Nontopological solitons in strongly coupled QED

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We review the pattern of symmetry breaking recently suggested for strongly coupled QED and point out that it is quite similar to that generally accepted as operative in the case of QCD. We suggest, therefore, that we can describe "hadronlike" objects in QED using methods developed for QCD. This discussion may be used to support our interpretation of the electron-positron peaks observed in heavy-ion collisions as representing the decay products of nontopological solitons (containing quasielectrons and quasipositrons) formed in a new vacuum phase of QED. We predict, however, that the existence of such objects will not lead to any anomalies in electron-positron scattering with center-of-mass energies of about 1.7 MeV.

I. INTRODUCTION

There has been much interest in the past few years in the observation of narrow electron and positron lines in the spectrum of particles emitted in heavy-ion collisions.[1–6] The interpretation of these peaks has defied analysis. In particular, the fact that several peaks are seen for the same system suggests that we are not seeing the decay of a single elementary object. (The radiation appears to come from a source moving with the velocity of the center of mass of the heavy ions;[6] this is one of the facts which is difficult to explain in a scenario involving new elementary particles.[7,8]

In a recent Letter[9] we suggested that we are seeing the decay of a kind of "pseudopositronium"[10] formed as a nontopological soliton in a new vacuum phase of QED (Ref. 11) created over a region which might have a characteristic radius of about 300 fm. Indeed, if we scaled our results for the spectrum of charmonium, or of b quarkonium,[12] so that we obtained a particle of mass of about 1.7 MeV, we obtained objects of the characteristic size mentioned above, with excited states with a level spacing of about 100 keV (Ref. 9).

For example, let us consider the field equations

$$(i\gamma^\mu\partial_\mu - m_q)\psi(x) = g_\chi\psi(x)\chi(x) \ , \tag{1.1}$$

$$(\Box + m_\chi{}^2)\chi(x) = -g_\chi\bar{\psi}(x)\psi(x) \ , \tag{1.2}$$

which have (nontopological) soliton solutions. (These solutions have been extensively studied in earlier works.[12]) Now introduce the dimensionless fields

$$X(x) = \chi(x)/m_q \ , \tag{1.3}$$

$$P(x) = \psi(x)/m_q{}^{3/2} \ , \tag{1.4}$$

and the dimensionless coordinate $\rho^\mu = m_q x^\mu$. With the definition $\eta^2 = m_\chi{}^2/m_q{}^2$, we have the (dimensionless) equations

$$\left[i\gamma^\mu \frac{\partial}{\partial\rho^\mu} - 1 \right] P(\rho) = g_\chi P(\rho) X(\rho) \ , \tag{1.5}$$

and

$$(\Box_\rho + \eta^2) X(\rho) = -g_\chi \bar{P}(\rho) P(\rho) \ . \tag{1.6}$$

Thus we see that the mass of the soliton is given by

$$m = f(g_\chi, \eta) m_q \ , \tag{1.7}$$

where $f(g_\chi,\eta)$ is a dimensionless function. Similarly, the radius of the soliton is given by

$$R = h(g_\chi, \eta)/m_q \ , \tag{1.8}$$

where $h(g_\chi,\eta)$ is also dimensionless. Using this formalism, it is a simple matter to change the mass scale (set by m_q) and make use of previous calculations of meson structure[12] to describe solitons in strong-coupling QED.

In this paper we wish to provide some justification for the use of Eqs. (1.1) and (1.2) in QED. Note that $\psi(x)$, in the case of QCD, is the quark field, and $\chi(x)$ is an order parameter of QCD vacuum condensates. In QED, $\psi(x)$ will represent the electron field and $\chi(x)$ will again be an order parameter, in this case associated with the broken chiral symmetry of QED in the strong-coupling phase. To support this interpretation we will review some recent work concerning the pattern of symmetry breaking in strong-coupling QED (Ref. 13). As we will see, in the strong-coupling limit, QED may contain certain scalar and pseudoscalar modes. The 0^- states are the Goldstone bosons associated with the breaking of chiral symmetry and remain massless. The 0^+ states obtain a mass.[13] A scale associated with these modes is governed by a cutoff parameter Λ. We can think of these modes as bound states of massless electrons and positrons. If Λ is large, the bound state could have a very small radius as compared to the characteristic size of our solitons, R. (That is, we consider the case $R \gg \Lambda^{-1}$.) In that limit, the internal wave function of the lepton condensate pairs is not important for our study of soliton structure. [Thus $\chi(x)$ will be a measure of the condensate strength, that is, the number of condensed pairs, in some region about the point x.] Alternatively, one notes that Λ sets the scale for the momentum dependence of the electron self-energy $\Sigma(p^2)$. The study of soliton dynamics usually involves the introduction of a coordinate-dependent self-energy $\Sigma(x)$, with a simultaneous neglect of the momentum dependence. This is a restatement of the approximation $R \gg \Lambda^{-1}$. *All in all, it is probably less confusing to concentrate on the electron self-energy and discard the reference to condensed pairs since chiral-symmetry breaking will take place in the quenched approximation and can be generated by considering the dynamics of a single electron interacting with a gauge field.* (The condensed fermions have zero mass and are therefore not the relevant quasiparticles of the theory after chiral-symmetry breaking takes place.)

