An uniformly accelerated quantum counting detector in Minkowski and Fulling vacuum states

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Abstract

In this work we discuss the process of measurements by a detector in an uniformly accelerated rectilinear motion, interacting linearly with a massive scalar field. The detector model for field quanta is a point-like system with a ground state and a continuum of unbounded states. We employ the Glauber theory of photodetection. In an uniformly accelerated reference frame, the detector, interacting with the field prepared in an arbitrary state of the Rindler Fock space, is excited only by absorption processes. For the uniformly accelerated detector prepared in the ground state, we evaluate the transition probability rate in three important situations. In the first one the field is prepared in an arbitrary state of \(n\)-Rindler quanta, then we consider a thermal Rindler state at a given temperature \(\beta^{-1}\), and finally the case in which the state of the field is taken to be the Minkowski vacuum. The well-known result that the latter excitation rates are equal is recovered. Accelerated or inertial observer interpretations of the measurements performed by the accelerated detector is presented. Finally, we investigate the behaviour of the detector in a frame which is inertial in the remote past but in the far future becomes uniformly accelerated. For the massless case, we obtain that the transition probability rate of the detector in the far future is tantamount to the analogous quantity for the detector at rest in a non-inertial reference frame interacting with the field prepared in an usual thermal state.

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1 Introduction

Quantum mechanics is a mathematical and conceptual framework that provides an algorithm to compute probabilities for the various possible outcomes of a set of measurements. This formalism is able to discuss the behaviour of systems modelled by a finite number of degrees of freedom. After choosing the algebra of operators and constructing the states space, one key point is that the description of the system cannot depend on the choice of representation. The textbook example is given by the Schrödinger and Heisenberg representations. Since in quantum mechanics experimental results are expressed in terms of matrix elements of linear operators between state vectors, the Schrödinger and Heisenberg pictures are alternative ways to describe how observables evolve in time. One can freely change from one representation to the other by performing unitary transformations. There are no physical consequences in choosing one (or the other) of such representations of the operator algebra. It is well-known that, for systems described by a finite number of degrees of freedom, all the irreducible representations of the algebra (technically a $C^*$-algebra) are unitary equivalent.

For the description of systems with countably infinite number of degrees of freedom, one must define the so-called canonical variables. Then one typically uncovers the algebra determined by such variables and, equipped with such an algebra, one is able to build up the Hilbert space of states of the theory. Nevertheless, it is possible to present unitarily inequivalent representations of the canonical commutation relations [1, 2]. The result is that each representation may depict different physical situations. Hence a fundamental point regarding systems described by an infinite number of degrees of freedom is the representation problem [3, 4]. In this situation new phenomena such as collective excitations, which can arise in physical systems undergoing spontaneous symmetry breaking, may appear, as is the case of the vacuum of a superconductor with a condensate of Cooper pairs.

The particle interpretation of field quantum theory offers a link between the mathematical structure of the theory and particle observables as mass and spin. For free fields in Minkowski spacetime, it is possible to have a corpuscular interpretation from a choice of states that span the Hilbert space. Such an interpretation that emerges from a quantum field theory is based in the commutation of fields at spacelike separated points, Poincaré invariance and spectral condition. However, one can abandon the isometries generated by the Poincaré invariance and still have the corpuscular content. For instance, cavity quantum electrodynamics is the study of radiative properties of atoms in the presence of boundaries. The presence of boundaries results in modified vacuum fluctuations, since one has to work with an infinite countable number of normal modes in order to set up the Hilbert space of states associated with the electromagnetic field. In this case, vacuum fluctuations can inhibit or enhance atomic radiative processes [5, 6]. Nevertheless, one can conceive another situation where the presence of boundaries yields to a dramatic change in the Hilbert space of quantum fields. This can be achieved through the investigation of a generic quantum field in the presence of an infinite plane distribution of matter [7, 8, 9, 10]. This idealised matter configuration is able to generate a global static gravitational field.

If one evokes the equivalence principle, an alternative route is to study quantum fields from the perspective of observers in rectilinear uniformly accelerated motion. This leads us to the Fulling quantization [11, 12] and the Unruh-Davies effect [13, 14]. According to the Unruh-Davies effect, a detector at rest in an uniformly accelerated frame of reference can be excited by absorption of a Rindler field quantum from the Minkowski vacuum. Our purpose here is to re-examine the Unruh-Davies effect using the Glauber theory of photodetection.
Although this theory was constructed to formulate models of detectors that respond to the electromagnetic field by absorption of photons, it is straightforward to implement the same idea to an apparatus device that measures quanta associated with a massive scalar field. In the Glauber theory of photodetection, measurements are absorption processes which are registered, e.g., by the detection of emitted photoelectrons. Such a kind of measurements is known as measurement of second kind, inasmuch as it changes the state being measured. In measurements of first kind, which we will not discuss here, the system is not modified by the measurement procedure.

Closely connected to the problem of how to obtain the excitation spectrum of quantum field theory with coupled Green’s functions, we shall address the question of how to define a suitable detector. From an operational point of view, measurements are means to obtain information about the physical reality using a measurement apparatus. The measurement apparatus sends a sign of information, a signal that must be decoded by some receiver that is able to give us a measurable value at the output of the system. When measured systems are described by quantum mechanics, the dimension of the Hilbert space of the measuring device is usually larger than that of the measured system. In this regard, it is possible to discuss how to optimize extraction of information from a finite system. We consider the case in which the dimension of the Hilbert space of the measured system is infinite. In this situation, an ideal detector of field quanta is a dimensionless system that can make a transition between two energy levels, decreasing the number of quanta of the field in some measured state. Therefore, a detector of field quanta is an experimental device coupled with the field that gives no signal if the state of the field is the ground state. We should keep in mind that the definition of a vacuum state is related to the choice of time-translation Killing vector fields used by the observers that quantize a classical field.

