Quantum electrodynamics based on self-fields: On the origin of thermal radiation detected by an accelerating observer

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We continue with our series of papers concerning a self-field approach to quantum electrodynamics that is not second quantized. We use the theory here to show that a detector with a uniform acceleration a will respond to its own self-field as if immersed in a thermal photon bath at temperature $T_a = \hbar a/2\pi kc$. This is the celebrated Unruh effect, and it is closely related to the emission of Hawking radiation from the event horizon of a black hole. Our approach is novel in that the radiation field is classical and not quantized; the vacuum field being identically zero with no zero-point energy. From our point of view, all radiative effects are accounted for when the self-field of the detector, and not the hypothetical zero-point field of the vacuum, acts back on the detector in a quantum-electrodynamic analog of the classical phenomenon of radiation reaction. When the detector is accelerating, its transformed self-field induces a different back reaction than when it is moving inertially. This process gives rise to the appearance of a photon bath, but the photons are not real in the sense that the space surrounding the accelerating detector is truly empty of radiation, a fact that is verified by the null response of an inertially moving detector in the same vicinity. The thermal photons are in this sense fictitious, and they have no independent existence outside the detector.

I. INTRODUCTION

In the wake of the discovery by Hawking of the apparent thermal emission from the event horizon of a black hole,[1] there came a related calculation by Unruh[2] that indicated that a uniformly accelerating particle detector would perceive a thermal bath of photons. If an idealized point detector is accelerating at a rate a, then the photon spectral distribution is Planckian at a temperature $T_a = \hbar a/2\pi ck$, where k is the Boltzmann constant. This thermal radiation is not picked up by an inertially moving detector, and the vacuum expectation of the normal ordered stress energy tensor $T_{\mu\nu}$ is identically zero in both the inertial and accelerated or unprimed and primed frames, respectively;[3] $\langle 0|{:}T_{\mu\nu}{:}|0\rangle = \langle 0|{:}T'_{\mu\nu}{:}|0\rangle = 0$. In what sense then can one say that these thermal photons are physically real if they do not alter the above expectation values? Davies argues that these results are indicative of a breakdown of the traditional quantum field theoretical notion of a particle when space-time is curved.[3] The present authors contend that the problem is not with the concept of particle but rather with the quantum field treatment of the vacuum field. Boyer has given an account of the Unruh effect in the framework of stochastic electrodynamics, which lends credence to the viewpoint that the acceleration somehow turns the virtual quanta of the Minkowski vacuum into real quanta.[4] In stochastic electrodynamics the zero-point field is taken to be a very real thing, responsible for many quantum-electrodynamical phenomena. The idea is that a classical vacuum with a spectrum proportional to $\frac{1}{2}\hbar\omega$ per normal mode is permissible on the grounds of Lorentz invariance. If one chooses the proportionality constant appropriately, one recovers a classical vacuum field that is nearly identical to that predicted by the second quantization procedure in field theory. Boyer then shows that under acceleration, the zero-point term is deformed into a zero-point plus Planckian spectrum at the Unruh temperature $T_a = \hbar a/2\pi kc$. The transformation is

$$\hbar\omega \to \hbar\omega \left[\frac{1}{2} + \frac{1}{e^{\hbar\omega/kT_a} - 1} \right], \tag{1}$$

where we will from now on set $\hbar = c = a = 1$.