In the next section we discuss one approach to dynamical mass generation in a gauge theory. In Sec. III we indicate how we can describe nontopological solitons using the order parameters of the phase with broken chiral symmetry. Sections IV and V contains some comments on heavy-ion collisions and some further discussion.

II. BILOCAL FIELDS AND THE EFFECTIVE ACTION

In this section we review the recent work of Morozumi and So.[13] These authors write the effective action for QED and proceed to integrate out the photon field. They obtain the action in the Landau gauge and in a Euclidean metric:

$$S[\psi,\bar\psi]=\int d^4x\,\bar\psi(x)\gamma^\mu\partial_\mu\psi(x)+\frac{e^2}{2}\int d^4x\int d^4y\,J_\mu(x)D^{\mu\nu}(x,y)J_\nu(y)\,, \tag{2.1}$$

with

$$J_\mu(x)=i\bar\psi(x)\gamma_\mu\psi(x) \tag{2.2}$$

and

$$D^{\mu\nu}(x-y)=-\int\frac{d^4p}{(2\pi)^4}e^{ip\cdot(x-y)}\times(\delta^{\mu\nu}p^2-p^\mu p^\nu)/p^4\,. \tag{2.3}$$

The next step in the analysis is a Fierz transformation to yield the action

$$S[\psi,\bar\psi]=\int d^4x\,\bar\psi(x)\gamma^\mu\partial_\mu\psi(x)+\tfrac{3}{8}e^2\int d^4x\int d^4y\{[\bar\psi(x)\psi(y)]D_0(x-y)[\bar\psi(y)\psi(x)]+\bar\psi(x)i\gamma_5\psi(y)D_0(x-y)\bar\psi(y)i\gamma_5\psi(x)+\cdots\}\,. \tag{2.4}$$

Here we have not written axial-vector, vector, and tensor terms which also appear. Note that

$$D_0(x-y)=-\int\frac{d^4p}{(2\pi)^4}\frac{e^{ip\cdot(x-y)}}{p^2} \tag{2.5}$$

$$=-\frac{1}{4\pi^2}\frac{1}{(x-y)^2}\,. \tag{2.6}$$

Further, auxiliary bilocal fields $\sigma(x,y)$ and $\pi(x,y)$, are introduced without changing the value of the action.[14] One adds to the action, with $\lambda\equiv\frac{3}{4}e^2$,

$$S_A=-\frac{1}{2\lambda}\int d^4x\int d^4y\{[\sigma(x,y)-\lambda\bar\psi(x)\psi(y)D_0(x-y)]D_0^{-1}(x-y)[\sigma(y,x)-\lambda D_0(x-y)\bar\psi(y)\psi(x)]+[\pi(x,y)-\lambda\bar\psi(x)i\gamma_5\psi(y)D_0(x-y)]D_0^{-1}(x-y)[\pi(y,x)-\lambda\bar\psi(y)i\gamma_5\psi(x)D_0(x-y)]+\cdots\}\,. \tag{2.7}$$

At this point the fermion fields are integrated out leaving an effective action $S_{\rm eff}$ expressed as a functional of the bilocal fields. One may write[13]

$$\sigma(x,y)=\int\frac{d^4p}{(2\pi)^4}\int\frac{d^4q}{(2\pi)^4}e^{iP\cdot X+iq\cdot r}\sigma_q(P)\,, \tag{2.8}$$

with

$$r_\mu=x_\mu-y_\mu\,, \tag{2.9}$$

$$X_\mu=\tfrac{1}{2}(x_\mu+y_\mu)\,, \tag{2.10}$$

and a similar equation for $\pi(x,y)$. Again, following Ref. 13, we set

$$\sigma_q(P)=(2\pi)^4\delta^{(4)}(P)\sigma(q^2)\,, \tag{2.11}$$

$$\pi_q(P)=(2\pi)^4\delta^{(4)}(P)\pi(q^2)\,. \tag{2.12}$$

Further, the effective action is evaluated in a two-loop approximation and one sets $\sigma(q^2)=B(q^2)$ (Ref. 13). The equation for $B(q^2)$ is

$$\frac{-4\pi^2}{\lambda}\Box_q B(q^2)-\frac{4B(q^2)}{q^2+B^2(q^2)}=0\,, \tag{2.13}$$

$$\Box_q \equiv \frac{\partial}{\partial q_\mu} \frac{\partial}{\partial q_\mu} \ . \tag{2.14}$$

