the case of energy nonconservation, e.g., see /10/ ch. VIII, § 29. Let us call such transitions (and the corresponding elements of T) "virtual".

In order to deliver the theory from unphysical virtual transitions, one must remove from H the corresponding interaction terms replacing them by others corresponding to physical processes. For this purpose, new operators α_n and $\gamma_{\bar{k}}$ will be introduced instead of the old a_n and $c_{\bar{k}}$ so that the Hamiltonian H when expressed in terms of α , γ would not give rise to the virtual transitions. Let the old operators be expressed in terms of the new ones by means of a unitary transformation

$$a_n = W \alpha_n W^{+} \quad , \quad c_{\bar{k}} = W \gamma_{\bar{k}} W^{+} \tag{7}$$

(here W is a function of the new operators). Then, α, γ would satisfy the same commutation relations as a , c do. Replacing a, c in eq. (6) by the r.h.s. of eq. (7) we find the expression of H in terms of α, γ (let us denote the expression by K)

$$H(a,c) = H(W\alpha W^{+}, W\gamma W^{+}) = W\, H(\alpha,\gamma)\, W^{+} \equiv K(\alpha,\gamma).$$

The relations of the kind $Wa^{+}W^{+}WaW^{+} = Wa^{+}aW^{+}$ were used. Eq. (2) means that in order to get $K(\alpha,\gamma)$ one must substitute α, γ for a, c in eq. (6) and then caloulate the produot $W\, H(\alpha,\gamma)\, W^{+}$.

The following W will suit our purpose

$$W = \exp i S_m S_n S_{\bar{k}} \left[Z_{mn}(\bar{k}) \alpha^{+}_m \alpha_n \gamma_{\bar{k}} + h.c. \right] =$$

$$= \exp i S_{\bar{k}} \left[\alpha^{+} Z_{\bar{k}} \alpha \gamma_{\bar{k}} + \alpha^{+} Z^{+}_{\bar{k}} \alpha \gamma^{+}_{\bar{k}} \right] \equiv \exp A . \tag{8}$$

$Z_{mn}(\bar{k})$ is supposed to be of the order e^{1} along with $T_{mn}(\bar{k})$. To calculate WHW^{+} , one may use the formula

$$e^{A} H e^{-A} = H + [A,H] + \frac{1}{2}[A,[A,H]] + \frac{1}{6}[A,[A,[A,H]]] + \ldots \qquad (9)$$

The following two kinds of commutators occur when calculating WHW^{+}. The first one is

$$[\alpha^{+} Z_{\vec{k}} \alpha \gamma_{\vec{k}}^{-}, \alpha^{+} E \alpha] = [\alpha^{+} Z_{\vec{k}} \alpha, \alpha^{+} E \alpha] \gamma_{\vec{k}}^{-}. \qquad (10)$$

The result of the calculation of the commutator in the r.h.s. of (10) can be represented as $\alpha^{+}[Z_{\vec{k}}, E]\alpha$, where $[Z_{\vec{k}}, E] = Z_{\vec{k}} E - E Z_{\vec{k}}$ is the matrix with the elements

$$[Z_{\vec{k}}, E]_{mn} = Z_{mn}(\vec{k})(E_n - E_m).$$

The second kind occurs, e.g., when calculating $W\alpha^{+} T_{\vec{k}}^{+} \alpha \gamma_{\vec{k}}^{+} W^{+}$:

$$[\alpha^{+} Z_{\vec{k}} \alpha \gamma_{\vec{k}}^{-}, \alpha^{+} T_{\vec{k}'}^{+} \alpha \gamma_{\vec{k}'}^{+}] = \alpha^{+} Z_{\vec{k}} \alpha \; \alpha^{+} T_{\vec{k}'}^{+} \alpha \delta_{\vec{k},\vec{k}'} + \alpha^{+} [Z_{\vec{k}}, T_{\vec{k}'}^{+}] \alpha \gamma_{\vec{k}'}^{+} \gamma_{\vec{k}}^{-}. \qquad (11)$$

After the normal ordering the first term in the r.h.s. eq. (11) becomes

$$-\alpha_{m}^{+} \alpha_{n}^{+}, \alpha_{n} \alpha_{n'}, Z_{mn}(\vec{k}) T_{m'n'}^{+}(\vec{k}) \delta_{\vec{k},\vec{k}'} + \alpha^{+} Z_{\vec{k}} \; T_{\vec{k}'}^{+} \alpha \; \delta_{\vec{k},\vec{k}'}. \qquad (12)$$

Here summation over m,m',n,n' is implied; $Z_{\vec{k}} \; T_{\vec{k}}^{+}$ is the matrix product of $Z_{\vec{k}}$ and $T_{\vec{k}}^{+}$.

The final result for $K = e^{A} H e^{-A}$ can be represented as

$$K = K_2 + K_3 + K_4 + \ldots \qquad (13)$$

Here K_2 is a sum of bilinear terms of the kind $\alpha^{+}\alpha$, $\gamma^{+}\gamma$ including terms of the order e^2, originating from (11), see the last term in (12); K_3 is trilinear and of the order e^1; K_4 denotes tetralinear terms $\sim e^2$. The explicit expressions for K_2, K_3, K_4 are given below.

