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Accelerated Multi-Level Atoms in an Electromagnetic Vacuum and Fulling–Davies–Unruh Effect *

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We consider, from the point of view of a coaccelerated frame, a uniformly accelerated multi-level atom in interaction with vacuum quantum electromagnetic fields in the multi-polar coupling scheme, and calculate the rate of change of the atom's energy assuming a thermal bath at a finite temperature T in the Rindler wedge. Comparison with the spontaneous excitation rate of the atom calculated in the instantaneous inertial frame of the atom shows that both the inertial and coaccelerated observer would agree with each other only when the temperature of the thermal bath equals the FDU value $T_{\rm FDU} = a/2\pi$.

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In a flat spacetime, it is well known that a uniformly accelerated (Rindler) observer in the Minkowski vacuum of quantum fields will feel as if immersed in a thermal bath of Rindler particles at a temperature proportional to the observer's proper acceleration (see[1,2] for recent reviews). This result was first discussed in detail by Fulling,[3] Davies[4] and Unruh,[5] called the Fulling–Davies–Unruh (FDU) effect. The effect is associated with the existence of an event horizon for the accelerated observer, and closely related to the Hawking radiation from black holes.

So far, the FDU effect has been examined in terms of different physical processes.[2] For the weak decay of uniformly accelerated protons, the same lifetime can be obtained in the inertial frame and the coaccelerated frame considering the FDU effect.[6] For the bremsstrahlung effect associated with a uniformly accelerated point electric charge and a point scalar source in the Minkowski vacuum, the same response rate in both the inertial and coaccelerated frame can be obtained only assuming the FDU thermal bath with an infinite number of zero-energy Rindler photons in the Rindler wedge,[7−9] and the same conclusion can also be obtained with some boundary conditions imposed on the quantum field.[10] Recently, for a two-level atom interacting with a massless quantum scalar field both in an unbounded space and in a space with a reflecting plane boundary,[11,12] we have demonstrated that the same spontaneous excitation rate for a uniformly accelerated atom in the Minkowski vacuum obtained in the inertial frame can be recovered in the coaccelerated frame assuming a thermal bath at the Fulling–Davies–Unruh temperature $T_{\rm FDU} = a/2\pi$.[13] All these studies show that the existence of the FDU effect is mandatory for the consistency of quantum field theory.

As a further step along the line of Ref. [13] for a better understanding of the FDU effect, we consider, in this study, from the point of view of a coaccelerated frame, the spontaneous excitation of a uniformly accelerated multi-level atom in interaction with vacuum quantum electromagnetic fields in the multi-polar coupling scheme. It should be noted that a multi-level atom interacting with quantized electromagnetic fields is a much more realistic model than a two-level atom in interaction with scalar fields considered in Ref. [13,14]. Assuming that there exists a thermal bath of infinite number of Rindler particles at a finite temperature T (Rindler thermal bath), we separately calculate, in the coaccelerated frame, the contributions of thermal fluctuations and radiation reaction to the rate of change of the atom's energy with the quantization of electromagnetic fields in the Rindler wedge. Comparison with the spontaneous excitation rate obtained by an instantaneous inertial observer[15] shows that these two observers can reach an agreement only when the temperature equals the FDU value $T_{\rm FDU} = a/2\pi$.

We consider a multilevel atom interacting with quantum electromagnetic fields, and assume that the atom has a constant proper acceleration a along the x direction in the Minkowski coordinates. In the Rindler wedge, the line element[16] is described by

$$ds^2 = e^{2a\xi}(d\tau^2 - d\xi^2) - dy^2 - dz^2. \qquad (1)$$

The Rindler coordinates are related to the usual Minkowski coordinates by means of the transformation

$$t = (e^{a\xi}/a)\sinh a\tau, \quad x = (e^{a\xi}/a)\cosh a\tau. \qquad (2)$$

The world line with constant ξ, y, z describes an observer who is accelerated uniformly with the proper acceleration $ae^{-a\xi}$. Thus in the coaccelerated frame