But are these thermal photons really real? Indeed, one may ask if even the virtual Minkowski photons with the spectrum of $\frac{1}{2}\hbar\omega$ have any real existence apart from the detector that appears to register them, say, as the apparent "trigger" for spontaneous emission. In stochastic electrodynamics the choice of a nonzero proportionality constant for the spectrum proportional to $\frac{1}{2}\hbar\omega$ is permissible, but not required. The other obvious choice is to set the spectrum of the vacuum identically equal to zero as is done in classical electrodynamics. Where then would radiative effects such as spontaneous emission and the Lamb shift originate if not driven by the vacuum fluctuations, as is usually assumed in quantum electrodynamics (QED)? In classical electrodynamics there are two perfectly respectable phenomena which should correspond to the classical limit of spontaneous emission and the Lamb shift in atoms; they are line breadth and level shift in the energy, for instance, of a harmonically bound charge.[5] These radiative corrections to the otherwise unperturbed motion of the charge arise not from any interaction with a zero-point field—the classical vacuum field is identically zero—but rather from the radiation reaction on the charge from its own self-field. The scale of the electromagnetic field fluctuations is set by the constant $\hbar c$, so in the classical limit of $\hbar \to 0$, one would expect spontaneous emission and the Lamb shift to vanish and to have no classical analog since the causative agent, the zero-point field, has vanished. This is clearly not the case in that we are actually left with the classical line breadth and level shift of an oscillating charge. Barut and his co-workers have shown that it is possible to formulate QED in terms of self-fields, so that such phenomena as spontaneous emission and the Lamb shift are viewed as natural generalizations of their classical counterparts in radiation reaction theory.[6] This is the approach that we shall use in the present work.

In QED it is usual to renormalize the free electromagentic field through the normal ordering of the operators in order that the zero-point energy of $\frac{1}{2}\hbar\omega$ per normal mode vanishes. This is done primarily because the keeping of the $\frac{1}{2}\hbar\omega$ in the Hamiltonian would lead to an infinite energy density of empty space since $\langle 0|T_{\mu\nu}|0\rangle$ would diverge. The rationalization usually given for this procedure is that only energy differences have physical meaning, and hence a transfinite translation of the energy scale cannot have physical consequences. But the energy density $T_{\mu\nu}$ does indeed have an absolute meaning when coupled to the gravitational field, in the sense that it determines the curvature of the metric via the Einstein field equations. It is not possible to change the curvature or to flatten out space-time simply by adjusting the energy scale. If we accept the electromagnetic zero-point energy as real, then by implication we must accept the infinite energy density of empty space. This implies an infinite curvature for the universe and infinite value for the cosmological constant—unless we are saved in some unforeseen fashion by a fortuitous cancellation of all the vacuum fields in some unified field theory. The cosmological constant Λ is the most accurately determined physical quantity in all of physics; the observations by Sandage[7] of distant galaxies puts $|\Lambda|=0$ with an upper limit of $|\Lambda|<10^{-56}\,\mathrm{cm}^{-1}$. It certainly is not infinite.

It is common to say that the vacuum fluctuations are the physical cause of spontaneous emission, the Lamb shift, the nonzero value of $g-2$, the Casimir effect, long-range Casimir-Polder van der Waals forces, apparatus dependent contributions to these radiative effects, and now the thermal response of an accelerating detector. This view is perhaps that of the majority. It is not as well appreciated that all of these effects may be equally well explained at least to order α, in terms of the fields which originate in the charged particles themselves.[6,8]

Jaynes[9] has given a nice example of why the zero-point fluctuation interpretation of radiative effects in QED makes many of us uneasy. Suppose we believe that the electromagnetic zero-point energy is physically real, right up to the Compton cutoff frequency $\omega=m$, which is used in nonrelativistic calculations of the Lamb shift to get the correct Bethe logarithm. If one computes the turbulent energy flow associated with the zero-point field at this cutoff, one gets a Poynting vector of about 6×10^{20} $\mathrm{MW\,cm^{-2}}$. (The total luminosity of the sun is about 2×10^{20} MW.) One feels that physically real radiation of this intensity would have slightly more of an effect than to shift the $2s$ level of hydrogen by 4 μV.

Much work has been done in the past few years to show that there is a deep and fundamental connection between the vacuum fluctuation and the self-field approaches to QED.[10] The duality between these two methods of doing QED does not necessarily prove, however, that the zero-point field is real. It is possible that self-field effects are the same *as if* vacuum fluctuations were the causative agent. Jaynes has shown that the energy density of the radiation reaction field over the spectral interval of the natural linewidth is exactly the same as that of the vacuum field.[9] In the present paper we will support this idea that the vacuum field approach is a mathematical subterfuge which gives the correct answer some of the time.

