Is there emitted radiation in Unruh effect?

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Abstract

The thermal radiance felt by a uniformly accelerated detector/oscillator/atom— the Unruh effect— is often mistaken to be some emitted radiation detectable by an observer/probe/sensor. Here we show by an explicit calculation of the energy momentum tensor of a quantum scalar field that, at least in 1+1 dimension, while a polarization cloud is found to exist around the particle trajectory, there is no emitted radiation from a uniformly accelerated oscillator in equilibrium conditions. Under nonequilibrium conditions which can prevail for non-uniformly accelerated trajectories or before the atom or oscillator reaches equilibrium, there is conceivably radiation emitted, but that is not what Unruh effect entails.

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I. INTRODUCTION

The title question has realistic significance in light of recent experimental proposals on the detection of 'Unruh radiation' emitted by linear uniformly accelerated charges. Earlier findings of ours and others had already addressed this issue, with the results that, at least in 1+1 dimension model calculations, there is no emitted radiation from a linear uniformly accelerated oscillator in a steady state, even though there exists a polarization cloud around it. There could be radiation emitted in nonequilibrium conditions, which arise for non-uniformly accelerated atoms (for an example of finite time acceleration, see [1]), or during initial transient time for a uniformly accelerated atom, when its internal states have
not yet reached equilibrium through interaction with the field. For a review of earlier work on accelerated detectors, see e.g., [6]. For a discussion of nonequilibrium processes beyond the Unruh effect, see [7,8].

After Unruh and Wald’s [9] explication of what a Minkowski observer sees, Grove [10] questioned whether an accelerated detector actually emits radiated energy. Raine, Sciama and Grove [11] (RSG) analyzed what an inertial observer placed in the forward light cone of the accelerating detector would measure, and concluded that the detector does not radiate. Unruh [12], in an independent calculation, basically concurred with the findings of RSG but he also showed the existence of extra terms in the two-point function of the field which could contribute to the excitation of a detector placed in the forward light cone. Massar, Parantani and Brout [13] (MPB) pointed out that the missing terms in RSG constitute a polarization cloud around the accelerating detector. Further discussion were conducted by Hinterleitner [14], Audretsch, Müller and Holzmann [15], Massar and Parantani [13].

Both RSG and RHA treated particle–field interaction as a quantum dissipative system. RSG attributed the lack of radiation from the accelerated detector to the existence of a fluctuation-dissipation relation (FDR) governing its dynamics. Using the open system concept RHA constructed the influence functional and derived a set of coupled stochastic equations for a system of n-detectors in arbitrary (yet prescribed) states of motion in a quantum field. One subcase they studied related to Unruh effect was the influence of an accelerated detector on a probe (which is not allowed to causally influence the accelerated detector itself) via the quantum field. They found that most of the terms in the correlations of the stochastic force acting on the probe cancel each other, owing to the existence of a correlation-propagation relation, related to the fluctuation-dissipation relation for the accelerated detector. The remaining terms, which contribute to the excitation of the probe, are shown to represent correlations of the free field across the future horizon of the accelerating detector.

Here, we will use the simpler Heisenberg operator method to calculate the two point function and the energy momentum tensor of a massless quantum scalar field in a $1 + 1$ -dimensional Minkowski spacetime minimally coupled to an accelerated particle with internal oscillator coordinates. Our analysis (based on Chapter 2 of Alpan Raval’s thesis [3]) is more

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1There is a common misconception that a FDR can be used to explain the cancellation of radiation reaction by vacuum fluctuations, not realizing that the former is classical in nature while the latter is a quantum entity. The FDR in our work exists at the quantum stochastic level and relates the quantum dissipation in a particle’s trajectory or an atom’s internal degrees of freedom to the vacuum fluctuations in a field. It does not involve radiation reaction, which vanishes for a uniformly accelerated charge because of special conditions existing for the classical acceleration fields.