2.3. Trilinear terms K_3

Let us show that the chosen transformation (8) is capable of removing from K the virtual trilinear interaction terms which are present in K_3 :

$$K_3 = S_{\bar{k}}\, \alpha^+ \{ T_{\bar{k}} + i[Z_{\bar{k}}, E] + i\kappa Z_{\bar{k}} \} \alpha \gamma_{\bar{k}} + h.c. \qquad (14)$$

It is implied that $P = V - V_0$ is of the order e^2 and, therefore, the term $\alpha^+[Z_{\bar{k}}, P]\alpha\gamma_{\bar{k}}$ is not included in (14). The matrix elements

$$T_{mn}(\bar{k}) + i Z_{mn}(\bar{k})(E_n - E_m + \kappa) \equiv \mathcal{T}_{mn}(\bar{k}) \qquad (15)$$

of the curly brackets in (14) can be made zero by setting

$$Z_{mn}(\bar{k}) = i\, T_{mn}(\bar{k}) / (E_n - E_m + \kappa). \qquad (16)$$

In contrast with the models considered in /7,8/ the denominator in (16) can be zero or small. Then, the relative $Z_{mn}(\bar{k})$ cannot be considered as being of the order e^1, as we have supposed. If Z is large, then one should consider along with (14) the trilinear terms resulting from normal ordering of pentalinear (and other multilinear) terms, see below subsection 2.5.

If $E_m < E_n$, then the denominator in (16) exceeds $\omega_{nm} = E_n - E_m$. Let us assume that $T_{mn}(\bar{k})\,\omega_{nm}^{-1} \sim e^1$. Then, $\mathcal{T}_{mn}(\bar{k})$ can be made zero in this case which corresponds to the strongly virtual transition. Suppose also that $T_{nn} = \mathcal{T}_{nn} = 0$. I use below the mnemonic notation $\mathcal{T}_{du} = 0$, d and u denoting "down" and "up".

If $E_m > E_n$, then the denominator module can be, e.g., smaller than the "transition width" Γ_{mn}, which corresponds to the probability of the transition $m \longrightarrow nk$, see /11/, §12. In a typical case, one has $\Gamma_{mn} \sim \omega_{mn} \sim e^2 (\omega_{mn}\, d_{mn})^2 \sim 10^{-6}$. (see /11/, d_{mn} is the transition dipole moment). If $|\kappa - \omega_{mn}| \sim \Gamma_{mn}$, then $Z_{mn} = i\, T_{mn}(\kappa - \omega_{mn})^{-1}$ is of the order $e/e^2 \sim e^{-1}$. If only $Z_{mn} \sim e^1$ is allowed, then it follows from (15) that $\mathcal{T}_{mn}(\bar{k})$ is equal exactly to $T_{mn}(\bar{k})$ at $\kappa = \omega_{mn}$. Moreover, $\mathcal{T}_{mn}(\bar{k}) \cong T_{mn}(\bar{k})$ if $|\kappa - \omega_{mn}| < \varepsilon_{mn}$, provided $\Gamma_{mn} \ll \varepsilon_{mn} \ll \omega_{mn}$. For example, one can let $\varepsilon_{mn} = \omega_{mn}\,(\omega_{mn}\, d_{mn})^2$. To simplify the following exposition, let us assume the Ansatz:

$\mathcal{T}_{mn}(\bar{k}) = 0$ at $|\kappa - \omega_{mn}| > \varepsilon_{mn}$, i.e., in the case of sizable energy nonconservation in the transition $u \longrightarrow dk$ (though the nonconservation is much less than in the strongly virtual transition $d \longrightarrow uk$).

2.4. Radiative correction to the electron Hamiltonian

Consider the bilinear terms of K

$$K_2 = \alpha^+ E \alpha + \alpha^+ P \alpha + S_{\bar{k}}\, \kappa \gamma_{\bar{k}}^+ \gamma_{\bar{k}} + \qquad (17)$$

$$+ \alpha^+ \{ S_{\bar{k}}\, Z_{\bar{k}} \left(-\tfrac{1}{2}[Z_{\bar{k}}^+, E] + \tfrac{1}{2}\kappa Z_{\bar{k}}^+ + i T_{\bar{k}}^+ \right) + h.c. \} \alpha .$$

The second line has its origin from the normal ordering of tetralinear terms in $K = WHW^{+}$, see (12). It is convenient to write the matrix elements $\mathcal{U}_{mn}$ of the curly brackets in eq.(17) by using $\mathcal{T}$ instead of $\mathcal{Z}$, see (15)

$$\mathcal{U}_{mn} = -\frac{1}{2}\sum_{\bar{k}}\sum_{\ell}\Big\{\frac{1}{\kappa-\omega_{n\ell}}\big[T_{m\ell}(\bar{k})+\mathcal{T}_{m\ell}(\bar{k})\big]\big[T^{+}_{\ell n}(\bar{k})-\mathcal{T}^{+}_{\ell n}(\bar{k})\big]+ \\ +\frac{1}{\kappa-\omega_{m\ell}}\big[T_{m\ell}(\bar{k})-\mathcal{T}_{m\ell}(\bar{k})\big]\big[T^{+}_{\ell n}(\bar{k})+\mathcal{T}^{+}_{\ell n}(\bar{k})\big]\Big\}. \tag{18}$$