The aim of this paper is to explore second-kind measurements in a situation extensively discussed in the literature: Measurements of field quanta associated with a (massive) field performed by observers in rectilinear uniformly accelerated motion. In the quantum theory of photodetection, a perturbative treatment of the detector-field interaction using the rotating-wave approximation (RWA) is usually adopted. In this context, some issues involving causality have already been raised. However, a thorough analysis reveals that such concerns are unwarranted. For a related discussion of causality in disordered settings, see Ref. [23].

In this work we discuss the process of measuring Rindler field quanta by an uniformly accelerated detector interacting with a massive scalar field using the Glauber theory of photodetection. We evaluate the transition probability rates of the accelerated detector prepared in the ground state for the cases in which the state of the field is taken to be an arbitrary state of $n$-Rindler quanta, a Rindler thermal state with temperature $\beta^{-1}$ and finally the Minkowski vacuum. We recover the well-known result that a detector in a rectilinear uniformly accelerated motion interacting with a field in the Minkowski vacuum has the same transition probability rate of an accelerated detector interacting with the field in a Rindler thermal state at a temperature $\beta^{-1}$. Since, there is no agreement in the literature concerning the interpretation of the processes described by observers at rest in an inertial reference frame and at rest in a non-inertial reference frame, using the Glauber model, we intend to shed some light on this issue. Ultimately, we investigate a model in the Minkowski spacetime where the vacuum defined in the remote past, the state $|0,\text{in}\rangle$, and in the far future, the state $|0,\text{out}\rangle$, are unitarily non-equivalent. We show that a Glauber detector in the far future can be excited when interacting with the field in the vacuum state $|0,\text{in}\rangle$. 

[References]

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The organization of this paper is as follows. In Sec. 2, the Fulling quantization, suitable for uniformly accelerated observers and for addressing the Unruh-Davies effect, is briefly discussed. In Sec. 3 the Glauber theory of detection of field quanta, interpreting measurements by absorption of field quanta, is presented. In Sec. 4 we perform the analysis of the radiative processes of the detector in an uniformly accelerated frame. In Sec. 5, we discuss a Glauber detector for a Kalnis-Miller coordinate system. In Sec. 6 we present our conclusions. We use units such that $\hbar = c = k_B = 1$ and $\eta_{\mu\nu} = \text{diag}(+1, -1, -1, -1)$.

## 2 The Fulling quantization and the Unruh-Davies effect

In this section, we briefly review the discussion on the Unruh-Davies effect, with the Fulling quantization of a scalar field in Rindler spacetime [28]. See also Refs. [29, 30, 31]. The main result discusses the behaviour of an Unruh-DeWitt detector at rest in a frame of reference with a rectilinear uniformly accelerated motion. One can show that this detector interacting with a scalar field in the Minkowski vacuum $|0, M\rangle$ has the same transition probability rate of the detector in a rectilinear uniformly accelerated motion interacting with the field in a Rindler thermal state at a temperature $\beta^{-1}$.

Since in an arbitrary globally hyperbolic stationary spacetime one can always find a Killing vector field corresponding to the time direction of some family of inertial observers, in Minkowski spacetime one can use this feature to decompose the free scalar field operator into a sum of its positive and negative frequency parts. Both contributions are given by

$$\varphi(t, x) = \varphi^+(t, x) + \varphi^-(t, x)$$

where

$$\varphi^+(t, x) = \frac{1}{(2\pi)^{3/2}} \int \frac{d^3k}{\sqrt{2\omega_k}} a(k)e^{-i(\omega_k t - k \cdot x)}$$

and

$$\varphi^-(t, x) = \frac{1}{(2\pi)^{3/2}} \int \frac{d^3k}{\sqrt{2\omega_k}} a^\dagger(k)e^{i(\omega_k t - k \cdot x)}$$

for $\omega_k = \sqrt{k^2 + m_0^2}$, $m_0$ being the mass of the field. Given a set of solutions of the Klein-Gordon equation, $\{\phi_i(x)\}$, these modes are normalized using the following inner product

$$(\phi_1, \phi_2) = i \int_{\Sigma} \phi_1^\dagger \partial_\alpha \phi_2 n^\alpha d\Sigma,$$

where $n^\alpha$ is future pointing normal vector to a spacelike hypersurface $\Sigma$ and $d\Sigma$ is the volume element of $\Sigma$ [32]. The annihilation and creation operators for quanta of the field, $a(k)$ and $a^\dagger(k)$ satisfy the usual commutation relations

$$[a(k), a(k')] = [a^\dagger(k), a^\dagger(k')] = 0,$$

$$[a(k), a^\dagger(k')] = \delta(k - k').$$

The vacuum state built by observers at rest in an inertial reference frame that have a timelike Killing vector field is the state in which field quanta are absent. It is also the lowest energy state of the respective Fock space of states, which are infinite direct sums of tensor products of
single-excitation Hilbert spaces. In conclusion, there is a translational-invariant vacuum state $|0, M\rangle$ such that

$$a(k)|0, M\rangle = 0 \quad \forall k.$$  \hfill (7)

For instance, although the zero-point energy can be eliminated from the Hamiltonian by Wick normal ordering, one can show that this procedure is correct only in Minkowski spacetime in the absence of boundaries. The Casimir effect shows that one must be careful when imposing the normal ordering \cite{33, 34, 35, 36}.

