Background thermal contributions in testing the Unruh effect

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Abstract

We consider inertial and accelerated Unruh-DeWitt detectors moving in a background thermal bath and calculate their excitation rate. It is shown that for fast moving detectors such a thermal bath does not affect substantially the excitation probability. Our results are discussed in connection with a possible proposal of testing the Unruh effect in high energy particle accelerators.

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I. INTRODUCTION

It is already two decades since Hawking discovered the striking result that quantum mechanics may induce black holes to evaporate \([1]\). Many different questions related with such an effect have been clarified since that time, and a number of other ones are presently under investigation \([2]\). The so-called Unruh effect \([3,4]\) has played an outstanding role in emphasizing some of the exquisite features present in the Hawking effect. According to the Unruh effect, a detector uniformly accelerated through the *inertial* vacuum responds as being in a thermal bath characterized by a temperature proportional to its proper acceleration. We call *inertial* vacuum the no-particle state as described by a family of observers following an *inertial* timelike Killing field \(\partial_t\) in Minkowski space. The Unruh effect is a direct consequence of the fact that the particle content of a field theory is very frame dependent \([3,5]\). Recently, Bell and Leinaas \([6]\) raised the very interesting possibility of interpreting the observed depolarization of electrons in particle accelerators in terms of the Unruh effect. They argued that electrons could be used as sensitive thermometers because of the fact that the coupling between the spin and the magnetic field induces a splitting between the “spin up” and “spin down” levels. Thus, the observed depolarization of electron beams might be interpreted in the electron’s rest frame as due to the thermal bath predicted by Unruh. Since electrons in *linear accelerators* do not have time to reach equilibrium in the polarization distribution, Bell and Leinaas decided to consider electrons in *storage rings* \([6]\). In this case, ultra-relativistic electrons following uniform circular motion experience in their rest frame an “effective” thermal bath characterized by a temperature which differs from the original Unruh temperature by a numerical factor, \(\pi/\sqrt{3}\) \([7,8]\). Although other effects \([3]\) may play an important role in this phenomenon, Bell and Leinaas’ results are in good agreement with experiment. At LEP conditions, electrons are typically accelerated at \(a = 2.9 \times 10^{23} \text{m/s}^2\), which corresponds to an Unruh temperature of \(\hbar a/2\pi c k = 1200 K\). This is only 4 times larger than typical laboratory temperatures. In this vein, it would be desirable to consider a detector accelerated in a background thermal bath rather than in the inertial vacuum in or-
der to investigate in what extent finite-temperature corrections should be taken into account when testing the Unruh effect under real laboratory conditions. The detector excitations will represent the electron depolarizations, since both ones share the common feature of being two-level systems [6]. We show that in real accelerator conditions the major contribution to the detector’s response comes from the inertial vacuum rendering the contribution due to the presence of the background thermal bath unimportant. It corroborates the usual assumption, when testing the Unruh effect in storage rings, of considering the electrons as being accelerated in the Minkowski vacuum [3,10].

The paper is organized as follows: In Section II we study inertial detectors evolving in a background thermal bath and show that due to time dilatation the faster the detector moves, the less it interacts with the thermal bath. In Section III we replace the inertial detectors by uniformly accelerated ones, and calculate finite-temperature corrections in the detector’s excitation rate due to the external thermal bath. In Section IV we consider detectors moving circularly with constant speed, and discuss our results in connection with the proposal of testing the Unruh effect in storage rings. Final conclusions are summarized in Section V. Natural units will be used (ħ = c = k = 1) unless stated otherwise, and the signature adopted is (+ − − −).

II. INERTIAL DETECTORS IN A BACKGROUND THERMAL BATH

We will show in this section that the faster a detector moves in a background thermal bath the less the detector interacts with the bath. This is so because time dilatation induces a fast moving detector to interact preferentially with low frequency modes. Although a thermal bath is rich of low frequency modes, the phase space volume element (∝ ω^2 dω) suppresses infrared contributions.

Let us begin considering an Unruh-DeWitt detector [3,11]. It is basically a two-level device which may be either in the ground state |E₀⟩ or in the excited state |E⟩. The detector will be described by a monopole Ām(τ) coupled to a massless scalar field ð(χμ) through the
interaction action

\[ S_I = \int_{-\infty}^{+\infty} d\tau \, c(\tau) \hat{m}(\tau) \phi [x^\mu(\tau)] , \quad (2.1) \]

where \( x^\mu(\tau) \) is the detector’s world line and \( \tau \) is its proper time. Here \( c(\tau) \) is a switching function through which the detector is turned on/off, and plays the role of a small coupling parameter. In this section it will be enough to consider a permanently switched on detector, i.e., \( c(\tau) = c_0 = \text{const.} \). In the Heisenberg picture the monopole operator is time evolved as

\[ \hat{m}(\tau) = e^{i\hat{H}_0 \tau} \hat{m}(0) e^{-i\hat{H}_0 \tau} , \quad (2.2) \]

where \( \hat{H}_0 |E\rangle = E |E\rangle \) for any detector’s energy eigenstate \(|E\rangle\).