The second line in eq.(17), i.e., $\alpha^{+}\mathcal{U}\alpha$, is of the same character as $\alpha^{+}P\alpha$ and gives rise to the nonzero contribution of the order e^{2} to the amplitude of the transition from the state "$|n\rangle$, no photons", to the state "$|m\rangle$, no photons". In particular, it gives rise to an additional electron scattering as compared to the scattering which is due to V_0 and P. The residual trilinear interaction

$$\mathcal{T} = \sum_{\bar{k}}\alpha^{+}\mathcal{T}_{\bar{k}}\alpha\gamma_{\bar{k}} + h.c. \tag{19}$$

can also give rise to transitions of that type in the second order of the perturbation theory. Let us show that the matrix elements $\mathcal{T}_{mn}(\bar{k})$ specified in subsection 2.3 give zero transition amplitudes for virtual transitions $n \to m, n \neq m$. These amplitudes are given by eq. (29.19) from /10/

$$a^{(2)}_{mn}(t) = \sum_{\alpha}\frac{\langle m|\mathcal{T}|\alpha\rangle\langle\alpha|\mathcal{T}|n\rangle}{W_\alpha - E_n}\Big[\frac{e^{i(E_m-E_n)t}-1}{E_m-E_n} - \frac{e^{i(E_m-W_\alpha)t}-1}{E_m-W_\alpha}\Big]. \tag{20}$$

Here $|m\rangle = \alpha^{+}_{m}\Omega_0$, $|\alpha\rangle = \alpha^{+}_{\ell}\gamma^{+}_{\bar{k}}\Omega_0$; $\sum_\alpha$ is $\sum_{\ell}\sum_{\bar{k}}$, $W_\alpha = E_\ell + K$. Using (19) one obtains

$$a^{(2)}_{mn}(t) = S_\ell S_{\bar{k}} \frac{\mathcal{T}_{m\ell}(\bar{k})\mathcal{T}^{+}_{\ell n}(\bar{k})}{E_\ell + K - E_n}[\dots]. \qquad (21)$$

Due to $\mathcal{T}_{dd} = \mathcal{T}^{+}_{dd} = 0$, see subsection 2.3, one has $\ell < \min(m,n)$. When $m \neq n$, the surviving terms of the sum S_ℓ vanish because the supports of $\mathcal{T}_{m\ell}(\bar{k})$ and $\mathcal{T}^{+}_{\ell n}(\bar{k})$ do not intersect. Indeed, $\mathcal{T}_{m\ell}(\bar{k}) \neq 0$ at $|K - \omega_{m\ell}| < \varepsilon_{m\ell}$, $\mathcal{T}^{+}_{\ell n}(\bar{k}) \neq 0$ at $|K - \omega_{\ell n}| < \varepsilon_{\ell n}$ and the distance ω_{mn} between $\omega_{m\ell}$ and $\omega_{n\ell}$ is assumed to be much larger than $\varepsilon_{m\ell}$ and $\varepsilon_{\ell n}$.

Consider the shift of the energy of the state $\alpha^{+}_n \Omega_0$ which is due to $\mathcal{T}$ in the order e^2 :

$$\Delta_{\mathcal{T}}(n) = P \Sigma_\alpha \frac{\langle n|\mathcal{T}|\alpha\rangle\langle\alpha|\mathcal{T}|n\rangle}{W_\alpha - E_n} = \qquad (22)$$

$$= -S_\ell P S_{\bar{k}} |\mathcal{T}_{\ell n}(\bar{k})|^2 (K - \omega_{n\ell})^{-1},$$

e.g., see /1/ §35, and §53; /11/, §13 . For the ground level n=0 one obtains $\Delta_{\mathcal{T}}(0) = 0$ becoause $\mathcal{T}_{0\ell} = 0$ for all ℓ (remember that $\mathcal{T}_{dd} = 0$). For the same reason one has for n=1

$$\Delta_{\mathcal{T}}(1) = - P S_{\bar{k}} |\mathcal{T}_{10}(\bar{k})|^2 (K - \omega_{10})^{-1}. \qquad (23)$$

The shift $\Delta_{\mathcal{T}}(n)$ for n=2,3,... in represented by the sum of several terms ($\ell = 0,1,\dots,n-1$), each term being similar to (23). Due to Ansatz, see subsection 2.3, each such term can be represented as follows:

$$\int_{\omega_{n\ell} - \varepsilon_{n\ell}}^{\omega_{n\ell} - \delta} \frac{K^2 dK}{K - \omega_{n\ell}} \int d\Omega_k \Sigma_\lambda |\mathcal{T}_{n\ell}(\bar{k},\lambda)|^2 + \qquad (24)$$

$$+ \int_{\omega_{n\ell} + \delta}^{\omega_{n\ell} + \varepsilon_{n\ell}} \frac{K^2 dK}{K - \omega_{n\ell}} \int d\Omega_k |\mathcal{T}_{n\ell}(\bar{k},\lambda)|^2.$$

Here $\delta \to 0$ originates from the principal value prescription P in eqs. (22),(23). The first integral in (24) is negative, the second one is positive. This allows one to make $\Delta_{\mathcal{T}}(n)$ zero. For this purpose, let $\mathcal{T}_{n\ell}(\bar{k})$ be equal to $T_{n\ell}(\bar{k})$ in that term of (24) which has the least module. Suppose this is the first term. In the second term of (24) let $\mathcal{T}_{n\ell}(\bar{k})$ be equal to $T_{n\ell}(\bar{k})$ in an interval $(\omega_{n\ell}+\delta, \omega_{n\ell}+\varepsilon')$, $\varepsilon' < \varepsilon_{n\ell}$ (and be zero outside it). One has to choose ε' so that the term module would be equal to the first term module.