Davies has emphasized that the meaning of the concept of a particle and the codependent concept of the vacuum depends crucially on the state of motion and history of the particle detector.[3] This is a fact which is often overlooked in Minkowski space, but which cannot be ignored in curved space-time where the decomposition of the field into positive and negative frequency normal modes is not unique for all observers. In general, different detectors will disagree on what constitutes the vacuum. If one detector registers no particles, a different detector on a different world line, in general, *will* register particles. This is because a Bogoliubov transformation between the two Fock spaces used to define the vacuum for each detector will not give identical vacuums for the two spaces. Davies concludes that because of this the concept of a particle, say a photon in the electromagnetic case, is not well defined. The present authors would like to use the same evidence to support a different conclusion: It is the standard notion of the vacuum in quantum field theory that is not well defined, a fact which seems obvious when one begins to consider quantum fields in curved space.

The stochastic electrodynamics theory of Boyer also develops pathological problems in curved space-time. Boyer chooses a classical zero-point spectrum proportional to $\frac{1}{2}\hbar\omega$ per normal mode because this is the only nonzero spectrum permitted on the grounds of Lorentz invariance. This means that in Minkowski space the stochastic electrodynamic vacuum is permitted since it is an invariant of the Poincaré group. The Poincaré group is not a symmetry group of a general curved space-time, however, and apparently the most compelling reason for choosing a stochastic zero-point spectrum proportional to $\frac{1}{2}\hbar\omega$, rather than to zero, completely disappears. The only choice of such a proportionality constant consistent with the demands of a space-time of arbitrary curvature is one that is zero. By Boyer's own reasoning we must conclude the only allowable classical vacuum field in curve space-time is the same as that used in classical electrodynamics—namely, the vacuum field must be chosen to be identically equal to zero in all its moments of the Wightman correlation function.

In discussing quantum fields in curved space-time the concept of a detector plays a central role: It is impossible to discuss properties of the quantum vacuum field in the absence of a detector to observe those properties. One cannot speak of the absence or presence of a vacuum without a detector to register deviations or nondeviations from that vacuum state. The concept of a vacuum in the absence of a detector is meaningless in both a philosophical and operational sense. But by the definition of a detector, it must couple to the vacuum field whose presence or absence we are trying to measure, and hence introduce its own self-field into the measurement process.

In the present paper we shall show that the self-field of a uniformly accelerating point detector responds to the acceleration in such a way as to drive the detector, via a quantum generalization of radiation reaction, into a superposition of states which when thermodynamically analyzed yields the Planck spectrum given in Eq. (1). But now the interpretation is different. The $\frac{1}{2}\hbar\omega$ corresponds to the spectral distribution of the detector's own field over the natural linewidth. For an inertially moving detector this is the only term which occurs, and it is responsible for the usual free-space atomic spontaneous emission as well as for the Lamb shift. Transforming to the Rindler coordinates of a uniformly accelerating detector we obtain the full result of Eq. (1). The interpretation now is *not* that the detector is immersed in a bath of thermal photons, but rather that the self-field of the detector is responding to the work being done on it by the accelerating agent in such a fashion so as to make it appear *as if* the detector were immersed in such a bath. There are no physical photons present, a fact which would be confirmed by a neighboring inertially moving detector. The particle concept is hence rescued, but only with the sacrifice of the notion of a dynamic vacuum state in the absence of a detector. Since one cannot discuss the vacuum without the detector, it would seem compelling to want to set the vacuum field identically equal to zero for all observers and then to attribute differences between detectors totally to the response of the self-field of the detector to its own worldline.

Notice that by the Einstein equivalence principle, a uniformly accelerating detector is equivalent to a detector at rest in a gravitational field. From the point of view of general relativity the thermal radiation seen by the detector seems to originate from a neighboring event horizon; in our calculation it would be the event horizon of Rindler coordinates. But this event horizon is related to that of a black hole by a conformal transformation. Hence Unruh and Hawking radiations are similar, and from the self-field point of view they are both equally interpreted in the sense that the thermal radiation effects are confined to within the detector, which is responding to the gravitational field directly.