Before discussing the quantization in a non-inertial frame, let us introduce a function associated with the microscopic causality, \textit{i.e.}, at spacelike separated points $(x-y)^2 < 0$ the scalar field operators commute. The invariant Pauli-Jordan function $\Delta(x; m^2_0)$ is defined as \cite{37, 38}

$$[\varphi(x), \varphi(y)] = i\Delta(x - y; m^2_0).$$  \hfill (8)

Using the decomposition of the field operator into a sum of its positive and negative frequency parts, one finds that

$$i\Delta(x; m^2_0) = \Delta^{(+)}(x; m^2_0) + \Delta^{(-)}(x; m^2_0)$$  \hfill (9)

where

$$\Delta^{(+)}(x - y; m^2_0) = [\varphi^{(+)}(x), \varphi^{(-)}(y)],$$

$$\Delta^{(-)}(x - y; m^2_0) = [\varphi^{(-)}(x), \varphi^{(+)}(y)].$$  \hfill (10, 11)

Consider now a family of observers at rest in a non-inertial reference frame, \textit{e.g.}, with a rectilinear uniformly accelerated motion. Starting from the usual Cartesian coordinates $x^\mu = (t, x^1, x^2, x^3)$, one defines the curvilinear coordinates $X^\mu = (\eta, \xi, y, z)$ using

$$
\begin{align*}
t &= \xi \sinh \eta \\
x^1 &= \xi \cosh \eta \\
x^2 &= y \\
x^3 &= z,
\end{align*}
$$  \hfill (12)

where $0 < \xi < \infty$ and $-\infty < \eta < \infty$. Therefore $\xi^2 = (x^1)^2 - t^2$ and $\eta = \tanh^{-1}(t/(x^1))$. An observer travelling in the wordline $\xi = 1/a = \text{constant}$ and $y, z$ are also constants has proper acceleration given by $a$. This coordinate system covers only a wedge of the Minkowski spacetime, \textit{i.e.} the region $|x| > t$, where there is a global timelike Killing vector field $\partial/\partial \eta$. The line element using the above curvilinear coordinates can be written as

$$ds^2 = \xi^2 d\eta^2 - d\xi^2 - dy^2 - dz^2.$$  \hfill (13)

There are two causally-independent wedges called the left and right Rindler wedge. Since the wedges are causally disconnected, we restrict our considerations to the right wedge where the detector is travelling. The same approach was put forward in Ref. \cite{39}. Define the vectors $\mathbf{q} = (k_y, k_z)$ and $\mathbf{y} = (y, z)$. The Klein-Gordon equation after separation of variables is written as a Bessel equation. From the general solutions of the Klein-Gordon equations $\phi_{vq}(\eta, \xi, y)$ and $\phi_{vq}^*(\eta, \xi, y)$, the scalar field operator can be written as sum of positive and negative frequency contribution with respect to $\partial/\partial \eta$ as $\varphi(\eta, \xi, y) = \varphi^{(+)}(\eta, \xi, y) + \varphi^{(-)}(\eta, \xi, y)$. Defining $m^2 = m^2_0 + q^2$, we have

$$\varphi^{(+)}(\eta, \xi, y) = \frac{1}{2\pi^2} \int_0^{\infty} d\nu \int d^2 \mathbf{q} \sqrt{\sinh \pi \nu} K_{\nu}(m\xi) e^{-i(\nu y - \mathbf{q} \cdot \mathbf{y})} b(\nu, \mathbf{q}),$$  \hfill (14)
where $K_\nu(z)$ is the Macdonald function of imaginary order. The annihilation and creation operator of Rindler field quanta $b(\nu, \mathbf{q})$ and $b^\dagger(\nu, \mathbf{q})$ satisfy the usual commutation relations over a constant $\eta$ hypersurface

$$[b(\nu, \mathbf{q}), b(\nu', \mathbf{q}')] = [b^\dagger(\nu, \mathbf{q}), b^\dagger(\nu', \mathbf{q}')] = 0, \quad (16)$$

$$[b(\nu, \mathbf{q}), b^\dagger(\nu', \mathbf{k}')] = \delta(\nu - \nu')\delta(\mathbf{q} - \mathbf{q}'). \quad (17)$$

Eqs. (14) and (15) can be seen as the consequence of a general result which states that it is possible to expand an arbitrary function as an integral with respect to cylinder functions of imaginary order and arguments [40]. It is clear that in the Hilbert space of states there is a vacuum state $|0, R\rangle$, known as the Fulling vacuum state, such that

$$b(\nu, \mathbf{q})|0, R\rangle = 0, \quad \forall \nu \in [0, \infty), -\infty < k_y, k_z < \infty. \quad (18)$$

From the Fulling vacuum one can generate the Rindler Fock space. The construction is similar to the inertial case where one starts with the basis state for the one excitation Hilbert space. Although this field is massive, this situation is similar to the case of a zero mass theory where the vacuum is not an isolated state but rather a unit point of continuum energy excitation. One solution of the infrared problem at least in quantum electrodynamics is the fact that in the experiments, it is not possible to detect soft photons [41].

The annihilation operator of Rindler field quanta in the mode $(\nu, \mathbf{q})$ can be expanded, using Bogoliubov transformations, into a linear combination of creation and annihilation of Minkowski field operators $a^\dagger(\mathbf{k})$ and $a(\mathbf{k})$. We have

$$b(\nu, \mathbf{q}) = \int d^3k \left[ U(\nu, \mathbf{q}, \mathbf{k})a(\mathbf{k}) + V(\nu, \mathbf{q}, \mathbf{k})a^\dagger(\mathbf{k}) \right], \quad (19)$$

where

$$U(\nu, \mathbf{q}, \mathbf{k}') = \frac{1}{\left[2\pi\omega_k(1 - e^{-2\pi\nu})\right]^{\frac{3}{2}}} \left( \frac{\omega_k + |\mathbf{k}|}{m} \right)^{i\nu} \delta(\mathbf{q} - \mathbf{q}'), \quad (20)$$

$$V(\nu, \mathbf{q}, \mathbf{k}') = \frac{1}{\left[2\pi\omega_k(e^{2\pi\nu} - 1)\right]^{\frac{3}{2}}} \left( \frac{\omega_k + |\mathbf{k}|}{m} \right)^{i\nu} \delta(\mathbf{q} - \mathbf{q}'). \quad (21)$$

So we see that, in the construction of the mathematical formalism, we obtain two inequivalent Fock spaces with two vacuum states $|0, M\rangle$ and $|0, R\rangle$. It turns out that the Poincaré invariant vacuum $|0, M\rangle$ has an infinite number of Rindler quanta in the mode $(\nu, \mathbf{q})$ associated with the massive scalar field. To show this result, we can use the Bogoliubov coefficients of (19) and compute the number of Rindler quanta in the Minkowski vacuum. One obtains

$$\langle 0, M | b^\dagger(\nu, \mathbf{q})b(\nu', \mathbf{q}') | 0, M \rangle = \frac{1}{e^{2\pi\nu} - 1} \delta(\nu - \nu') \delta(\mathbf{q} - \mathbf{q}'). \quad (22)$$

In order to provide a physical interpretation, one must address the query of how energy can be transferred between the system under study and the measurement devices. In the quantum
theory of measurements, the information received from the state of the field by registering clicks corresponds to the energy transfer from the field system to the detector. To complete the measurement process one must obtain information from the measuring device.