Now I return to the disoussion of eq.(18). Note that the terms of the r.h.s. of (18), which contain $\mathcal{T}_{m\ell}(\bar{k})\mathcal{T}^{+}_{\ell n}(\bar{k})$, are zero on the same ground as $\mathcal{U}^{(2)}_{mn}(t)$ and $\Delta_{\mathcal{T}}(n)$ are. The terms containing $T_{m\ell}(\bar{k})\mathcal{T}^{+}_{\ell n}(\bar{k})$ are zero if n=0 because of $\mathcal{T}_{\alpha\alpha}=0$. In the case $n > 0$, the corresponding integrals of these terms are expected to be much smaller than the integrals of $T_{m\ell}(\bar{k})T^{+}_{\ell n}(\bar{k})$. The reason is that the integrals of $T_{m\ell}(\bar{k})\mathcal{T}^{+}_{\ell n}(\bar{k})$ over $|\bar{k}|$ are taken over small intervals $\omega_{n\ell}-\varepsilon_{n\ell} < |\bar{k}| < \omega_{n\ell}+\varepsilon_{n\ell}$, cf. (24). In particular, all terms in eq. (18) which contain $\mathcal{T}$ do not surely contribute to the possible divergencies of $\mathcal{U}_{mn}$. So

$$\mathcal{U}_{mn} \cong -S_\ell P S_{\bar{k}} \frac{1}{2}\left[\frac{1}{\kappa-\omega_{n\ell}} + \frac{1}{\kappa-\omega_{m\ell}}\right] T_{m\ell}(\bar{k})\, T^{+}_{\ell n}(\bar{k}). \qquad (25)$$

Let us comment on the appearance of the principal value P in (25). The r.h.s. of eq. (18) has no singularities at $\kappa = \omega_{n\ell}$ and $\kappa = \omega_{m\ell}$ because according to Ansatz the differences $T^{+}_{\ell n}(\bar{k}) - \mathcal{T}^{+}_{\ell n}(\bar{k})$ and $T_{m\ell}(\bar{k}) - \mathcal{T}_{m\ell}(\bar{k})$ vanish at these values of k. Therefore, (18) does not alter if $\int_0^\infty d\kappa$ is understood as $P\int_0^\infty d\kappa$. The introduction of the principal value P allows one to give sense to the separate terms of eq.(18) which have the above singularities, the r.h.s. of eq. (25) or $S_\ell P S_{\bar{k}}\, T_{m\ell}\, \mathcal{T}^{+}_{\ell n}(\kappa-\omega_{m\ell})^{-1}$ being the example.

Note that the transformation W, eq. (8), with an arbitrary Z (not satisfying eq. (16), i.e., $\mathcal{T}$ being also arbitrary) leads to

$$\mathcal{U}_{nn} = -S_\ell S_{\bar{k}} \left\{ |T_{n\ell}(\bar{k})|^2 - |\mathcal{T}_{n\ell}(\bar{k})|^2 \right\} (\kappa-\omega_{n\ell})^{-1}. \qquad (26)$$

Like P_{nn} this is a correction to the energy E_n of the electron state $|n\rangle$. Using (22) one has

$$\mathcal{U}_{nn} + \Delta_{\mathcal{T}}(n) = -S_e \, P S_{\bar{k}} \, | \, T_{en}(\bar{k}) |^2 (k - \omega_{ne})^{-1} \equiv \Delta_T(n). \qquad (27)$$

Here $\Delta_T(n)$ is the well known level shift $\sim e^2$, resulting from the original interaction T, eq. (5). So the sum $\mathcal{U}_{nn} + \Delta_{\mathcal{T}}(n)$ is invariant under the transformation W. With $\mathcal{T}$ chosen as above one has $\Delta_{\mathcal{T}}(n) = 0$ and $\mathcal{U}_{nn}$ becomes the total radiative correction of the order e^2 to E_n.

The term $\alpha^+ \mathcal{U} \alpha$ can be called the radiative correction to the electron Hamiltonian $\alpha^+ E \alpha + \alpha^+ P \alpha$. It can be shown that the most divergent part of $\mathcal{U}_{mn}$ is proportional to $(\bar{p})^2_{mn}$, $\bar{p}$ being the electron momentum operator (the demonstration is similar to the nonrelativistic calculation of $\Delta_T(n)$, see, e.g., [12]). So, the name "potential" or even "quasi-potential" is unsuitable for $\alpha^+ \mathcal{U} \alpha$. The term $\alpha^+ \mathcal{U} \alpha$ gives also all radiative corrections of the order e^2 to the electron scattering on the external potential V and to the electron level energies.