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our accelerated atom is stationary with the trajectory $x(\tau) = [\tau, \xi(\tau) = y(\tau) = z(\tau) = 0]$. The atom is assumed to have a series of stationary states $|n\rangle$, with energies ω_n. The Hamiltonian that governs the time evolution of the atom with respect to the proper time τ is written as

$$H_A(\tau) = \sum_n \omega_n \sigma_{nn}(\tau), \tag{3}$$

where $\sigma_{nn}(\tau) = |n\rangle\langle n|$. The free Hamiltonian of the quantum electromagnetic field is

$$H_F(\tau) = \int d\omega \int dk_y \int dk_z \sum_{\lambda=1}^{4} \omega b^{\dagger}_{(\lambda,\omega,k_y,k_z)}(\tau) \cdot b_{(\lambda,\omega,k_y,k_z)}(\tau), \tag{4}$$

where λ labels the polarization degrees of freedom; $b^{\dagger}_{(\lambda,\omega,k_y,k_z)}(\tau)$ and $b_{(\lambda,\omega,k_y,k_z)}(\tau)$ denote the creation and annihilation operators of the Rindler particle. In the Feynman gauge, the quantization of the electromagnetic field in the Rindler wedge can be carried out by the expansion in terms of the creation and annihilation operators as

$$A_\mu(x(\tau)) = \int d\omega \int dk_y \int dk_z \sum_{\lambda=1}^{4} \big(b_{(\lambda,\omega,k_y,k_z)}(\tau) \cdot A^{(\lambda,\omega,k_y,k_z)}_{\mu}(x(\tau)) + \text{H.c.}\big), \tag{5}$$

which satisfies the field equation

$$\nabla_\mu \nabla^\mu A_\nu = 0, \tag{6}$$

and the Lorenz condition

$$\nabla_\mu A^\mu = 0. \tag{7}$$

We couple the multi-level atom and the quantum electromagnetic field in the multipolar coupling scheme[17]

$$\begin{aligned} H_I(\tau) =& -e\boldsymbol{r}(\tau) \cdot \boldsymbol{E}(x(\tau)) \\ =& -e \sum_{mn} \boldsymbol{r}_{mn} \cdot \boldsymbol{E}(x(\tau)) \sigma_{mn}(\tau), \end{aligned} \tag{8}$$

where e is the electron electric charge, $e\boldsymbol{r}$ the atomic electric dipole moment. Now we can obtain the Heisenberg equations of motion for the dynamical variables of the atom and the electromagnetic field from the Hamiltonian $H = H_A + H_F + H_I$,

$$\frac{d}{d\tau}\sigma_{mn}(\tau) = i(\omega_m - \omega_n)\sigma_{mn}(\tau) - ie\boldsymbol{E}(x(\tau)) \cdot [\boldsymbol{r}(\tau), \sigma_{mn}(\tau)], \tag{9}$$

$$\frac{d}{d\tau} b_{(\lambda,\omega,k_y,k_z)}(\tau) = -i\omega b_{(\lambda,\omega,k_y,k_z)}(\tau) - ie\boldsymbol{r}(\tau) \cdot [\boldsymbol{E}(x(\tau)), b_{(\lambda,\omega,k_y,k_z)}(\tau)]. \tag{10}$$

The solutions of the equations of motion can be split into two parts: a free part, which is present even in the absence of the coupling, and a source part, which is caused by the interaction of the atom and the field.

We assume that the initial state of the electromagnetic field is not the Rindler vacuum, but a thermal state $|\beta\rangle$ at a finite temperature T and the density matrix $\rho = e^{-\beta H_F} = e^{-H_F/T}$, while the atom is prepared in the state $|a\rangle$. We will separately discuss the two physical mechanisms that contribute to the rate of change of atomic observables: the contribution of thermal fluctuations and that of radiation reaction. For this purpose, we follow the method of Dalibard, Dupont–Roc and Cohen–Tannoudji (DDC)[18,19], and choose a symmetric ordering between the atom and field variables. We can separate the two contributions of thermal fluctuations and radiation reaction to the rate of change of H_A as,