II. SELF-FIELD APPROACH TO QED

In classical electrodynamics one usually computes the zeroth-order motion of the charges first and then adds on the self-field or radiation reaction effects as a perturbation to the original motion. (Although, in special circumstances, one may find exact solutions to the nonlinear Lorentz-Dirac equation of motion.) The philosophy in the self-field approach to QED is precisely the same. Conceptually we may separate the electromagnetic four-potential A_μ surrounding a point charge into and external field A^e_μ, which is prescribed as part of the initial conditions, and a self-field term A^s_μ, which originates from the charge. The total field is then $A_\mu := A^e_\mu + A^s_\mu$. (The notation $a := b$ indicates that a is being defined as being equal to b with the colon on the same side as the quantity to be defined.) With this separation, we shall assume that the coupling of the Dirac spinor field Ψ to A^e_μ alone determines the bulk, zeroth-order motion of the electron, while the coupling to A^s_μ is responsible for all radiative corrections. As in previous work we find it convenient to proceed from the standpoint of an action formalism,[6,8] with the action given by

$$W = \int dx\, w[x;\Psi;A] , \tag{2}$$

where W is the total action and $w(x)$ is the corresponding action density. Variation of Eq. (2) with respect to Ψ yields the Dirac equation of motion, while variation with respect to A_μ gives Maxwell's equations. The electromagnetic field tensor $F_{\mu\nu}$ is defined as usual as

$$\begin{aligned} F_{\mu\nu} &:= A_{[\nu,\mu]} \\ &= A^e_{[\nu,\mu]} + A^s_{[\nu,\mu]} \\ &=: F^e_{\mu\nu} + F^s_{\mu\nu} , \end{aligned} \tag{3}$$

where $[\,,]$ indicates commutation with respect to the indices. Far from the source of the external field we have Maxwell's equations as

$$(F^{\mu\nu})_{,\mu} = (F^{\mu\nu}_s)_{,\mu} = ej^\nu , \tag{4}$$

where j_μ is the usual Dirac current $\Psi\gamma_\mu\Psi$. The action density $w(x)$ can be written now as

$$\begin{aligned} w &= \overline{\Psi}(i\gamma^\mu\partial_\mu - m)\Psi + eA_\mu j^\mu + \tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} \\ &=: w_0 + w_1 + w_2 , \end{aligned} \tag{5}$$

where w_0 is the free particle density, w_1 the particle-field coupling, and w_2 the free electromagnetic (EM) field action density. It is evident that w_0 and w_1 taken together are equivalent to the canonical coupling $i\partial_\mu \to i\partial_\mu - eA_\mu$. At this point the external and self-electromagnetic field have not yet been separated. We proceed now with an analysis of w_2,

$$w_2 = \tfrac{1}{4}(F^e_{\mu\nu}F^{\mu\nu}_e + F^e_{\mu\nu}F^{\mu\nu}_s + F^s_{\mu\nu}F^{\mu\nu}_e + F^s_{\mu\nu}F^{\mu\nu}_s) . \tag{6}$$

The two middle terms can be converted into surface integrals under $\int dx$, which vanish if A^s_μ is sufficiently localized. The first term of this expression is the invariant

$$\tfrac{1}{4}F^e_{\mu\nu}F^{\mu\nu}_e = -\tfrac{1}{2}(E_e^2 - B_e^2) , \tag{7}$$

which is an additive constant that does not effect the equations of motion, and so we may drop it from the action. We are left with the last term, which can be transformed as

$$\begin{aligned} \tfrac{1}{4}F^s_{\mu\nu}F^{\mu\nu}_s &= \tfrac{1}{4}A^s_{[\nu,\mu]}F^{\mu\nu}_s \\ &= \tfrac{1}{4}(A^s_{[\nu}F^{\mu\nu}_s)_{,\mu]} - \tfrac{1}{4}A^s_{[\nu}(F^{\mu\nu}_s)_{,\mu]} \\ &= -\frac{e}{2}A^s_\nu j^\nu , \end{aligned} \tag{8}$$

where we have used the inhomogeneous Maxwell equations (2). Equality here is with respect to integration by parts and the vanishing of possible surface integrals under the application of $\int dx$. (Surface terms are, however, needed in the discussion of processes in which radiation goes to infinity such as in bremsstrahlung, the Compton effect, etc.) With these results the total density of expression (5) becomes

$$w=\bar{\Psi}(i\gamma^{\mu}\partial_{\mu}-m)\Psi+eA^{e}_{\mu}j^{\mu}+\frac{e}{2}A^{s}_{\mu}j^{\mu}$$
$$=:w_0+w_i+w_s \ . \tag{9}$$

Together w_0+w_i are responsible for the zeroth-order motion of the electron in the external field, while w_s induces radiative corrections to that motion.