2Such a relation can be equivalently viewed as a construction of the free field two-point function for each point on either trajectory from the two-point function along the uniformly accelerated trajectory alone.
general than that of MPB in that the two-point function is calculated for the two points lying in arbitrary regions of Minkowski space, and not restricted to lie to the left of the accelerated oscillator trajectory. We show where the extra terms in the two point function are which were ignored in the RSG analysis. More relevant to answering the title question, we show that at least in two dimensions the energy momentum tensor vanishes everywhere except on the horizon. This means that beyond the initial transient, there is no net flux of radiation emitted from the uniformly accelerated oscillator.

II. CORRELATIONS AND STRESS ENERGY OF QUANTUM FIELD

A. Minimal coupling particle-field model

As in RHA, we consider the scalar electrodynamic or “minimal” coupling of oscillators to a scalar field in $1+1$ dimensions. This coupling provides a positive definite Hamiltonian, and is of interest because it resembles the actual coupling of charged particles to an electromagnetic field. We assume that the field and the detector are initially decoupled from each other, and that the field is initially in the Minkowski vacuum state.

The complete action of the minimally coupled field-particle system is

$$S = \frac{1}{2} \int d\tau \left\{ \left( \frac{dQ}{d\tau} \right)^2 - \Omega_0^2 Q^2 \right\} + e \int d\tau \frac{dQ}{d\tau} \phi(x(\tau), t(\tau))$$

$$+ \frac{1}{2} \int dx \int dt \left\{ \left( \frac{\partial \phi}{\partial t} \right)^2 - \left( \frac{\partial \phi}{\partial x} \right)^2 \right\}.$$ (2.1)

where $\Omega_0$ is the bare frequency of the oscillator and $e$ its coupling constant to the field.

Under uniform acceleration, the particle trajectory parametrized by the proper time $\tau$ is

$$x(\tau) = a^{-1} \cosh a\tau; \quad t(\tau) = a^{-1} \sinh a\tau.$$ (2.2)

Variation of the action leads to the following equations of motion:

$$\frac{d^2 Q}{d\tau^2} + \Omega_0^2 Q = -e \frac{\partial \phi}{\partial \tau}(x(\tau), t(\tau))$$ (2.3)

$$\frac{\partial^2 \phi}{\partial t^2} - \frac{\partial^2 \phi}{\partial x^2} = e \int d\tau \frac{dQ}{d\tau} \delta(x - x(\tau)) \delta(t - t(\tau)) + e \int d\tau \frac{dQ}{d\tau} G_{ret}(x, t; x(\tau), t(\tau)).$$ (2.4)

Because the action is a quadratic functional of the field and oscillator variables, these are also the Heisenberg operator equations of motion for the system. We shall thus view the above equations as operator equations from now on.

The field equations are solved by introducing the retarded Green function of a massless scalar field in $1+1$ dimensions:

$$\phi(x, t) = \phi_0(x, t) + e \int_{-\infty}^{\infty} d\tau \frac{dQ}{d\tau} G_{ret}(x, t; x(\tau), t(\tau))$$ (2.5)
where $\phi_0$ is a solution to the homogenous field equations corresponding to $Q = 0$. We will find it convenient to introduce the null coordinates $u = t - x$ and $v = t + x$. Correspondingly, we also find it convenient to define the regions $F, P, R$ and $L$ of Minkowski space as ($R$ is called the Rindler wedge)

\[ F : u > 0, v > 0 \quad P : u < 0, v < 0 \]
\[ R : u < 0, v > 0 \quad L : u > 0, v < 0. \quad (2.6) \]

In terms of the $(u,v)$ coordinates, the retarded Green function for a massless scalar field in $1 + 1$ dimensions takes the form:

\[ G_{ret}(x,t; x(\tau), t(\tau)) = \frac{1}{2} \theta(t - t(\tau) - x + x(\tau))\theta(t - t(\tau) + x - x(\tau)) \]
\[ = \frac{1}{2} \theta(u + a^{-1}e^{-a\tau})\theta(v - a^{-1}e^{a\tau}). \quad (2.7) \]