2.5. Tetralinear and other interaction terms

Using the matrix $i(T_{\bar{k}} + \mathcal{T}_{\bar{k}})/2$ instead of the combination $-\frac{i}{2}[Z_{\bar{k}}, E] - \frac{K}{2} Z_{\bar{k}} + i T_{\bar{k}}$, see (15), one can write the terms K_4 from eq. (13) in the form

$$K_4 = \alpha^+ \left[Z_{\bar{k}}, \tfrac{i}{2}(T_{\bar{k}'} + \mathcal{T}_{\bar{k}'}) \right] \alpha \, \gamma_{\bar{k}} \gamma_{\bar{k}'} + h.c. +$$

$$+ \alpha^+ \left[Z_{\bar{k}}, \tfrac{i}{2}(T^+_{\bar{k}'} + \mathcal{T}^+_{\bar{k}'}) \right] \alpha \, \gamma^+_{\bar{k}'} \gamma_{\bar{k}} + h.c. + \qquad (28)$$

$$+ \alpha^+_m \alpha^+_{m'} \alpha_n \alpha_{n'} \left\{ Z_{mn}(\bar{k}) \left(-\tfrac{i}{2}\right) (T^+_{\bar{k}} + \mathcal{T}^+_{\kappa}) \right\}_{m'n'} + h.c. \Big\}.$$

Summations over all repeated indices are implied. The terms of the kind $\alpha^+ \alpha \gamma^+ \gamma$ describe the photon scattering on the electron. The terms $\alpha^+ \alpha \gamma \gamma + h.c.$ give rise to the virtual transition of the electron from the ground state $\alpha^+_0 \Omega_0$ to an excited one with emission of two photons. The terms $\alpha^+ \alpha^+ \alpha \alpha$ contribute to the electron-electron scattering in the two-electron sector (they vanish in the one-electron sector).

In the simplified consideration presented here only several first terms of the infinite series for $K = W H W^+$ were discussed.

The triple and higher repeated commutators from eq.(9) give rise to pentalinear terms of the order e^3 and other multilinear terms. Normal ordering of $\alpha^+\alpha\gamma\gamma\gamma^+$ or $\alpha^+\alpha\alpha^+\alpha\gamma$ gives trilinear terms of the order e^3 which were not taken into account here.

3. Radiative correction to the particle-particle interaction

One can add to the model Hamiltonian (I) the Coulomb interaction of electrons

$$I_c = \frac{1}{2}\int d^3x \int d^3y\, \psi^+(\bar{x})\psi^+(\bar{y})\frac{e^2}{4\pi|\bar{x}-\bar{y}|}\psi(\bar{y})\psi(\bar{x}). \tag{29}$$

The last line in eq.(28) is similar to (29) and can be named the radiative correction to I_c which results from the transversal photon exchange.

Now assume for simplicity that $V_0 = P = 0$. Then, one can take the electron momenta $\bar{p}$ as the indices n. One can show that the main part of the radiative correction renormalizes the electron mass: $p^2/2\mu_0$ in eq.(I) is replaced by $p^2/2\mu$, where μ corresponds to the observable electron mass [a]. Therefore $\alpha^+ U\alpha + \alpha^+ \mathcal{U}\alpha$ assumes the form $\int d^3p\, \alpha^+_{\bar{p}}(p^2/2\mu)\alpha_{\bar{p}}$. Let us collect together all the terms from K_2 and K_4, eq. (28), which contain the electron operators only and call them K_{part}

$$K_{part} = \alpha^+_{\bar{p}}\frac{p^2}{2\mu}\alpha_{\bar{p}} + \alpha^+_{\bar{p}_1}\alpha^+_{\bar{p}_2} I(\bar{p}_1,\bar{p}_2;\bar{p}_1',\bar{p}_2')\alpha_{\bar{p}_1'}\alpha_{\bar{p}_2'}. \tag{30}$$

The integration over $\bar{p}, \bar{p}_1, \bar{p}_2, \bar{p}_1', \bar{p}_2'$ is implied; I is the sum of (29) and the last line from eq.(28), In the two-electron sector, the second-quantized operator (30) is equivalent to the first-quantized operator having the elements

$$\langle \bar{p}_1, \bar{p}_2 \mid K_{part} \mid \bar{p}_1', \bar{p}_2' \rangle =$$

$$= \left(\frac{p_1^2}{2\mu} + \frac{p_2^2}{2\mu}\right)\delta(\bar{p}_1-\bar{p}_1')\,\delta(\bar{p}_2-\bar{p}_2') + 2I(\bar{p}_1,\bar{p}_2;\bar{p}_1'\bar{p}_2'). \tag{31}$$

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a) It is implied that a cut-off is introduced into the interaction term in eq.(I) and the main contribution to $\alpha^+\mathcal{U}\alpha$, which is proportional to p^2, is small as compared with $p^2/2\mu_0$. In addition, $\alpha^+\mathcal{U}\alpha$ contains, e.g., a smaller term proportional to p^4.

Here $|\bar{p}_1, \bar{p}_2\rangle = \frac{1}{\sqrt{2}} \alpha^+_{\bar{p}_1} \alpha^+_{\bar{p}_2} \Omega_0$. The equation for proper vectors of the operator (31) is

$$\left(\frac{p_1^2}{2\mu} + \frac{p_2^2}{2\mu} \right) \varphi_E(\bar{p}_1, \bar{p}_2) + \int d^3p_1' \int d^3p_2' \, 2I(\bar{p}_1, \bar{p}_2; \bar{p}_1', \bar{p}_2') \varphi_E(\bar{p}_1', \bar{p}_2') = E \varphi_E(\bar{p}_1; \bar{p}_2). \quad (32)$$

The second-quantized expression for the proper vector φ_E is

$$\varphi_E = \int d^3p_1 \int d^3p_2 \, \varphi_E(\bar{p}_1, \bar{p}_2) \frac{1}{\sqrt{2}} \alpha^+_{\bar{p}_1} \alpha^+_{\bar{p}_2} \Omega_0. \quad (33)$$

The vector φ_E can be used to describe the initial and final two-electron states instead of $\alpha^+_{\bar{p}_1} \alpha^+_{\bar{p}_2} \Omega_0$. In other words, (33) can be used as the zeroth approximation functions for the perturbation theory with the interaction $K - K_{part} - K_{ph}$. This approach allows one, e.g., to take exactly into account the Coulomb interaction of electrons.