One may formally solve Eq. (4) for A^{s}_{μ} in terms of the current j_{μ} through the use of an electromagentic Green's function $D_{\mu\nu}$ via

$$A^{s}_{\mu}(x)=e\int dy\, D_{\mu\nu}(x-y)j^{\nu}(y) \ . \tag{10}$$

If we define $W_s:=\int dx w_s(x)$ as the contribution to the total action W from the self-field correction, then inspection of expressions (9) and (10) yields immediately that

$$W_s=\frac{e^2}{2}\int\int dx\, dy\, j^{\mu}(x)D_{\mu\nu}(x-y)j^{\nu}(y) \ . \tag{11}$$

This single nonlinear addition to the usual action contains information about all radiative effects, e.g., spontaneous emission, the Lamb shift, and the electron $g-2$ value.[6,8] The interpretation here is, once again, that these radiative corrections arise as an effect of the back reaction of the self-field on the motion of the electron in a manner analogous to the classical phenomenon of radiation reaction.

To insure the boundary conditions that provide the correct combination of retarded and advanced potentials, we choose for the Green's function $D_{\mu\nu}(x-y)$ the causal Feynmann propagator $D_{\mu\nu}(x-y):=-\eta_{\mu\nu}D_{\mu\nu}(x-y)$, where $\eta_{\mu\nu}$ is the Minkowski metric tensor with signature $(+---)$ and $D_{\mu\nu}(x-y)$ has the equivalent forms

$$D(x-y)=\frac{i}{4\pi^2}\left[\frac{1}{(x-y)^2}-i\pi\delta((x-y)^2)\right] \tag{12a}$$
$$=\frac{i}{4\pi^2}\left[\frac{1}{(x-y)^2+i\epsilon}\right] \tag{12b}$$
$$=\frac{1}{(2\pi)^4}\int dk\frac{e^{-ik(x-y)}}{k^2+i\epsilon} \tag{12c}$$

in the Feynman gauge.

Now in order to further analyze W_s in Eq. (11) we perform a Fourier expansion of the Ψ in terms of quasibound-state energies E_n, which are to be determined, since we anticipate using a bound electron as our Unruh detector. The expansion is

$$\Psi(x)=\sum\!\!\!\!\!\!\int_n \Psi_n(\mathbf{x})e^{-iE_n t} \ , \tag{13}$$

where the sum runs over positive and possibly negative energy levels. To a first iteration we assume that the Ψ_n with associated eigenvalues E_n exactly minimize the action W_0+W_i. We are then using these zeroth-order wave functions to evaluate the W_s radiative correction term in an iterative fashion. Inserting expression (12c) for the Green's function into the expression (11) for W_s; expanding each of the Ψ as per Eq. (13); and carrying out the dx_0, dy_0, and dk_0 integrations yields

$$W_s=-\frac{e^2}{2}(2\pi)^4\sum_{n,m,p,q}\int\int\int dx^3dy^3dk^3[\bar{\Psi}_n(\mathbf{x})\gamma_{\mu}\Psi_m(\mathbf{x})]$$
$$\times[\bar{\Psi}_p(\mathbf{y})\gamma^{\mu}\Psi_q(\mathbf{y})]$$
$$\times\frac{e^{ik\cdot(x-y)}}{\omega_{nm}^2-|\mathbf{k}|^2+i\epsilon}$$
$$\times\delta(\omega_{nm}+\omega_{pq}) \ , \tag{14}$$

where $\omega_{nm}:=E_n-E_m$. The δ function can be satisfied by either of the two choices

$$n=m, \quad p=q \tag{15a}$$
$$n=q, \quad m=p \ . \tag{15b}$$

The condition (15a) leads to a vacuum polarization term,[6] and we shall not consider it here. The condition (15b) leads to spontaneous emission and the Lamb shift, here interpreted as quantum analogs of the classical radiation reaction effects of line broadening and level shift.[6] We will consider in the present work how these phenomena are effected by boosting the detector into an accelerating frame.