With this substitution, an integration by parts in Eq. (2.5) yields:

\[ \phi(x,t) = \phi_0(x,t) + e^\frac{2}{2} \left[ \theta(-u)\theta(-\lambda)Q(-a^{-1}ln(|au|)) \right. \]
\[ \left. + \theta(v)\theta(\lambda)Q(a^{-1}ln(|av|)) \right] \quad (2.8) \]

where we have also defined $\lambda = 1 + a^2uv$. The oscillator trajectory satisfies $\lambda = 0$. The quantities $-a^{-1}ln(|au|)$ and $a^{-1}ln(|av|)$ are just the retarded times of the point $(x,t)$, according to whether it lies to the right or the left of the accelerated trajectory, respectively. These two cases are distinguished by the appearance of the step functions with argument $\mp \lambda$. The step functions in $u$ and $v$ distinguish the cases when the point lies anywhere in the past light cone or anywhere in the forward light cone of the accelerated particle (these two conditions are simultaneously satisfied only in the Rindler wedge). With this in mind, we see that the first term linear in the coupling constant contributes only for points to the right of the oscillator trajectory, whereas the second term contributes only for points to the left of the oscillator trajectory and within the forward light cone of the oscillator. In particular, as expected, there is no correction to the field operator in the region $L \cup P$, which cannot be causally influenced by the accelerated trajectory.

Along the accelerated trajectory, the solution for $\phi$ reduces to

\[ \phi(x(\tau), t(\tau)) = \phi_0(x(\tau), t(\tau)) + e^\frac{2}{2}Q(\tau). \quad (2.9) \]

Putting this back to the equation of motion for $Q$, (2.3), we obtain:

\[ \frac{d^2Q}{dt^2} + \frac{e^2}{2} \frac{dQ}{d\tau} + Q^2Q = -e^\frac{d\phi_0}{dt}(x(\tau), t(\tau)). \quad (2.10) \]

The term linear in the proper velocity of the oscillator degree of freedom arises from the oscillator - field interaction and corresponds to dissipation of a quantum origin in the oscillator.
B. Equation of motion for the detector

The above equation of motion is easily solved. If the oscillator field interaction has always been switched on, the oscillator has reached a steady state at any finite time. We can then ignore transient terms in the solution for $Q$ and obtain:

$$Q(\tau) = -\frac{e}{\Omega} \int_{-\infty}^{\tau} d\tau' \sin \Omega(\tau - \tau') e^{-\gamma(\tau - \tau')} \frac{d\phi_0}{d\tau'}(x(\tau'), t(\tau'))$$  \hspace{1cm} (2.11)

where we have defined the dissipation constant $\gamma = \frac{\varepsilon^2}{\eta}$, and the frequency $\Omega = \sqrt{\Omega_0^2 - \gamma^2}$. We may also solve equation (2.10) in frequency space. Ignoring transients as before, we obtain

$$\tilde{Q}(\omega) \equiv \frac{1}{2\pi} \int_{-\infty}^{\infty} d\tau e^{-i\omega\tau} Q(\tau) = \chi_\omega \tilde{J}(\omega)$$  \hspace{1cm} (2.12)

where

$$\tilde{J}(\omega) = -\frac{e}{2\pi} \int_{-\infty}^{\infty} d\tau e^{-i\omega\tau} \phi_0(x(\tau), t(\tau))$$  \hspace{1cm} (2.13)

and $\chi_\omega$ is the impedance function of the oscillator, given by

$$\chi_\omega = i\omega(-\omega^2 + \Omega_0^2 + 2i\omega\gamma)^{-1}.$$  \hspace{1cm} (2.14)