Let us briefly discuss basic results concerning a two-level Unruh-DeWitt detector. It consists of an idealized point-like two-level system with a ground state $|g\rangle$ and an excited one $|e\rangle$ interacting linearly with a scalar field. The detector-field interaction Hamiltonian is given by

$$H_{\text{int}} = c_1 m(\tau) \phi(x(\tau)),$$

where $m(\tau)$ is called the monopole operator and $c_1$ is a small coupling constant. Here $\tau$ is the proper time of the detector. In interacting field theory the interaction between different fields is switched on and off slowly. This is known as the adiabatic hypothesis. Here with a quantum mechanical treatment of the apparatus device we are interested to discuss the behaviour of the detector interacting with the field in prepared states. One can find the transition probability from the initial state $|\tau_i\rangle = |g\rangle \otimes |\Phi_i\rangle$ at $\tau_i$ to a final state $|\tau_f\rangle = |e\rangle \otimes |\Phi_f\rangle$ at $\tau_f$, where $|\Phi_i\rangle$ is the initial field state and $|\Phi_f\rangle$ is the final field state. After summing over a complete set $\{|\Phi_f\rangle\}$, one obtains the so-called response function, in first-order perturbation theory:

$$F(\omega_{eg}, \tau_f, \tau_i; x) = \int_{\tau_i}^{\tau_f} d\tau \int_{\tau_i}^{\tau_f} d\tau' e^{-i\omega_{eg}(\tau-\tau')} \langle \Phi_i | \varphi(\tau, x) \varphi(\tau', x) | \Phi_i \rangle,$$

where $\omega_{eg}$ is the energy gap between the energy levels of the detector and we used that $\sum_f |\Phi_f\rangle\langle\Phi_f| = 1$. The response function reveals the bath of quanta that the detector may experience. For the case of an uniformly accelerated Unruh-DeWitt detector and a real massless scalar field in the Minkowski vacuum, the transition probability was evaluated for a finite proper time in Ref. [42]. The asymptotic rate of spontaneous and induced emission and absorption of Rindler field quanta is given by

$$R(\omega_{eg}) = \frac{|\omega_{eg}|}{2\pi} \left[ \theta(-\omega_{eg}) \left( 1 + \frac{1}{e^{2\pi\sigma\omega_{eg}} - 1} \right) + \theta(\omega_{eg}) \frac{1}{e^{2\pi\sigma\omega_{eg}} - 1} \right],$$

where $\sigma^{-1}$ is the proper acceleration of the detector. See also Refs. [43, 44]. Although this is a fairly simple calculation, the interpretation of how the process of measurement is described by different observers is far from being obvious. Hence a fundamental question would be how one should interpret detector measurements in different frames. This is the topic of the next sections.

3 The Glauber theory of photodetection

In this section we are interested to discuss measurements in the scenario of radiative processes of prepared states. We wish to consider the case of real atoms for which, e.g., the continuum of excited states (in a single electron atom) corresponds to electronic states above the ionization threshold.

The detector Hilbert space of states is the ground state $|\psi_b\rangle$ and a continuum of unbounded states $|\psi_j\rangle$. The time independent detector Hamiltonian satisfies

$$H_d|\psi_l\rangle = E_l|\psi_l\rangle, \quad l = b \text{ or } j$$

where $E_b$ ($E_j$) is the energy associated with the ground state (excited states). The detector-field interaction Hamiltonian is still given by Eq. (23), but now $\tau$ is not necessarily the detector's
proper time; here it is just a parameter with respect to which we describe the time evolution of the system. One can find the transition amplitude from the initial state $|\tau_i\rangle = |\psi_i\rangle \otimes |\Phi_i\rangle$ at $\tau_i$ to $|\tau_f\rangle = |\psi_f\rangle \otimes |\Phi_f\rangle$ at $\tau_f$, where $|\Phi_i\rangle$ and $|\Phi_f\rangle$ are arbitrary field states. For instance, they can be Fock states constructed with $a^\dagger(k)$ operators acting in the Poincaré invariant vacuum state $|0, M\rangle$. Let us define $\omega_{jb} = E_j - E_b$. For excitation processes, $\omega_{jb} > 0$. This is the case we are interested in. As discussed one can derive the response function $F(\omega_{jb}, \tau_f, \tau_i, x)$. One finds

$$F(\omega_{jb}, \tau_f, \tau_i, x) = \int_{\tau_i}^{\tau_f} d\tau' \int_{\tau_i}^{\tau_f} d\tau'' e^{-i\omega_{jb}(\tau' - \tau'')} \langle \Phi_i | \varphi(\tau', x) \varphi(\tau'', x) | \Phi_i \rangle. \quad (27)$$

The response function can be written as

$$F(\omega_{jb}, \tau_f, \tau_i, x) = F_1(\omega_{jb}, \tau_f, \tau_i, x) + F_2(\omega_{jb}, \tau_f, \tau_i, x), \quad (28)$$

where

$$F_1(\omega_{jb}, \tau_f, \tau_i, x) = \int_{\tau_i}^{\tau_f} d\tau' \int_{\tau_i}^{\tau_f} d\tau'' e^{-i\omega_{jb}(\tau' - \tau'')} \langle \Phi_i | \varphi(-)(\tau', x) \varphi(+)(\tau'', x) | \Phi_i \rangle, \quad (29)$$

and

$$F_2(\omega_{jb}, \tau_f, \tau_i, x) = \int_{\tau_i}^{\tau_f} d\tau' \int_{\tau_i}^{\tau_f} d\tau'' e^{-i\omega_{jb}(\tau' - \tau'')} \langle \Phi_i | \varphi(+)(\tau', x) \varphi(-)(\tau'', x) | \Phi_i \rangle. \quad (30)$$