In the case of two electrons, there are no proper vectors of K_{part} which would correspond to bound states. Our model must be enlarged in order to discuss two-particle bound states. One can consider besides the electron field ψ another nonrelativistic field φ describing the positively charged particles (protons, μ^+ mesons) which have an analogous trilinear interaction with photons. The relevant generalized unitary transformation W would give rise to an operator K_{part} containing the negative and positive particle creation-destruction operators. In the sector "one electron, one proton", one would obtain the equation of the kind (32) with different masses μ_e and μ_p . It would have solutions describing bound states (of the hydrogen atom).

To be able to consider the positronium, one needs the relativistic spinor QED. Its Hamiltonian also has trilinear interaction term, Coulomb interaction being added in the Coulomb gauge. Both these interactions now give rise to the electron-positron pair creation and annihilation, so that the number of electrons is not conserved. The papers [7,8] have dealt with such a case and have provided a suitable transformation W which removes virtual interaction terms from $K = WHW^+$. Let us consider that part K_{part} of K which contains only electron-positron operators. Now K_{part} gives rise to pair creation. To obtain the first-quantized equation of the kind (32) for the electron-positron bound states, one can project this K_{part} onto the corresponding sector "one electron, one positron" by eliminating

from K_{part} the terms creating pairs, see, e.g., /13/. So only a part of K_{part} (which was called $K_{part}^{no\ pair}$ in /13/) should be used for determination of the positronium states.

Another way is possible which is in the line of our "dressing". One must perform an additional unitary transformation $W_4 = \exp A_4$, where $A_4 \sim e^2$ and is tetralinear in the electron-positron operators $\alpha, \beta, \alpha^+, \beta^+$. The task of W_4 is to remove from K_{part} the tetralinear terms giving rise to pair creation, e.g., $\alpha^+\beta^+\alpha^+\alpha$. The remaining tetralinear terms would commute with operators of the electron and positron numbers and can play the role of the second term in the r.h.s. of eq. (30).

Let us note in conclusion that one can add to K_{part}, eq. (30), sextalinear terms of the kind $\alpha^+\alpha^+\alpha^+\alpha\alpha\alpha$ resulting from higher commutators in eq.(9). These terms describe three-particle interactions irreducible to the two-particle interactions, discussed above. In the three-particle sector, one would obtain the equation for proper vectors $\varphi_E(\bar{p}_1, \bar{p}_2, \bar{p}_3)$ which takes into account these irreducible three-particle forces.

4. Radiative potential-like correction in the multipolar form of QED

The radiative correction $\alpha^+ U \alpha$, considered in sect.2, can be used to discuss a trouble with the multipolar form (MF) of QED. The form has been suggested by Power and Zienau /14/ and was developed in subsequent papers, e.g., see references 2 - 9 cited in /15/. In MF the interaction term may be expanded in a series of electric and magnetic multipoles (moments) of atoms or molecules. MF has calculational advantages and is widely used, see references 10-16, cited in /15/.

But the Hamiltonian of MF has a specific quadratically divergent potential-like term which is denoted here by $P_1(\bar{x})$, see eq.(35) below. A regularized modification of MF has been suggested in /15/ in which P_1 becomes a finite potential. But then it reveals its another deficiency: it is confinement potential, see /16/ sect.2.4 and "Note added in proof" in the end of sect.3.2 in /15/. As is known, there is no confinement in the realm of electromagnetic phenomena.

It will be shown in this section that the radiative correction $\alpha^+ U \alpha$ contains a part which cancels P_1.

In the long-length wave approximation, the second-quantized version of the nonrelativistic MF Hamiltonian is of the form (e.g., cf. /15/)

$$H = \int d^3x\, \psi^+(\vec{x}) \Big\{ \frac{\hat{\vec{p}}^2}{2\mu_0} + V_0(\vec{x}) + P_\perp(\vec{x}) + e \int_0^{\vec{x}} d\vec{\ell}\cdot\vec{E}(\vec{\ell}) \Big\} \psi(\vec{x}) + H_{ph}. \tag{34}$$

I consider here only this nonregularized version, the regularized one can be dealt with in a similar way. The last term in the curly brackets in (34) contains the integral $\int d\vec{\ell}\cdot\vec{E}(\vec{\ell})$ taken along the straight line connecting the centre of the potential V_0 with the point $\vec{x}$, the electron coordinate.

$$P_\perp(\vec{x}) = \frac{e^2}{2} \sum_{m,n} \int_0^{\vec{x}} d\ell_m \int_0^{\vec{x}} d\ell'_n \, \delta^\perp_{mn}(\vec{\ell}-\vec{\ell}'), \quad m,n = 1,2,3. \tag{35}$$