III. RESPONSE OF AN INERTIALLY MOVING DETECTOR TO ITS SELF-FIELD

To illustrate the self-field method of approach we will now confirm that a pointlike detector on an inertial worldline experiences at zero temperature only the effects of the usual spontaneous emission and Lamb shift in free space, which occur via the interaction of the detector with its own field. (This is, of course, clear form the Lorentz invariance of the theory, but it is instructive as an illustrative example of how to apply the self-field approach to this kind of problem.)

The trajectory for an inertial detector can be written as

$$\mathbf{x}=\mathbf{x}_0+\mathbf{v}t=\mathbf{x}_0+\mathbf{v}\gamma\tau \ ,$$
$$\mathbf{y}=\mathbf{y}_0+\mathbf{v}u=\mathbf{y}_0+\mathbf{v}\gamma\nu \ , \tag{16}$$

where τ and ν are the proper times which correspond to the x_{μ} and the y_{μ} time components $t=\gamma\tau$ and $u=\gamma\nu$, respectively. The velocity $\mathbf{v}$ is constant, with $v:=\beta<1$ and $\gamma:=(1-\beta^2)^{-1/2}$ as usual. The coordinates $\mathbf{x}_0$ and $\mathbf{y}_0$ are those of the electron in a detector based system. For a pointlike detector we may take $\mathbf{x}_0-\mathbf{y}_0\approx\mathbf{0}$ or $\exp[\mathbf{k}\cdot(\mathbf{x}_0-\mathbf{y}_0)]\approx 1$ in the dipole approximation.

Inserting the expressions (16) into the Green's function of (12b), we first notice that

$$(x-y)^2+i\epsilon=(\xi+i\epsilon)^2 \ , \tag{17}$$

where we have defined $\xi:=\tau-\nu$, and we have absorbed a positive function of ξ into the ϵ. The self-field action of Eq. (11) can now be written as

$$W_s^{\text{inertial}} = -\frac{e^2}{8\pi^2}\sum_{n,m}\int\int d\tau\, d\nu \frac{\langle n|\gamma_\mu|m\rangle\langle m|\gamma^\mu|n\rangle}{(\xi+i\epsilon)^2} \times e^{i\omega_{nm}\xi}, \quad (18)$$

where we have adopted the Dirac bra ket notation, and used the relation (15b) after expanding the Ψ as per the prescription of Eq. (13). The action W is formally infinite, but it can be related to the transition probability G_n for the nth energy level via[6,8]

$$W = 2\pi\delta(\omega_{nm})G_n ,$$

and the identification $\int\int d\tau\, d\nu \to \int d\xi$ gives us the finite transition probability per unit time for this nth state as

$$G_n = \frac{\alpha}{2\pi}\sum_m \langle n|\gamma_\mu|m\rangle\langle m|\gamma^\mu|n\rangle(\tfrac{1}{2}\omega_{nm}) . \quad (19)$$

In our units $\alpha = \theta^2/4\pi$, and the contour integral was carried out on an infinite semicircle in the lower ξ plane. A single pole of order 2 located at $\xi = -i\epsilon$ contributes

$$\tfrac{1}{2}2\pi i \operatorname{res}\left[\frac{e^{i\omega_{nm}\xi}}{(\xi+i\epsilon)^2}\right] = -2\pi(\tfrac{1}{2}\omega_{nm}) ,$$

where we note that if n is the ground state then $\omega_{nm} < 0$. A detailed analysis of (19) shows that this corresponds to the usual spontaneous-emission transition rate and Lamb shift in free space.[6] Notice how the factor of $\frac{1}{2}\hbar\omega_{nm}$ enters here, not as a consequence of any electromagnetic zero-point energy, but rather through the Fourier spectrum of the detector's self-field. Once again it looks *as if* there is a vacuum field which is stimulating the spontaneous emission or Lamb shift; in reality it is the detector's own field that is responsible.