$$\frac{dH_A(\tau)}{d\tau} = \Big(\frac{dH_A(\tau)}{d\tau}\Big)_{TF} + \Big(\frac{dH_A(\tau)}{d\tau}\Big)_{RR}, \tag{11}$$

where

$$\Big(\frac{dH_A(\tau)}{d\tau}\Big)_{TF} = -\frac{ie}{2}\big(\boldsymbol{E}^f(x(\tau)) \cdot [\boldsymbol{r}(\tau), H_A(\tau)\,] + [\boldsymbol{r}(\tau), H_A(\tau)(\tau)] \cdot \boldsymbol{E}^f(x(\tau))\big), \tag{12}$$

representing the contribution of the thermal fluctuations and

$$\Big(\frac{dH_A(\tau)}{d\tau}\Big)_{RR} = -\frac{ie}{2}\big(\boldsymbol{E}^s(x(\tau)) \cdot [\boldsymbol{r}(\tau), H_A(\tau)] + [\boldsymbol{r}(\tau), H_A(\tau)] \cdot \boldsymbol{E}^s(x(\tau))\big), \tag{13}$$

denoting that of the radiation reaction. Separating $r_i(\tau)$ and $\sigma_{nn}(\tau)$ into their free part and source part and taking the expectation value in the field and atom states, we can obtain, in a perturbative treatment up to order e^2,

$$\Big\langle \frac{dH_A(\tau)}{d\tau} \Big\rangle^R_{TF} = 2ie^2 \int_{\tau_0}^{\tau} d\tau' (C^F_{ij})_\beta(x(\tau), x(\tau')) \frac{d}{d\tau}(\chi^A_{ij})_a(\tau, \tau'), \tag{14}$$

$$\Big\langle \frac{dH_A(\tau)}{d\tau} \Big\rangle^R_{RR} = 2ie^2 \int_{\tau_0}^{\tau} d\tau' (\chi^F_{ij})_\beta(x(\tau), x(\tau')) \frac{d}{d\tau}(C^A_{ij})_a(\tau, \tau'), \tag{15}$$

where the superscript R denotes the calculation performed in the Rindler coordinates and $|\rangle = |a, \beta\rangle$. Here the statistical functions, $(C^F_{ij})_\beta$ and $(\chi^F_{ij})_\beta$, of the field at a finite temperature $T = 1/\beta$ are defined as

$$(C^F_{ij})_\beta(x(\tau), x(\tau')) = \frac{1}{2}\langle\{E^f_i(x(\tau)), E^f_j(x(\tau'))\}\rangle_\beta, \tag{16}$$

$$(\chi^F_{ij})_\beta(x(\tau), x(\tau')) = \frac{1}{2}\langle[E^f_i(x(\tau)), E^f_j(x(\tau'))]\rangle_\beta. \tag{17}$$

$(C^F_{ij})_\beta$ and $(\chi^F_{ij})_\beta$ are called the symmetric correlation function and linear susceptibility of the field in the

thermal state $|\beta\rangle$. Analogously the symmetric correlation function and linear susceptibility of the atom are defined as

$$(C^A_{ij})_a(\tau,\tau') = \frac{1}{2}\langle a|\{r^f_i(\tau), r^f_j(\tau')\}|a\rangle, \tag{18}$$

$$(\chi^A_{ij})_a(\tau,\tau') = \frac{1}{2}\langle a|[r^f_i(\tau), r^f_j(\tau')]|a\rangle, \tag{19}$$

which are independent of the trajectory of the atom and characterize only the atom itself. The explicit forms of the statistical functions of the atom are given by

$$(C^A_{ij})_a(\tau,\tau') = \frac{1}{2}\sum_b[\langle a|r_i(0)|b\rangle\langle b|r_j(0)|a\rangle e^{i\omega_{ab}(\tau-\tau')} + \langle a|r_j(0)|b\rangle\langle b|r_i(0)|a\rangle e^{-i\omega_{ab}(\tau-\tau')}], \tag{20}$$

$$(\chi^A_{ij})_a(\tau,\tau') = \frac{1}{2}\sum_b[\langle a|r_i(0)|b\rangle\langle b|r_j(0)|a\rangle e^{i\omega_{ab}(\tau-\tau')} - \langle a|r_j(0)|b\rangle\langle b|r_i(0)|a\rangle e^{-i\omega_{ab}(\tau-\tau')}], \tag{21}$$

where $\omega_{ab} = \omega_a - \omega_b$ and the sum extends over a complete set of atomic states.