Our definition of a bona fide detector is a device that goes to an excited state by decreasing the number of quanta of some state. It is clear that $F_1(\omega_{jb}, \tau_f, \tau_i, x)$ describes absorption processes. In the Glauber theory of photodetection only this term contributes to the transition rate. On the other hand, $F_2(\omega_{jb}, \tau_f, \tau_i, x)$ is commonly associated with emission processes, accompanied by the decay of the detector. However, since we are considering $\omega_{jb} > 0$, one may ask whether one should really maintain such a contribution. Indeed, as can be easily checked, the $F_2$ contribution vanishes in the asymptotic limits $\tau_i \to -\infty, \tau_f \to \infty$ when $\varphi(+)x$ and $\varphi(-)x$ are given by Eqs. (2) and (3) and $|\Phi_i\rangle$ is the Minkowski vacuum state. Notwithstanding, one can still rewrite this term in such a way that it displays a contribution to virtual processes. Using the Pauli-Jordan function we can write

$$\varphi(+)(x) \varphi(-)(y) = \varphi(-)(y) \varphi(+)(x) + \Delta(+)x - y; m^2. \quad (31)$$

The last term of the above equation (which is independent of the state of the field) accounts for spontaneous emission processes. Since we are interested in absorption processes, we disregard this term. Collecting our results, one finds that

$$\tilde{F}_2(\omega_{jb}, \tau_f, \tau_i, x) = \int_{\tau_i}^{\tau_f} d\tau' \int_{\tau_i}^{\tau_f} d\tau'' e^{-i\omega_{jb}(\tau' - \tau'')} \langle \Phi_i | \varphi(-)(\tau'', x) \varphi(+)(\tau', x) | \Phi_i \rangle. \quad (32)$$

where $\tilde{F}_2$ is the $F_2$ contribution after neglecting the term associated with the Pauli-Jordan function. In any case, this anomalous contribution also vanishes in the asymptotic limits $\tau_i \to -\infty, \tau_f \to \infty$ in the same situation considered above and hence it does not contribute to the asymptotic rate in this case.

Henceforth let us assume a broadband detector. For simplicity we choose $\tau_i = 0$ and $\tau_f = \tau$. The density of final excited states of the detector is defined by $\rho(\omega_{jb})$ [45]. The probability of excitation is

$$P(\tau, x) = \int \rho(\omega_{jb}) P_{|\psi_i\rangle \to |\psi_f\rangle}(\tau, x) d\omega_{jb}. \quad (33)$$
Considering that $\rho(\omega_{jb})$ is a slowly-varying function, one finds
\[
\int_{-\infty}^{\infty} d\omega_{jb} e^{i\omega_{jb}(\tau''-\tau')} \rho(\omega_{jb}) = 2\pi \delta(\tau''-\tau') \rho(\bar{\omega}_{jb}).
\] (34)

In conclusion, for a broadband detector, with only the $F_1(\omega_{jb}, \tau_f, \tau_i, \mathbf{x})$ contribution, one can show that the excitation probability is given by
\[
P(\tau, \mathbf{x}) = C \int_0^\tau d\tau' \langle \Phi_i | \varphi^(-)(\tau', \mathbf{x}) \varphi^+(\tau', \mathbf{x}) | \Phi_i \rangle,
\] (35)
where $C$ represents the efficiency of the detector. The probability of excitation per unit time at time $\tau$ is defined by $dP(\tau, \mathbf{x})/d\tau$. Therefore, the transition probability rate of the detector reads
\[
\frac{dP}{d\tau} = W(\tau, \mathbf{x}; |\Phi_i\rangle) = C \langle \Phi_i | \varphi^(-)(\tau, \mathbf{x}) \varphi^+(\tau, \mathbf{x}) | \Phi_i \rangle.
\] (36)
This is the standard result of the Glauber quantum counting model. The transition probability rate comes from absorption processes, where $\varphi^+(\tau, \mathbf{x})$ and $\varphi^-(\tau, \mathbf{x})$ are given by Eqs. (2) and (3).

4 The Glauber theory of photodetection in the Unruh-Davies effect

We are now in the position to investigate radiative processes of the Glauber detector at rest in an uniformly accelerated frame of reference. The important question that emerges now is how accelerated or inertial observers should interpret measurements performed by the uniformly accelerated detector. In the Glauber theory the detector works only by absorption of field quanta. In order to understand the essential point of our approach recall that in, stationary spacetimes such as the Minkowski or Rindler spacetime, observers use different choices of timelike Killing fields. Therefore one must define two different orderings, given by $\varphi^(-)(x) \varphi^+(x') : m$, using Eqs. (2) and (3) or $\varphi^-(x) \varphi^+(x') : a$, using Eqs. (14) and (15) related to each Killing vector field respectively.