IV. RESPONSE OF A UNIFORMLY ACCELERATING DETECTOR TO ITS SELF-FIELD

In the self-field approach to QED spontaneous emission occurs as a back reaction of the field on the detector. In curved space-time, such as in the Rindler coordinates of a uniformly accelerating observer, one would expect the self-field to become modified by the curvature and by a non-Minkowskian event horizon. Any change in the configuration of the self-field would be transmitted through the radiation reaction effects and would surface as a modification of the spontaneous emission rate, as well as of other radiative effects. We now compute this, the Unruh effect, from the self-field point of view.

Let us suppose that our detector is accelerating uniformly with acceleration $a := 1$. The worldline is hyperbolic and can be written in Rindler coordinates as

$$\begin{aligned} &x_0 =: t = \sinh(\tau), \quad y_0 =: u = \sinh(\nu) , \\ &x_3 = \cosh(\tau), \quad y_3 = \cosh(\nu) , \\ &x_1 = x_2 = 0, \quad y_1 = y_2 = 0 , \end{aligned} \quad (20)$$

with τ and ν the proper times as before. Hence we have, in the dipole approximation,

$$\begin{aligned} (x-y)^2 + i\epsilon &= 4\sinh^2\left[\frac{\xi}{2} + i\epsilon\right] \\ &= -4\sin^2\left[\frac{i\xi}{2} - \epsilon\right] . \end{aligned} \quad (21)$$

Once again we have absorbed a positive valued function of ξ into the ϵ. If we now insert expression (21) into Eq. (12b) for the Green's function we get

$$D(x-y) = -\frac{1}{16\pi^2}\csc^2\left[\frac{i\xi}{2} - \epsilon\right] . \quad (22)$$

Making use of the Laurent expansion for cosecant, namely,

$$\csc^2(z) = \sum_{p=-\infty}^{\infty}\frac{1}{(z-\pi p)^2} , \quad (23)$$

we can expand the Green's function of (22) as

$$D(x-y) = \frac{1}{4\pi^2}\sum_{p=-\infty}^{\infty}\frac{1}{(\xi + 2\pi i p + i\epsilon)^2} . \quad (24)$$

If we use this expression for the Green's function in the self-field action of (11) we eventually arrive at

$$W_s^{\text{accelerating}} = -\frac{\alpha}{2\pi}\sum_{n,m}\langle n|\gamma_\mu|m\rangle\langle m|\gamma^\mu|n\rangle \times \sum_{p=-\infty}^{\infty}\int\int d\tau\, d\nu \frac{e^{i\omega_{nm}\xi}}{(\xi+2\pi i p + i\epsilon)^2} , \quad (25)$$

with $\xi := \tau - \nu$ as before. Converting, as in the inertially moving case, to the transition probability per unit time per energy level n, and carrying out the integral on the same contour as before, we get

$$G_n = \frac{\alpha}{2\pi}\sum_m \langle n|\gamma_\mu|m\rangle\langle m|\gamma^\mu|n\rangle \times \left[\frac{1}{2}\omega_{nm} + \frac{\omega_{nm}}{e^{2\pi\omega_{nm}} - 1}\right] , \quad (26)$$

where we have summed a convergent geometric series

$$\sum_{p=1}^{\infty} e^{2\pi p\omega_{nm}} = \frac{1}{e^{2\pi\omega_{nm}} - 1} ,$$

which arises as the sum of the residues contributed from the infinitude of poles enclosed by the contour and located at $\xi = -2\pi i p$ for $p = 1,2,3,\ldots$. (Recall that if n is the ground state, then $\omega_{nm} < 0$.) This is our primary result, reinserting the constants into the expression in parentheses on the last line of (26) yields for the contents of these parentheses

$$\frac{1}{2}\hbar\omega_{nm} + \frac{\hbar\omega_{nm}}{e^{2\pi c\omega_{mn}/a} - 1} , \quad (27)$$