It satisfies the identity

$$\chi_\omega + \chi_\omega^* = 4\gamma |\chi_\omega|^2$$  \hspace{1cm} (2.15)

which has the form of a fluctuation-dissipation relation. We now expand $\phi_0$ in Minkowski normal modes:

$$\phi_0(t, x) = \phi_0^{(+)} + \phi_0^{(-)} = \int_{-\infty}^{\infty} \frac{d^2k}{\sqrt{(2\pi)^22\omega_k}} a_k e^{i(kx - \omega_k t)} + h.c.$$  \hspace{1cm} (2.16)

where $h.c.$ denotes Hermitian conjugate and $\phi_0^{(-)}$ is the Hermitian conjugate of $\phi_0^{(+)}$. The operators $a_k$ annihilate the Minkowski vacuum. Based on this separation of the field into positive and negative frequency parts, we obtain the corresponding separation of the oscillator degree of freedom:

$$Q(\tau) = Q^{(+)}(\tau) + Q^{(-)}(\tau)$$  \hspace{1cm} (2.17)

where

$$Q^{(+)}(\tau) = -\frac{e}{\Omega} \int_{-\infty}^{\tau} d\tau' \sin \Omega(\tau - \tau') e^{-\gamma(\tau - \tau')} \frac{d\phi_0^{(+)}}{d\tau'}(x(\tau'), t(\tau'))$$  \hspace{1cm} (2.18)
and \( Q(-) \) is the Hermitian conjugate of \( Q(+) \).

On the accelerated trajectory, Eq. (2.16) gives

\[
\phi_0^+(x(\tau), t(\tau)) = \int_{-\infty}^{\infty} \frac{d^2k}{(2\pi)^2} \mathbf{a}_k \left[ e^{ikx} e^{-\alpha \tau} \theta(k) + e^{-ikx} e^{\alpha \tau} \theta(-k) \right].
\] (2.19)

Introducing the Fourier transforms of \( e^{ikx} e^{-\alpha \tau} \) and \( e^{-ikx} e^{\alpha \tau} \), we get

\[
\phi_0^+(x(\tau), t(\tau)) = \frac{1}{2\pi a} \int_{-\infty}^{\infty} \frac{d^2k}{(2\pi)^2} \mathbf{a}_k \int_{-\infty}^{\infty} d\omega e^{-i\omega \tau} e^{\frac{i\omega x}{a}} \times \left[ \Gamma(-\frac{i\omega}{a}) | \frac{k}{a} |^{\frac{i\omega}{a}} \theta(k) + \Gamma(\frac{i\omega}{a}) | -\frac{k}{a} |^{\frac{i\omega}{a}} \theta(-k) \right].
\] (2.20)

Differentiating with respect to \( \tau \) and substituting in the equation for \( Q^+(\tau) \), (2.18), we obtain, after carrying out the integration over \( \tau \), an expression of \( Q^+(\tau) \). Then we can substitute it back into the equation for the field operator (2.8) and get

\[
\phi_{int}(x, t) = -\frac{\gamma}{\pi a} \int_{-\infty}^{\infty} \frac{d^2k}{(2\pi)^2} \mathbf{a}_k \int_{-\infty}^{\infty} d\omega e^{\frac{i\omega x}{a}} \chi^*_\omega \times \left[ \Gamma(-\frac{i\omega}{a}) | \frac{k}{a} |^{\frac{i\omega}{a}} \theta(k) + \Gamma(\frac{i\omega}{a}) | -\frac{k}{a} |^{\frac{i\omega}{a}} \theta(-k) \right] \times \left[ | au |^{\frac{i\omega}{a}} \theta(-u) - \lambda + | av |^{\frac{i\omega}{a}} \theta(v) \theta(\lambda) \right]
\] (2.21)

where \( \phi_{int}(x, t) = \phi(x, t) - \phi_0^+(x, t) \) accounts for the interaction of the quantum field with the oscillator.