Let us study the broadband detector introduced above, but now we explicitly take into account an uniformly accelerated motion. This detector is interacting with a massive scalar field. In this case the transition probability rate must be given by products of $\varphi^+(x)$ and $\varphi^-(x)$ defined in Eqs. (14) and (15), respectively. First we will consider that the field is in an arbitrary state $|\phi_i\rangle$ constructed with the $\hat{b}^\dagger(\nu)$ operators acting on the Fulling vacuum state $|0, R\rangle$. Since we are using the above defined ordering associated with the Rindler timelike Killing vector field, similar steps as those previously outlined lead us to the following expression for the transition probability rate $W(\eta, \xi; |\phi_i\rangle)$:
\[
W(\eta, \xi; |\phi_i\rangle) = \langle \phi_i | \varphi^-(\eta, \xi, y) \varphi^+(\eta, \xi, y) | \phi_i \rangle
\] (37)
where we are employing the Rindler time $\eta$ in order to describe the time evolution of the system and for simplicity we have set $C = 1$. Observe that the average Rindler quantum counting rate is proportional to the expectation value of the normal ordered product of the negative and
positive frequency parts contributions of the field operator at the worldline $\xi = \text{constant}$ and $y, z$ also constant. Using Eqs. (14) and (15), one finds that

$$W(\xi; |\phi_i\rangle) = \frac{1}{4\pi^4} \int_0^\infty d\nu \int_0^\infty d\nu' \int d^2q \int d^2q' e^{i[\eta(\nu-\nu') - y(a-a')]} \sqrt{\sinh \pi\nu} \sqrt{\sinh \pi\nu'} \times K_{\nu\nu}(m\xi) K_{\nu'\nu'}(m\xi) \langle \phi_i\rangle b^\dagger(\nu, q) b(\nu', q') |\phi_i\rangle. \quad (38)$$

It is now clear that, when the state of the Rindler Fock space is the Fulling vacuum, i.e., $|\phi_i\rangle = |0, R\rangle$, the rate of excitation vanishes, as expected.

Now let us assume a thermal bath in Rindler spacetime. Each observer at rest in an accelerated frame of reference measures a local temperature $T = \beta^{-1}$. Planck’s law shows that the transition rate in a specific worldline is

$$W_\beta(\xi) = \frac{1}{4\pi^4} \int d^2q \int_0^\infty d\nu \frac{\sinh \pi\nu}{e^{\beta\nu} - 1} K_{\nu\nu}^2 \left(\xi \sqrt{m_0^2 + q^2}\right). \quad (39)$$

Now, let us make the connection between temperature and proper acceleration. A family of observers can be defined by a set of timelike worldlines. In Rindler spacetime each observer travelling in the worldline $\xi = \text{constant}$ defines a uniformly accelerated frame with proper acceleration $a = \xi^{-1}$. One can show that there is a thermal equilibrium using the Tolman relation $\beta^{-1} \sqrt{g_{00}} = \text{constant}$ [46, 47]. Therefore, the local temperature in each wordline is $\beta^{-1} = (2\pi\xi)^{-1}$.

Now we want to understand how the detector behaves if one substitutes the Rindler Fock space by another Hilbert space that carry another representation of the field algebra. This is the fundamental point here. One is applying operators constructed to act in one representation of the operator algebra to states that belong to an unitarily inequivalent representation. In other words, we wish to comprehend how the uniformly accelerated Glauber detector behaves if the state of field is the Minkowski vacuum. In this case, from the above discussion, we have to study the contribution from $F_1(\omega_{jb}, \tau_f, \tau_i)$. By means of the same above procedure, one obtains an expression to the transition rate similar to Eq. (38), where now $|\phi_i\rangle = |0, M\rangle$. Let us call it $W_1(\xi; |0, M\rangle)$. Now using Eq. (22), one obtains

$$W_1(\xi; |0, M\rangle) = \frac{1}{4\pi^4} \int d^2q \int_0^\infty d\nu \frac{\sinh \pi\nu}{e^{\beta\nu} - 1} K_{\nu\nu}^2 \left(\xi \sqrt{m_0^2 + q^2}\right). \quad (40)$$

We obtain that $W_1(\xi; |0, M\rangle) = W_\beta(\xi)$. This is another version of the Bisognano-Wichmann theorem [48, 49] that states that the Minkowski vacuum expectation value of observables which are localized in the right Rindler wedge satisfies the Kubo-Martin-Schwinger condition [50, 51] with respect to the Rindler time variable $\eta$. As has been discussed in the literature, for the accelerated detector interacting with the Minkowski vacuum state, this state is equivalent to a density matrix since one has to trace all of degrees of freedom in the Rindler left edge. From the point of view of the accelerated observer, the contribution to the rate given by $W_1$ is a process of excitation of the Glauber detector with absorption of a Rindler quantum from the scalar field. In addition, for the case of the Unruh-DeWitt accelerated detector we have to properly take into account the contribution $\tilde{W}_2$ which comes from $\tilde{F}_2(\omega_{jb}, \tau_f, \tau_i)$. Note that this contribution in the Rindler vacuum vanishes.

From Eq. (40), for the case where $m_0^2 = 0$ we are able to obtain the well-known result

$$W_1(\xi; |0, M\rangle) = \frac{1}{(2\pi\xi)^2} \int_0^\infty d\nu \frac{\nu}{e^{\beta\nu} - 1}. \quad (41)$$
In the worldline \( \xi = \text{constant} \) the rate of absorption of energy in this square-law detector per unit of frequency is

\[
\Pi(\nu) = \frac{1}{(2\pi \xi)^2 e^{2\pi \nu} - 1}.
\] (42)

Now we are able to discuss the interpretation given by the inertial observer for the excitation process. Since for the inertial observer the state of the field is the Minkowski vacuum, the detector makes transitions to excited states and quanta of the field show up. The source of energy that allows for these processes comes from the agent that accelerates the detector. Applying Eq. (19) to the Minkowski vacuum state, we get

\[
b(\nu, \mathbf{q}) |0, M\rangle = \int d^3\mathbf{k} V(\nu, \mathbf{q}, \mathbf{k}) a(\mathbf{k}) |0, M\rangle.
\] (43)

The traditional way to interpret the above equation is to assert that counter-rotating processes \( \sigma^+ a(\mathbf{k}) (\sigma^+ = |e\rangle \langle g| \) for usual two-level detectors) are responsible for the right-hand side contribution. These represent a virtual process. For a detector at rest in a rectilinear uniformly accelerated moving frame, this process becomes a real one, which is commonly deemed as Unruh radiation. The same analysis performed by an accelerated observer discussing the role of \( F_1 \) can be realized by the inertial observer. The \( F_1 \) contribution is a processes of excitation of the detector with emission. There appears a field state as a continuous superposition of number states.