Here $\delta^\perp$ is the transversal part of the δ-function, e.g., see /15/. Only the electric-dipole part of the MF interaction is considered because just this part gives the radiative correction cancelling $P_\perp$. Using eq. (3) and standard expansion of the transverse electric field $\vec{E}(\vec{x})$ in photon operators one can reduce (34) to the form (4) with $P = P_\perp$ and

$$T_{mn}(\vec{k},\lambda) = \int d^3x\, \varphi_m^*(x) \Big[ie(2\pi)^{-3/2} \sqrt{\tfrac{k}{2}} \int_0^{\vec{x}} d\vec{\ell}\cdot\vec{\mathcal{E}}_\lambda(\vec{k}) e^{i\vec{k}\vec{\ell}} \Big] \varphi_n(\vec{x}).$$

It can be shown that

$$(P_\perp)_{mn} = \int d^3k \sum_\lambda \frac{1}{k} S_e T_{me}(\vec{k},\lambda) T^+_{en}(\vec{k},\lambda). \tag{36}$$

Now let us use eq. (25) for $\mathcal{U}_{mn}$. The identity $(\kappa-\omega)^{-1} = k^{-1} + \omega\kappa^{-1}(k-\omega)^{-1}$ allows one to separate from $\mathcal{U}_{mn}$ the part

$$-S_e P \int d^3k \sum_\lambda \frac{1}{k} T_{me}(\vec{k},\lambda) T^+_{en}(\vec{k},\lambda),$$

which is equal to $-(P_\perp)_{mn}$. So

$$\alpha^+_m (E + P_\perp + \mathcal{U})_{mn} \alpha_n = \alpha^+_m (E_{mn} + \mathcal{U}'_{mn}) \alpha_n,$$

$$\mathcal{U}'_{mn} = -S_e P S_{\vec{k}} \frac{1}{2} \Big[\frac{\omega_{ne}}{k(k-\omega_{ne})} + \frac{\omega_{me}}{k(k-\omega_{me})} \Big] T_{me}(\vec{k}) T^+_{en}(\vec{k}). \tag{37}$$

This means, in particular, that the shift $\mathcal{U}_{nn}$ of the energy of the

state $\alpha_n^+ \Omega_0$ has a contribution (which is the most divergent part of $\mathcal{U}_{nn}$) which cancels the level shift resulting from P_1 in the first order of perturbation theory. Just this cancellation has been proven in [17], see eqs. (6-51) – (6-66) there. It has been shown above that $\mathcal{U}_{mn}$ contains the contribution which calcels the potential P_1 as a whole, not only its diagonal elements.

Note that the cancellation occurs in the Hamiltonian $K = WHW^+$ containing besides residual trilinear terms $\mathcal{T}$ also tetralinear and other interaction terms. Let us show that the cancellation can be accounted for in the frames of the initial expression (34) for the hamiltonian. For this purpose we use the equivalence theorem stating the equality of two S-matrices, one being calculated by using $K_{int} = K - K_0$, $K_0 = \alpha^+ E \alpha + S_\xi \kappa \gamma_\xi^+ \gamma_\xi$, and the other by using $H_{int} = H - H_0$, $H_0 = a^+ E a - S_\xi \kappa c_\xi^+ c_\xi$. This theorem follows from the fact that "dressed" and "bare" Heisenberg operators tend as $t \to -\infty$ or $t \to +\infty$ to the same "in" or "out" operators [8,18]. Stress that K_{int} contains the term $\alpha^+ \mathcal{U}' \alpha$, see eq.(37), while H_{int} must contain P_1. This means that in order to calculate the S-matrix by using (34), one must treat P_1 as a perturbation term along with the trilinear interaction (though P_1 is surely "larger" that V_0 in all senses). This prescription gives grounds to the handling with P_1 which has been used in [14,17]. The term P_1 plays the role similar to the mass renormalization term $\delta m \bar{\psi} \psi$ in the spinor QED. Its sole function in MF is to cancel (the most divergent) part of the radiative correction to $\alpha^+ E \alpha$.

5. Some consequences of the "dressed" description of particles

New electron and photon operators α, γ, introduced in sect.2, can be represented in terms of the old ones a,c by using eqs. (7) and (9), e.g.,

$$\alpha_n^+ = a_n^+ - i S_\xi S_m [Z_{mn}(\bar{\kappa}) a_m^+ c_\xi^+ + Z_{mn}^+(\bar{\kappa}) a_m^+ c_\xi^+] + \ldots \qquad (38)$$

Now one- and many-particle states can be defined by using α, γ. For example, new one-electron states are described by the vectors $\alpha_n^+ \Omega_0 = W^+ a_n^+ \Omega_0$. One of the consequences of this description is the nonvanishing mean value of the transversal electric field operator in the states

$$\langle \alpha_n^+ \Omega_0, \bar{E}(\bar{x}) \alpha_n^+ \Omega_0 \rangle = \langle a_n^+ \Omega_0, W \bar{E}(\bar{x}) W^+ a_n^+ \Omega_0 \rangle \neq 0$$

unlike $\langle a_n^+\Omega_0, \bar{E}(\vec{x})\, a_n^+\Omega_0\rangle$ which vanish. This means that α_n^+ creates besides the charged particle also a relevant transversal electric field. Remind that in the Coulomb gauge, the one-electron state similarly possesses a nonvanishing electric field but the longitudinal (Coulomb) one, see [19]. By the above reason, the operator α^+ and the particle created by it can be called "dressed".