which we see is the Planck blackbody spectral distribution, complete with the so-called zero-point term. However, as we saw in Sec. III, the $\frac{1}{2}\hbar\omega_{nm}$ does not correspond to a vacuum spectrum but rather to the self-field spectrum of an inertially moving detector. When we boost into an accelerating frame we get back the inertial term plus the Planckian term at a temperature of

$$T_a = \frac{\hbar a}{2\pi k c} . \tag{28}$$

So then it appears as if the accelerating detector is exposed to a thermal bath of photons at temperature T_a; just as it appears as if it is also being exposed to a zero-point field embodied in the $\frac{1}{2}\hbar\omega$. Neither set of photons are physically real. Since any nearby inertially moving detector would detect no photons, the accelerating detector cannot be detecting real photons either. By the equivalence principle we can conclude the same thing about a detector placed in a uniform gravitational field of strength a. The field does not create a bath of photons, rather the detector is responding directly to the local curvature of space-time. The energy required to excite the detector into a thermal superposition of states is tapped directly from the metric without the intermediary of any electromagnetic radiation.

If our results generalize to the case of black holes, then we would conclude that although a black hole has the capability of directly exciting detectors in its neighborhood, it does not necessarily do so by emitting a flood of thermal radiation. If this is indeed the case, then black holes do not radiate in the ordinary sense of that word, i.e. they do not lose mass or energy via this mechanism unless a detector is actually present.

V. CONCLUSION

We have shown that the Unruh effect can be calculated within the context of a source-field theory; we conclude that the thermal response of the detector arises not through an interaction with real photons in the surrounding space, but from the spectrum of its self-field which has become altered by the change to a noninertial frame. This indicates that the detector is becoming excited directly by changes in the metric tensor. If all such responses of the detector can be attributed to modifications of the self-field, as opposed to modifications in a vacuum field, it would seem unlikely then that black holes are emitting real, physical radiation.

Davies has argued that, in particular, the concept of the photon is not well defined in curved space-time quantum electrodynamics, since in a curved space-time different detectors respond differently to the vacuum field and "detect" different photon spectra.[3] It is our contention that it is the vacuum field in QED that is not well defined; if it were identically zero to begin with it would not cause trouble in any space-time, curved or otherwise. Davies has persistently pointed out that any discussion of the vacuum field must always concern itself with the worldline of the detector which registers departures from that vacuum. Herein lies the key. Since in curved space-time the concept of vacuum and detector cannot be either conceptually or operationally separated, this is a clue that they are really two sides of the same coin. The fact that the source field and the vacuum field are closely related is well documented;[10] but in the framework of standard QED it appears as if both are always required for the internal consistency of the theory. The fact that the vacuum field is required in standard QED is a direct consequence of the second quantization procedure. In the present approach there is no quantization of the EM field, and yet we obtain correct results at least to order α, for radiative effects thought to require second quantization for their explanation.[6,8] So the question is as follows: Can we always get the correct results without recourse to some sort of vacuum fluctuation? If we can set the zero-point field identically equal to zero for all moments in the Wightman, two-point correlation function, then the concept of photon might be rescued. There are no photons in the Minkowski vacuum or any other vacuum to be counted by any detector—regardless of its state of motion or history.

The self-field approach to quantum electrodynamics as presented in this work has been used successfully to first order in α to account for nonrelativistic and relativistic formulas for spontaneous emission, the Lamb shift, and $g-2$, all in free space.[6,8] This approach has also been used to compute various apparatus dependent contributions to these effects; again to lowest order. By enlarging the notion of boundary to include the effects of a non-Minkowskian event horizon we have, in the present work, accounted for the Unruh effect. Work is in progress to show that the Unruh effect is essentially a classical phenomenon, and to analyze the response of an atom to Planck blackbody radiation in a more general setting from the point of view of self-fields. (Since spontaneous emission and the Lamb shift have the classical limits of line broadening and level shift, the Unruh effect should also have such a classical analog.) The extension of the self-field approach to higher orders of the fine-structure constant has until now been hampered by delicate and complicate calculations of the wave functions needed in the general nth order iteration of the self-field contribution to the total action. Some progress is now being made in this direction.

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