C. Two point function of the field

In order to evaluate the two-point function \( \langle \phi(x, t)\phi(x', t') \rangle \) in the Minkowski vacuum, we first recognize that it is equal to \( \langle \phi^+(x, t)\phi^-(x', t') \rangle \). This is because only operator products of the form \( \mathbf{a}_k \mathbf{a}^\dagger_k \) contribute when taking the expectation value in the Minkowski vacuum. Denoting \( \langle \phi(x, t)\phi(x', t') \rangle \) by \( G(x, t; x', t') \) and \( \langle \phi_0(x, t)\phi_0(x', t') \rangle \) by \( G_f(x, t; x', t') \), we obtain

\[
G(x, t; x', t') = G_f(x, t; x', t') + \langle \phi^+(x, t)\phi^-(x', t') \rangle + \langle \phi^+(x, t)\phi^-(x', t') \rangle + \langle \phi^-(x, t)\phi^-(x', t') \rangle.
\] (2.22)

Using the expression for \( \phi^+_0(x, t) \) and \( \phi^-_{int}(x', t') \) we obtain

\[
G(x, t; x', t') = -\frac{\gamma}{2\pi} \int_{-\infty}^{\infty} \frac{d\omega}{\omega} (1 - e^{-\frac{2\pi \omega}{a}})^{-1} \times \left[ \left| \frac{au}{au'} \right|^{\frac{i\omega}{a}} \theta(-u) \theta(-u') \right] \times \left\{ \chi^*_\omega \theta(-\lambda) + \chi_\omega \theta(-\lambda') - 4\gamma | \chi_\omega |^2 \theta(-\lambda) \theta(-\lambda') \right\}
\]

6
The role of the relation (2.15) in the cancellation of various terms in the first half of the above expression is thus made explicit. Different terms will vanish depending on which region of Minkowski space each of the two points is in, and according to whether these points are to the left or the right of the accelerated trajectory.

The above result can also be obtained via a different quantization procedure. Instead of expanding the field in Minkowski modes, as above, we can use Unruh modes which are linear combinations of Rindler modes and positive frequency with respect to Minkowski time (see, for example [16]). These modes are easier to handle in the manipulations involved. However, they have the disadvantage of being defined differently in each region (F, P, R and L). Although the Rindler modes are defined only in R and L, the Unruh modes, as linear combinations of Rindler modes, can be analytically extended to the entire spacetime. One then computes the two point function in a desired region by expanding the field in a complete set of Unruh modes as defined by analytic extension to that region. Of course, one always needs the mode decomposition in R, because the field operator at an arbitrary point depends both on the free field operator at that point as well as on the accelerated trajectory, which lies in R (see equations (2.8) and (2.18)). This procedure will not be repeated here, as it leads to the same result.

D. Energy Momentum Tensor

Let us first consider the coincidence limit of the two point function. In that case all terms involving $u$ and $u'$ or $v$ and $v'$ vanish as a consequence of the relation (2.15). The remaining terms can be simplified to give:

$$\langle \varphi^2(x, t) \rangle - \langle \varphi_0^2(x, t) \rangle = -\frac{\gamma}{2\pi} q(v) \int_{-\infty}^{\infty} \frac{d\omega}{\omega} (1 - e^{-2\pi\omega})^{-1} \times$$

$$\left[ |a^2uv| \frac{i\omega}{\omega} \left\{ \chi_+^\ast \theta(-u)\theta(-\lambda) + \chi_\omega \theta(\lambda)\theta(u)e^{-\frac{\pi\omega}{a}} + \theta(-u) \right\} \right] + |a^2uv| \frac{i\omega}{\omega} \left\{ \chi_\omega \theta(-u)\theta(-\lambda) + \chi_+^\ast \theta(\lambda)\theta(u)e^{-\frac{\pi\omega}{a}} + \theta(-u) \right\} \right].$$ (2.24)
This corresponds to a static polarization cloud confined to the region $F \cup R$, i.e. $v > 0$. It is static because it is a function of $uv = t^2 - x^2$ in each region. Thus it is constant along any accelerated world line in particular. In $F$, the curves $t^2 - x^2 = \text{constant}$ are spacelike curves and therefore do not correspond to world-lines of physical particles. Therefore any physical detector in $F$ will respond to the field in a non-trivial, time-dependent way.