Note that only a simple partial "dressing" was realized in sect.2. The relevant "dressed" ground state $\alpha_0^+\Omega_0$ turns out to be more stable than the "bare" one $a_0^+\Omega_0$ in the following sense: it is stable under the one-photon emission but it can virtually emit two photons due to the interaction $\alpha^+\alpha\gamma\gamma + h.c$, see eq.(28), the corresponding amplitude being of the order e^2.

Notwithstanding the difference between "dressed" and "bare" states, the above-mentioned fact of the equality of the corresponding S-matrices does hold in the sense specified in the preceding section. In other words, the S-matrix is invariant under the dressing transformation (as well as the total level shift $U_{nn} + \Delta_\gamma(n)$, see sect.2).

There exist observable phenomena which cannot be described by the S-matrix. The most known example is the time evolution of an unstable state. Consider the electron which at t=0 was in the excited "dressed" state $|n\rangle = \alpha_n^+\Omega_0$. The usual definition of the nondecay probability N(t) is

$$N(t) = |\langle n | \exp(-iHt) | n\rangle|^2.$$

N(t) is known to be equal approximately to exp $(-\Gamma t)$, e.g., see [11] §13; [20]. The one-photon transitions into the lower states m give the main contributions to Γ, so $\Gamma = \sum_m \Gamma_{mn}(x)$ at $x = E_n$, where

$$\Gamma_{mn}(x) = 2\pi \int_{\vec{k}} |\mathcal{T}_{nm}(\vec{k})|^2\, \delta(x - E_m + \kappa), \tag{39}$$

see [11], eq. (85). This value for Γ coinoides with the value which has Γ in the "bare" formalism (the initial state being described by $a_n^+\Omega_0$) with the interaction T, eq. (5). This follows from the equality $\mathcal{T}_{nm}(\omega_{nm}) = T_{nm}(\omega_{nm})$, see sect.2.3. But the function $\Gamma_{mn}(x)$ has a finite support in the "dressed" formalism: it follows from eq. (39) that $\Gamma_{mn}(x)$ vanishes if x is outside $(E_n - \mathcal{E}_{mn},\ E_n + \mathcal{E}_{mn})$ because $\mathcal{T}_{nm}(\vec{k})$ vanishes if κ is outside $(\omega_{nm} - \mathcal{E}_{mn},\ \omega_{nm} + \mathcal{E}_{mn})$. This changes the "line shape" P(k,t) which is defined as the probability that photon energy is equal to k, if measured at the moment t in the transition n $\longrightarrow$ mk. If $t \gg 1/\Gamma$, then

$$P(k,t) \cong \frac{1}{2\pi} \Gamma_{mn}(x) \left\{ [x - E_n + \Delta(x)]^2 + \Gamma^2(x)/4 \right\}^{-1}, \quad x \equiv k + E_m, \qquad (40)$$

e.g., see /11/ eq. (113); /20/ eq. (122b). The distribution (40) differs from the Lorentzian one

$$\frac{1}{2\pi} \Gamma_{mn}(E_n) \left\{ (k - \omega_{nm})^2 + \Gamma^2(E_n)/4 \right\}^{-1}$$

by cutting on the wings: (40) vanishes if k is outside the interval $(\omega_{nm} - \varepsilon_{mn},\ \omega_{nm} + \varepsilon_{mn})$. As to the value of ε_{mn}, it satisfies the inequalities $\Gamma_{mn} \ll \varepsilon_{mn} \ll \omega_{nm}$; no further restriction for ε_{mn} was found in sect.2. The cutting becomes more prominent in the distribution P(k,t) at $t < 1/\Gamma$, which is equal to the square of the module of the expression (119b) in /20/. For the explicit formula see /21/. This paper has reported a measurement of P(k,t) at small t by using the Mössbauer effect.

The cutting of $\mathcal{T}_{mn}(\bar{k})$ as compared with $T_{mn}(\bar{k})$ results also in the additional deviations of the exact nondecay law N(t) from $\exp(-\Gamma t)$, e.g., see /20/. The nonexponential behaviour of N(t) has not yet been observed, see the paper /22/ which contains references to theoretical and experimental papers on this topic.

6. Summary

The way of "dressing", suggested in /7,8/, is modified here for the case when "dressed" states can be unstable. The peculiarities of the "dressed" excited atom states were discussed in sect.5.

The "dressing" allows one to define new zeroth approximation $K_{zero} = K_{part} + K_{photon}$ see sect.3) to the total Hamiltonian K which differs from its standard free part K_0 by addition of terms renormalizing the particle mass and (quasi) potential terms. In the case of the electron bounded by an external potential, the latter give rise to radiative level shifts. The determination of K_{zero} allows one to propose a new approach to the problem of finding bound states such as the positronium in QED and hadrons in QCD. These bound states are defined as proper vectors of the Hermitian operator K_{part} in suitable sectors of the theory Hilbert space. The quantitative comparison of the obtained (quasi-)potential terms with analogous terms of other approaches was not the purpose of this paper. Let us mention only some qualitative peculiarities of the proposed approach.

The operators K_{zero} and K_{part} are parts of the total Hamiltonian. This is not the case in other approaches though its (quasi-) potentials are also calculated by using the field Lagrangian or Hamiltonian (more specifically the relevant Feynman diagrams).

When calculating (quasi-)potentials from Feynman diagrams, one may obtain different expressions coinciding on the "energy shell". The nonuniqueness like that is absent in the proposed approach.

The approach does not contain excess variables such as relative time or energy in the Bethe–Salpeter equation.

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