In order to calculate the statistical functions for the electromagnetic field, firstly we need to seek a set of normal modes solutions of the field equation (6) which satisfies the Lorenz condition (6). The positive frequency modes in Eq. (5) can be written as

$$A^{(\lambda,\omega,k_y,k_z)}_\mu(x(\tau)) = f^{(\lambda,\omega,k_y,k_z)}_\mu(\xi)e^{i(k_y y + k_z z - \omega\tau)}. \tag{22}$$

According to Refs. [7,8,20], there are two physical modes, which are not pure gauge and satisfy the Lorenz condition,

$$f^{(I,\omega,k_y,k_z)}_\mu = C^{(I,\omega,k_y,k_z)}(0,0,k_z\phi,-k_y\phi), \tag{23}$$

and

$$f^{(II,\omega,k_y,k_z)}_\mu = C^{(II,\omega,k_y,k_z)}(\partial_\xi\phi,-i\omega\phi,0,0), \tag{24}$$

where $f_\mu = (f_\tau, f_\xi, f_y, f_z)$, and $C^{(\lambda,\omega,k_y,k_z)}$ are normalization constants. The function ϕ has the form

$$\phi = K_{i\omega/a}\Big(\frac{k_\perp}{a}e^{a\xi}\Big), \tag{25}$$

where $K_\nu(x)$ is the Bessel function of imaginary argument and $k_\perp = \sqrt{k_y^2 + k_z^2}$. Imposing the following normalization condition:

$$(A^{(\lambda,\omega,k_y,k_z)}, A^{(\lambda',\omega',k'_y,k'_z)}) = \delta^{\lambda\lambda'}\delta(\omega-\omega')\delta(k_y-k'_y)\delta(k_z-k'_z), \tag{26}$$

we can obtain

$$|C^{(I,\omega,k_y,k_z)}| = |C^{(II,\omega,k_y,k_z)}| = \frac{1}{2\pi^2 k_\perp}\Big(\frac{\sinh(\pi\omega/a)}{a}\Big)^{1/2}. \tag{27}$$

Then, using Eq. (5), we can calculate the two point function for the electromagnetic field at a finite temperature $T = 1/\beta$ as follows:

$$\begin{aligned}&\langle A_\mu(x(\tau))A_\nu(x(\tau'))\rangle_\beta \\ &= \mathrm{tr}[\rho A_\mu(x(\tau))A_\nu(x(\tau'))]/\mathrm{tr}[\rho] \\ &= \int_0^\infty d\omega\int_{-\infty}^\infty dk_y\int_{-\infty}^\infty dk_z\sum_\lambda\Big[\frac{e^{\omega/T}}{e^{\omega/T}-1} \cdot A^{(\lambda,\omega,k_y,k_z)}_\mu(x(\tau))A^{*(\lambda,\omega,k_y,k_z)}_\nu(x(\tau')) \\ &\quad + \frac{1}{e^{\omega/T}-1}A^{*(\lambda,\omega,k_y,k_z)}_\mu(x(\tau)) \cdot A^{(\lambda,\omega,k_y,k_z)}_\nu(x(\tau'))\Big].\end{aligned} \tag{28}$$

The statistical functions of the field can be calculated using (28).