On the other hand, in order to understand why inertial observers should perceive the existence of radiation emitted by a uniformly accelerated detector, one may also resort to the familiar result derived from the interaction of a quantized electromagnetic field with a classical source [52]. Indeed, in this case one verifies that classical current creates a coherent state from the vacuum. The probability corresponding to the emission of photons is given by a Poisson distribution. Even though such physical situations should be plainly distinguished, one may argue that in both cases the physical interpretation according to the inertial-observer perspective should follow along similar lines.

5 The Glauber detector in Kalnis-Miller coordinate system

The aim of this section is to discuss the behaviour of a detector in a frame that in the remote past is inertial and in the far future becomes uniformly accelerated. A coordinate system adapted to this frame was obtained in Refs. [53, 54, 55]. The scalar field quantization was discussed by Costa, Svaiter and De Paola [56, 57, 58, 59, 60]. For the four-dimensional spacetime with the usual cartesian coordinates \( x^\mu = (t, x, y) \) and the curvilinear coordinates \( X^\mu = (\xi, \eta, y) \), we define the following mapping

\[
t + x = \frac{2}{a} \sinh a(\xi + \eta),
\] (44)

and

\[
t - x = -\frac{1}{a} e^{-a(\xi - \eta)},
\] (45)
for \(-\infty < \eta < \infty\) and \(-\infty < \xi < \infty\). This coordinate system \(X^\mu = (\xi, \eta, y)\) is valid only for \(t - x < 0\). Nevertheless, it is possible to extend it in order to cover the whole spacetime. The four-dimensional line element can be written using this curvilinear coordinates. We get
\[
ds^2 = (e^{-2a\eta} + e^{2a\xi}) (d\eta^2 - d\xi^2) - dy^2 - dz^2.
\]
The next step is to discuss the proper acceleration of an observer travelling in the worldline \(\xi, y = \text{constant}\). We have that
\[
\alpha = a \left( e^{-2a\eta} + e^{2a\xi} \right)^{-3/2} e^{2a\xi} |_{\xi=\xi,y=y}.
\]
Therefore, in the remote past and in the far future we have, respectively
\[
\left\{
\begin{align*}
\lim_{\eta \to -\infty} \alpha(\eta, \xi, y)|_{\xi=\xi,y=y} &= 0, \\
\lim_{\eta \to \infty} \alpha(\eta, \xi, y)|_{\xi=\xi,y=y} &= ae^{-a\xi} = a_\infty.
\end{align*}
\right.
\]
In these curvilinear coordinates the Klein-Gordon can be written as
\[
\left[ \frac{\partial^2}{\partial \eta^2} - \frac{\partial^2}{\partial \xi^2} + m_0^2 + m^2 \right] F(\eta, \xi, y) = 0.
\]
Let us define the variables \(\zeta = a^{-1}e^{-a\eta}\) and \(\chi = a^{-1}e^{a\xi}\) for \(\infty > \zeta > 0\) and \(0 < \chi < \infty\). Using the result obtained by Kalnias and Miller we write \(\phi(\zeta, \chi, y) = F(\zeta)G(\chi)H(y)\). The separation of variables of the above equations yields
\[
\left[ \frac{d^2}{d\zeta^2} + \frac{1}{\zeta} \frac{d}{d\zeta} + m_0^2 + \frac{\lambda^2}{\zeta^2} \right] F(\zeta) = 0,
\]
and
\[
\left[ \frac{d^2}{d\chi^2} + \frac{1}{\chi} \frac{d}{d\chi} - m_0^2 - \frac{\lambda^2}{\chi^2} \right] G(\chi) = 0.
\]
There are two complete orthonormal bases that can be used to expand the scalar field, \(\{u_\lambda(\zeta, \chi, y), u^*_\lambda(\zeta, \chi, y)\}\) and \(\{v_\nu(\zeta, \chi, y), v^*_\nu(\zeta, \chi, y)\}\) which are of the form
\[
u_\nu(\zeta, \chi, y) = \frac{1}{4\pi^{3/2}} \left[ \lambda (1 - e^{-2\pi \lambda}) \right]^{1/2},
\]
and
\[
u^*_\nu(\zeta, \chi, y) = \frac{1}{4\pi^{3/2}} \left[ J_{\lambda}(m\zeta)K_{\lambda}(m\chi)e^{i\pi y},
\right.
\]
where \(J_\nu(z)\), \(J^{(2)}_\nu(z)\) are Bessel functions of third kind or Hankel’s function and \(J_\nu(z)\) is a Bessel function of first kind. One can show that \(u_\lambda(\zeta, \chi, y)\) are \(u^*_\lambda(\zeta, \chi, y)\) are positive and
negative frequency modes in the remote past and \( \nu_\nu(\zeta, \chi, y) \) are respectively positive and negative frequency modes in the far future \cite{61, 62}. In the remote past, the field can be expanded as

\[
\phi(\zeta, \chi, y) = \int_0^\infty d\lambda \int d^2q \left[ a_{in}(\lambda, q)u_\lambda(\zeta, \chi, y) + a_{in}^\dagger(\lambda, q)u_\lambda^*(\zeta, \chi, y) \right],
\] (57)

where \( a_{in}(\lambda, q) \) and \( a_{in}^\dagger(\lambda, q) \) are annihilation and creation operator for quanta of the field in the remote past. It is now possible to define a vacuum state by

\[
a_{in}(\lambda, q) |0, in\rangle = 0, \quad \forall \lambda, q.
\] (58)

In the same way, the field expansion suitable for observers in the far future can be written as

\[
\phi(\zeta, \chi, y) = \int_0^\infty d\nu \int d^2q \left[ a_{out}(\nu, q)\nu_\nu(\zeta, \chi, y) + a_{out}^\dagger(\nu, q)\nu_\nu^*(\zeta, \chi, y) \right],
\] (59)

where \( a_{out}(\nu, q) \) and \( a_{out}^\dagger(\nu, q) \) are annihilation and creation operator for quanta of the field in the far future. The vacuum state in this case is defined by

\[
a_{out}(\nu, q) |0, out\rangle = 0, \quad \forall \nu, q.
\] (60)