However, it is simple to show that the renormalized energy-momentum tensor of the field vanishes everywhere except at the past null horizon $v = 0$ of the accelerated trajectory, and on the accelerated trajectory itself. The energy-momentum tensor is renormalized by subtracting out the free field contribution. It is thus given by

$$T_{uu} = \lim_{u' \to u, v' \to v} \partial_u \partial_{u'} (G(x, t; x', t') - G_f(x, t; x', t'))$$
$$T_{vv} = \lim_{u' \to u, v' \to v} \partial_v \partial_{v'} (G(x, t; x', t') - G_f(x, t; x', t'))$$
$$T_{uv} = 0. \quad (2.25)$$

Going back to the expression (2.23) for the two-point function, we find, in $P \cup L$ (i.e. $v, v' < 0$), that $G(x, t; x', t') - G_f(x, t; x', t') = 0$. Thus the renormalized energy momentum tensor trivially vanishes there. In the region $F \cup R$, and to the left of the trajectory, $\lambda, \lambda' > 0$, we have

$$G(x, t; x', t') - G_f(x, t; x', t') = -\frac{\gamma}{2\pi} \int_{-\infty}^{\infty} \frac{d\omega}{\omega} \left(1 - e^{-\frac{2\pi\omega}{a}}\right)^{-1}$$
$$\times \left[ a^2 uv \Gamma^{\omega \lambda} \chi_\omega \left(\theta(-u) + \theta(u) e^{-\frac{2\pi\omega}{a}}\right) + a^2 u' v \Gamma^{-\omega \lambda} \chi^*_\omega \left(\theta(-u') + \theta(u') e^{-\frac{2\pi\omega}{a}}\right) \right]. \quad (2.26)$$

The terms involving $u$, $u'$ and $v$, $v'$ all vanish as a consequence of (2.15). The remaining cross-terms do not contribute to the energy-momentum tensor, as can be checked by straightforward differentiation.

Similarly, to the right of the trajectory, $(\lambda, \lambda' < 0)$, we obtain

$$\langle \phi(x, t)\phi(x', t') \rangle - \langle \phi_0(x, t)\phi_0(x', t') \rangle = -\frac{\gamma}{2\pi} \int_{-\infty}^{\infty} \frac{d\omega}{\omega} \left(1 - e^{-\frac{2\pi\omega}{a}}\right)^{-1}$$
$$\times \left[ a^2 uv \Gamma^{\omega \lambda} \chi_\omega + a^2 u' v \Gamma^{-\omega \lambda} \chi^*_\omega \right]. \quad (2.27)$$

The energy-momentum tensor vanishes here as well, in a similar fashion.

If we therefore consider a world-tube formed by two accelerated world-lines with $\lambda > 0$ and $\lambda < 0$ in the Rindler wedge, then this tube encloses the accelerated trajectory $\lambda = 0$. Also the energy-momentum tensor vanishes everywhere on the boundary of the tube. Hence there is no flux of energy-momentum, or radiation from the oscillator at $\lambda = 0$.

The cross-terms in $u, v'$ and $u', v$ which appear in the above expressions are missing in RSG. Although we have found that they do not contribute to the energy-momentum tensor, they do signal the presence of a polarization cloud around the oscillator. These results
support those of Unruh [12] and MPB [13]. However, the above analysis has the advantage of clearly displaying the role of the “fluctuation-dissipation relation” (2.15) in the cancellation of terms which would naively be expected to contribute to the energy-momentum. Also, we have here computed an expression for the two-point function which is valid over the entire spacetime. This is a generalization of previous work.

Calculation is underway in four dimensional spacetime, which is certainly more physical. In this case [17] we expect to see the ordinary classical radiation of the Larmor type from a uniformly accelerated charge, but the question of interest to us is whether in 4D there is emitted radiation. If there were it should manifest in the content of the energy momentum tensor and, being of quantum nature, discernible from the classical radiation. This would further clarify any existing confusion on the nature of Unruh radiation.

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