Next, we apply the previously developed formalism to calculate the rate of change of the mean atomic energy for the atom in the coaccelerated frame (the Rindler wedge with a proper acceleration a) at a finite temperature $T = 1/\beta$. The atom's trajectory can be described, in Rindler coordinates, by

$$\tau\ , \xi(\tau) = y(\tau) = z(\tau) = 0. \tag{29}$$

Putting the trajectory (29) into Eqs. (22)–(24) and (28), and using the equation $E_i = \partial_i A_0 - \partial_0 A_i$, we can calculate the field correlation function to get

$$\langle E_i(x(\tau))E_j(x(\tau'))\rangle_\beta = \delta_{ij}\frac{1}{6\pi^2}\int_0^\infty d\omega\omega(a^2+\omega^2) \cdot\Big(\frac{e^{\omega/T}}{e^{\omega/T}-1}e^{-i\omega(\tau-\tau')} + \frac{1}{e^{\omega/T}-1}e^{i\omega(\tau-\tau')}\Big). \tag{30}$$

From Eq. (30), the statistical functions of the field, (16) and (17), can be shown to have the form

$$(C^F_{ij})_\beta(x(\tau),x(\tau')) = \delta_{ij}\frac{1}{12\pi^2}\int_0^\infty d\omega\omega(a^2+\omega^2) \cdot\Big(1+\frac{2}{e^{\omega/T}-1}\Big)\big(e^{-i\omega(\tau-\tau')} + e^{i\omega(\tau-\tau')}\big), \tag{31}$$

$$(\chi^F_{ij})_\beta(x(\tau),x(\tau')) = \delta_{ij}\frac{1}{12\pi^2}\int_0^\infty d\omega\ \omega\ (a^2+\omega^2) \cdot\big(e^{-i\omega(\tau-\tau')} - e^{i\omega(\tau-\tau')}\big). \tag{32}$$

With a substitution $u = \tau - \tau'$, we can now evaluate the contributions of thermal fluctuation (14) and radiation reaction (15) to the rate of change of the atom's energy

$$\Big\langle\frac{dH_A(\tau)}{d\tau}\Big\rangle^R_{TF} = -\frac{e^2}{12\pi^2}\sum_b\omega_{ab}|\langle a|\boldsymbol{r}(0)|b\rangle|^2 \cdot\int_0^\infty d\omega\ \omega(a^2+\omega^2)\Big(1+\frac{2}{e^{\omega/T}-1}\Big) \cdot\int_0^\infty du(e^{-i\omega u}+e^{i\omega u})(e^{-i\omega_{ab}u}+e^{i\omega_{ab}u}), \tag{33}$$

$$\Big\langle\frac{dH_A(\tau)}{d\tau}\Big\rangle^R_{RR}=-\frac{e^2}{12\pi^2}\sum_b\omega_{ab}|\langle a|\boldsymbol{r}(0)|b\rangle|^2\cdot\int_0^\infty d\omega\;\omega(a^2+\omega^2)\times\int_0^\infty du\cdot(e^{-i\omega u}-e^{i\omega u})(e^{-i\omega_{ab}u}-e^{i\omega_{ab}u}),\qquad(34)$$

where we have extended the range of integration to infinity for sufficiently long times $\tau-\tau_0$. It is easy to evaluate the integral to find

$$\Big\langle\frac{dH_A(\tau)}{d\tau}\Big\rangle^R_{TF}=-\frac{e^2}{6\pi}\Big[\sum_{\omega_a>\omega_b}\omega_{ab}^4|\langle a|\boldsymbol{r}(0)|b\rangle|^2\cdot\Big(\frac{a^2}{\omega_{ab}^2}+1\Big)\Big(1+\frac{2}{e^{\omega_{ab}/T}-1}\Big)-\sum_{\omega_a<\omega_b}\omega_{ab}^4\cdot|\langle a|\boldsymbol{r}(0)|b\rangle|^2\Big(\frac{a^2}{\omega_{ab}^2}+1\Big)\Big(1+\frac{2}{e^{|\omega_{ab}|/T}-1}\Big)\Big],\qquad(35)$$

for the contribution of the thermal fluctuations to the rate of change of atomic excitation energy, and