The boundary conditions are given in Ref. 13, where it is pointed out that Eq. (2.13) is a differential form of an (approximate) Schwinger-Dyson equation whose solutions are known.[15–17] There is a critical point at $\lambda = \lambda_c = \pi^2$. In the weak-coupling phase ($\lambda < \lambda_c$), there is a trivial solution, $B(q^2) = 0$. In the strong-coupling phase ($\lambda \geq \lambda_c$), there are an infinite number of nontrivial solutions. The lowest vacuum energy follows from adopting the *nodeless* solution for $B(q^2)$, $B_0(q^2)$ (Ref. 13).

By expanding the action in terms of the deviation field,

$$\sigma'_q(P) = \sigma_q(P) - (2\pi)^4 \delta^{(4)}(P) B_0(q^2) \ , \tag{2.15}$$

and $\pi_q(P)$, it is shown in Ref. 13 that the σ field is massive and the field π is massless for $\lambda > \lambda_c$. (At the critical point both fields are massless.[13])

We remark that the field $B_0(q^2)$ decreases monotonically from its value at $q^2 = 0$, $B_0(0)$. We can take $B_0(q^2) = 0$ for values of $q^2 > C^2\Lambda^2$, where $C \sim 4$ and Λ is a cutoff introduced in this model.[13] Thus we see that the choice of Λ will set the scale for the internal wave function of the condensed field or, alternatively, Λ sets the scale for the momentum variation of the self-energy $\Sigma(p^2)$. [It is important to note that if one assumes that λ_c is an ultraviolet fixed point, the value of $B_0(0)$ is independent of Λ. For example, $B_0(0)$ is proportional to κ, if $\lambda/\lambda_c = 1 + \pi \ln^{-2}(\Lambda/\kappa)$ for $\Lambda \to \infty$ (Refs. 13 and 17).]

III. NONTOPOLOGICAL SOLITONS

Let us now consider the situation where $C\Lambda$ is large compared to 1 MeV, since 1 MeV is the characteristic energy scale associated with the nontopological soliton we will consider ($R \sim 300$ fm). In that case we can take $B_0(q^2)$ to be a constant when discussing soliton structure. Therefore we may set

$$\sigma_q(P) \simeq \sigma_{\text{vac}}(P) = (2\pi)^4 \delta^{(4)}(P) B_0 \tag{3.1}$$

in the vacuum. For the purpose of studying soliton structure we now introduce

$$\chi(P) = \sigma(P) - (2\pi)^4 \delta^{(4)}(P) B_0 \tag{3.2}$$

as the field measuring the deviation from the vacuum value. In coordinate space this equation would read

$$\chi(x) = \sigma(x) - B_0 \ ,$$

which is meaningful if one does not probe distances of the order of Λ^{-1}.

At this point we need to bring back reference to the fermion field. We consider the equations

$$(i\gamma^\mu \partial_\mu - \tilde{m}_e)\psi_e(x) = g\psi_e(x)\chi(x) \ , \tag{3.3}$$

$$(\Box + m_\chi{}^2)\chi(x) = -g\bar{\psi}_e(x)\psi_e(x) \ , \tag{3.4}$$

which are the field equations obtained from the Lagrangian

$$\mathcal{L}(x) = \bar{\psi}_e(x)[i\gamma^\mu \partial_\mu - \tilde{m}_e - g\chi(x)]\psi_e(x) + \tfrac{1}{2}\partial^\mu \chi(x) \partial_\mu \chi(x) - m_\chi{}^2 \chi^2(x)/2 \ . \tag{3.5}$$

Here we have taken the effective potential to be

$$V(\chi) = \tfrac{1}{2} m_\chi{}^2 [\sigma(x) - B_0]^2 \tag{3.6}$$

$$= \tfrac{1}{2} m_\chi{}^2 \chi^2(x) \ , \tag{3.7}$$

with the zero of the potential at $\sigma(x) = B_0$. We can set $\tilde{m}_e = gB_0$, or $\tilde{m}_e = m_e + gB_0$. In the latter case we assume that the electron has a mass m_e when $B_0 = 0$. We also see that m_χ is the mass of the σ field. Note that m_χ is shown to be nonzero for $\lambda > \lambda_c$ in the model of Ref. 13.