The number of quanta associated with the \((\nu, q)-out\) modes in the vacuum defined in the remote past is given by

\[
\langle 0, in|a_{out}^\dagger(\nu, q)a_{out}(\nu, q)|0, in\rangle = \int d\lambda \int d^2q' |\beta_{vq,\lambda q'}|^2,
\] (61)

where the Bogoliubov coefficients are given by

\[
|\beta_{vq,\lambda q'}| = \frac{1}{a\sqrt{e^{2\pi \nu} - 1}} \delta(\nu - \lambda)\delta(q - q').
\] (62)

This model presents a behaviour similar to the Hawking effect. If one considers the quantum version of a field theory without interactions the vacuum state constructed in the infinity past appears as a thermal state in the far future. Let us discuss the behaviour of a Glauber detector in this model. We define the quantity

\[
W(x; \Phi) = \langle \Phi | :\varphi(-)^{(x)}\varphi^{(+)}(x) :_{out} | \Phi \rangle.
\] (63)

As discussed above, for a broadband detector this is proportional to the excitation rate. For \(|\Phi\rangle = |0, out\rangle\) we get \(W = 0\) as expected. We are interested in considering the case

\[
W(x; |0, in\rangle) = \langle 0, in| :\varphi(-)^{(x)}\varphi^{(+)}(x) :_{out} |0, in\rangle.
\] (64)

A straightforward calculation gives us

\[
W(\zeta, \chi; |0, in\rangle) = \frac{1}{4\pi^3} \int d^2q \int_0^\infty d\nu \nu J_{\nu}(m\zeta)J_{-\nu}(m\zeta)K_{\nu}^2(m\chi).
\] (65)

For the case where \(m_0^2 = 0\) and \(\zeta \to 0\), we get

\[
\lim_{\zeta \to 0} W(\zeta, \chi; |0, in\rangle) = \frac{1}{(2\pi \chi)^2} \int_0^\infty d\nu \frac{\nu}{e^{2\pi \nu} - 1}.
\] (66)
Hence the same interpretation given to $W_1$ and $W_\beta$ can be discussed in the Kalnis-Miller model. Previously we have obtained that the transition probability rate of an accelerated detector interacting with a thermal Rindler bath is equal to the transition probability rate of the accelerated detector coupled to the field in the Minkowski vacuum. In a similar fashion, we observe that, for the massless field, the transition probability rate of the detector in the far future is equivalent to the transition probability rate of the detector at rest in a non-inertial reference frame interacting with the field prepared in an usual thermal state.

One can ask how to generate the Rindler vacuum state. It is known that accelerated mirrors can be used for such a purpose [63]. The Kalnis-Miller model can be employed to investigate the accelerated mirror problem imposing a Dirichlet boundary conditions in one worldline $\chi = \text{constant}$ in a two dimensional toy model. See, for example, the approach presented in Ref. [64]. A natural continuation is to discuss the flux of radiation from this mirror.

6 Conclusions

Field quantum theory is a conceptual formalism that combines the probabilistic interpretation of quantum mechanics with the tenets of the special relativity. One fundamental problem that emerges is how to link the formalism of the coupled fields with measurable observables. In particle physics, when studying scattering processes we assume that in the remote past and far future, particles are non interacting. In this asymptotic region, one may replace the fields by free fields describing free particles carrying mass and other quantum numbers. Another useful approach is to introduce a detector of field quanta.

In this paper, we propose to study the Unruh-Davies effect using the Glauber theory of photodetection. There are some simplifications in our discussion such as the use of point-like interactions (fields are operator-valued distributions) and the absence of a switching function. We believe that it is possible to understand the effect even with such simplifications. We discussed the process of measuring Rindler field quanta by an uniformly accelerated detector interacting with a massive scalar field. Using the Glauber theory of photodetection, the rate of excitation of an accelerated detector interacting with the scalar field was evaluated. We rederive the result that a uniformly accelerated detector coupled to a field in the Minkowski vacuum state has the same transition rate of an accelerated detector interacting with the field in a Rindler thermal one at some temperature $\beta^{-1}$.

Our results show that an inertial interpretation of the radiative processes associated with observers at rest in the frame of the uniformly accelerated detector is available in the Glauber formalism. In other words, we showed how the inertial observer interprets the excitation of the accelerated detector. The fundamental difference between the Unruh-DeWitt detector and the Glauber detector is the contribution coming from $F_2$ (or $\bar{F}_2$). If one wishes to keep this term in excitation processes, then one is inescapably led to the following interpretation [65]. The detector is able to perform a transition to excited states since $\omega_{j\eta} > 0$, but with emission of an unphysical quantum with negative energy $-\nu$. In this situation we have energy conservation and such a term in the Glauber scenario also contributes to decrease the number of excitations in the environment. In any case we have given above detailed arguments that showed that such a term does not actually contribute to the excitation rate for large observation time intervals.

Also using the Glauber theory of photodetection we have discussed a model with unitary non-equivalent vacuum states in the remote past in the far future. This detector presents a behaviour similar to the detector interacting with a field in a scenario of gravitational collapse.
For the massless case, we obtained that the transition probability rate of the detector in the far future is equivalent to the transition probability rate of the detector at rest in a non-inertial reference frame interacting with the field prepared in an usual thermal state. In this scenario, the introduction of a boundary condition allows one to investigate how to generate the Fulling vacuum.

An interesting application of our analysis is the discussion of the equivalence principle in the quantum domain [66, 67]. One must show that the following physical situations are equivalent: The excitation rate of a detector at rest in a local inertial frame coupled to a field in the Boulware vacuum must be equal to the excitation rate of a detector travelling in an inertial world line in a Minkowski spacetime interacting with a field in the Fulling vacuum state [68]. Since Rindler operators span only a sub-algebra of the equal-time field algebra, it should be clear that it is not straightforward to obtain the desired result. We hope to fully explore this issue in a future work.

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