$$\Big\langle\frac{dH_A(\tau)}{d\tau}\Big\rangle^R_{RR}=-\frac{e^2}{6\pi}\Big[\sum_{\omega_a>\omega_b}\omega_{ab}^4|\langle a|\boldsymbol{r}(0)|b\rangle|^2\Big(\frac{a^2}{\omega_{ab}^2}+1\Big)+\sum_{\omega_a<\omega_b}\omega_{ab}^4|\langle a|\boldsymbol{r}(0)|b\rangle|^2\Big(\frac{a^2}{\omega_{ab}^2}+1\Big)\Big].\qquad(36)$$

for that of radiation reaction. From the point of view of a coaccelerated frame, we note that here the contribution of radiation reaction to the rate of change of the atomic energy, Eq. (36), has the same form as that obtained by an instantaneous inertial observer in the Minkowski vacuum[15], and is independent of the temperature, and always leads to a loss of atom's energy. While the contribution of thermal fluctuations depends on the temperature of the Rindler thermal bath, and equally renders the atom in the ground state to excite and the atom in the excited state to de-excite. Adding the contributions of thermal fluctuations and radiation reaction, we obtain the total rate of change of the atomic energy,

$$\Big\langle\frac{dH_A(\tau)}{d\tau}\Big\rangle^R_{\rm tot}=-\frac{e^2}{3\pi}\Big(\sum_{\omega_a>\omega_b}\omega_{ab}^4|\langle a|\boldsymbol{r}(0)|b\rangle|^2\cdot\Big(\frac{a^2}{\omega_{ab}^2}+1\Big)\Big(1+\frac{1}{e^{\omega_{bd}/T}-1}\Big)-\sum_{\omega_a>\omega_b}\omega_{ab}^4\cdot|\langle a|\boldsymbol{r}(0)|b\rangle|^2\Big(\frac{a^2}{\omega_{ab}^2}+1\Big)\frac{1}{e^{|\omega_{bd}|/T}-1}\Big).\qquad(37)$$

Since the rate of change of the atomic energy in the proper time must be a scalar, we should expect an agreement in both the inertial and the coaccelerated frames. Comparing Eq. (37) with the spontaneous excitation rate of the uniformly accelerated atom calculated in the inertial frame[15], one sees that only when the temperature of the Rindler thermal bath equals the FDU value, $T_{\rm FDU}=a/2\pi$, can the inertial and the coaccelerated observers reach an agreement. Hence the FDU effect is necessary for the consistency of quantum field theory.

An interesting feature worth to be noted of the above results, in contrast to the case where the accelerated atom is assumed to be in interaction with massless scalar fields in a more or less toy model,[13,14] is the appearance of the correction terms proportional to the proper acceleration squared, which are not in the Planckian form. These corrections, which deviate from a purely thermal form, lead the accelerated atom to radiate and absorb differently from an inertial atom immersed in the thermal bath. However, this discrepancy does not imply that the exact final thermal equilibrium is not achieved. As a matter of fact, since the correction term is an even function of ω_{ab}, one can show, by the same argument as that of Ref. [21], that thermal equilibrium will be established exactly at the FDU temperature $T_{\rm FDU}=a/2\pi$.

In conclusion, assuming that there exists a thermal bath of infinite number of Rindler particles at a finite temperature T in the Rindler wedge, and we have separately calculated, from the point of view of the coaccelerated frame, the contributions of thermal fluctuations and radiation reaction to the rate of change of the atomic energy for a uniformly accelerated multilevel atom interacting with vacuum electromagnetic fields in the multi-polar coupling scheme using DDC's method. It is shown that for a uniformly accelerated atom in the Minkowski vacuum, both the instantaneous inertial and the coaccelerated observers can obtain the same rate of change of the atom's energy only when the temperature of the Rindler thermal bath equals the FDU value. So, for a coaccelerated observer, the Minkowski vacuum is not vacant, but is equivalent to a thermal bath of Rindler particles at the FDU temperature $T_{\rm FDU}=a/2\pi$. This is the FDU effect.

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