Although strongly coupled QED is expected to be in a confined phase,[11,16,18] we see that Eqs. (3.3) and (3.4) make no reference to confinement. This situation can be remedied and one can discuss confined solutions for nontopological solitons; however, as can be seen from our previous work,[12,19,20] modeling confinement is not particularly important if one only wishes to describe low-lying states of hadrons. For example, we can describe the 1*S*, 2*S*, and 3*S* states of *b* quarkonium without having to introduce a model for confinement.[19] This matter will be discussed in a future work, and is not particularly relevant to our considerations here.

IV. HEAVY-ION COLLISIONS

The physical picture which we put forward is as follows. In some manner, which we do not understand, the intense fields of the heavy ions take one into a new phase of strongly coupled QED (Ref. 21). There is a vacuum condensate of bound e^+e^- pairs formed over some region of space, several hundred fm in size. The presence of a quasielectron and a quasipositron destroys the condensate over a finite region of radius $r_0 < R$, and "pseudopositronium" is formed. This object is metastable with a lifetime τ greater than 10^{-20} sec. The width $\Gamma \sim 1/\tau$ may be of the order of 20 keV and is due to the collapse of the soliton field [$\chi(x) \to 0$, or $\sigma(x) \to B_0$]. Note that this width is very much larger than the width which would be obtained by scaling the leptonic widths for charmonium or *b*-quarkonium decays. In the latter case (QCD) the vacuum condensates are *stable* and exist through all space exterior to the hadron.

One may ask how one can have a relatively small-mass object (~ 1.8 MeV) in QED and not destroy the excellent agreement achieved in QED for a broad range of observables. In our model, however, the results of standard QED calculations are unmodified. We may contrast our soliton model of bound states in QED with bound states in QCD. Consider the leptonic decay of charmonium ($J/\psi \to e^+ + e^-$), for example. One may then perform an electron-positron scattering experiment. At the appropriate energy, one readily forms charmonium. We suggest, however, that an electron-positron collision with a center-of-mass energy of about 1.8 MeV will not form a new vacuum phase and no anomalous peaks will be seen in such an experiment. In order to form a nontopological soliton in a new phase one would have to perform an experiment in which the momenta of all the final-state decay products, including the heavy ions, were reversed. In other words, one cannot create the metastable object, which we suggest is responsible for the e^+-e^- peaks, from the ordinary vacuum phase without the generation of the intense electric fields created by the heavy ions. Therefore, ordinary QED-based dynamics does not lead to any constraints on our model. This is in marked contrast to the various models which require the introduction of new (scalar or pseudoscalar) particles coupled to the electromagnetic field in the normal QED phase.[7,8]

V. DISCUSSION

It should be noted that chiral-symmetry breaking has been demonstrated in lattice electrodynamics (see Bartholomew *et al.*[18]). However, the critical coupling constant $\lambda_c^{\rm lat}$ is about a factor of 3 *smaller* than the value found by analytical methods ($\lambda_c^{\rm anal}=\pi^2$). [This could reflect the fact that the "ladder approximation" of Eq. (2.13) is inadequate. That is to say, the lattice calculation sums many diagrams which are neglected in the ladder approximation.] Another interesting point is that the lattice result indicates that there is no new scale developed dynamically in strongly coupled QED which would set a size for the condensed fermion-antifermion pair wave function. The scale appears to be set by the lattice spacing. This is in accord with the observation that, in the analytical analysis, the scale for $B(q^2)$ is set by the cutoff Λ if one does not assume the presence of an ultraviolet fixed point in strongly coupled QED (Ref. 17).

Of course, the real challenge for our model is to understand why a gauge theory such as QED should become strongly coupled in the unusual environment created in a heavy-ion collision. Researchers who have studied the problem of explaining the narrow peaks seen in heavy-ion experiments appear to now believe that one is seeing the decay of a composite object.[7] To the extent that the description of chiral-symmetry breaking in QED presented in Ref. 13 is correct, we can construct hadronlike states in QED making use of the order parameters of the condensate. (The situation in QED is, of course, different than that of QCD, where the condensates fill all of space outside a hadron. This difference would have to be taken into account in any detailed model.)

Further work is needed to see if the decay of such objects is indeed the source of the narrow electron-positron peaks observed. We do predict, however, that while exotic states will continue to be seen in heavy-ion collisions, these states will not be excited in electron-positron collisions.

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