The amplitude for the detector to be excited and simultaneously absorb a particle \(|k\rangle\) is

\[ A_{\text{abs}}^{\text{exc}} = \langle 0 \otimes \langle E | S_I | E_0 \rangle \otimes |k\rangle \] (2.3)

Using the expansion of the scalar field in plane waves (see, e.g., [12]), and assuming that our detector follows an inertial world line \( x = y = 0; z = vt; t = \tau/\sqrt{1 - v^2} \) (where \( v = |v| \) is the detector’s speed with respect to the background thermal bath), we obtain

\[ A_{\text{abs}}^{\text{exc}} = \frac{c_0}{\sqrt{4\pi\omega}} \delta \left[ \Delta E - \frac{\omega - k_z v}{\sqrt{1 - v^2}} \right] , \quad (2.4) \]

where \( \Delta E = E - E_0 \), and \( \omega = |k| \). We will assume the selectivity \( \langle E | \hat{m}(0) | E_0 \rangle \equiv 1 \) since it only depends on the internal details of the detector, and it can be always factorized out. The amplitude for the detector to be excited and simultaneously emit a particle into the vacuum vanishes due to energy conservation. Thus, at tree level, the total excitation rate per total proper time \( T^{\text{tot}} \) of the detector will be

\[ \frac{P^{\text{exc}}}{T^{\text{tot}}} = \frac{1}{T^{\text{tot}}} \int d^3k |A_{\text{abs}}^{\text{exc}}|^2 \left[ \frac{1}{e^{\beta\omega} - 1} \right] , \quad (2.5) \]

where \( T^{\text{tot}} = 2\pi\delta(0) \) [12], and we have added into brackets the usual absorption weight associated with a thermal bath at a temperature \( \beta^{-1} \).

As a consequence of (2.4), fast moving detectors will only interact with low frequency modes. The very behavior of the detector will be determined in (2.5) by the competition
between the thermal bath, which is rich of low frequency modes, and the phase space volume element \((\propto \omega^2 d\omega)\) which tends to suppress infrared contributions. Substituting (2.4) in (2.5), and performing the integrals, we obtain

\[
\frac{\mathcal{P}^{\text{exc}}}{T^{\text{tot}}} = \frac{c_0^2 \beta^{-1} \sqrt{1 - v^2}}{4 \pi v} \ln \left[ \frac{1 - e^{-\beta \Delta E \sqrt{1+v}/\sqrt{1-v}}}{1 - e^{-\beta \Delta E \sqrt{1-v}/\sqrt{1+v}}} \right].
\]

(2.6)

In the limit \(v \to 0\) the detector responds with a Planckian spectrum

\[
\frac{\mathcal{P}^{\text{exc}}}{T^{\text{tot}}} = \frac{c_0^2}{2 \pi} \frac{\Delta E}{e^{\beta \Delta E} - 1},
\]

as expected, while as \(v \to 1\) the excitation rate per total proper time vanishes (see Fig. 1). This suggests that in testing the Unruh effect in ultra-relativistic regimes as in particle accelerators \([6,10]\), background thermal contributions should be small. This issue will be investigated in detail next.

### III. UNIFORMLY ACCELERATED DETECTORS IN A BACKGROUND THERMAL BATH

The total excitation probability for a uniformly accelerated detector evolving in a background thermal bath characterized by a temperature \(\beta^{-1}\) can be written as \([13]\)

\[
\mathcal{P}^{\text{exc}} = \int_{-\infty}^{+\infty} d\tau c(\tau) \int_{-\infty}^{+\infty} d\tau' c(\tau') e^{-i \Delta E (\tau - \tau')} G^+_{\beta}[x^\mu(\tau), x^\mu(\tau')],
\]

(3.1)

where

\[
G^+_{\beta}[x^\mu(\tau), x^\mu(\tau')] = -\sum_{n=-\infty}^{+\infty} \frac{(4\pi^2)^{-1}}{(t - t' - i \beta n - i \varepsilon)^2 - |x - x'|^2},
\]

(3.2)

is the Wightman function, \(x^\mu(\tau)\) is the world line of the accelerated detector, and \(\tau\) is its proper time. The world line of a detector moving in the \(z\)-axis with constant proper acceleration \(a\) is

\[
t = \frac{1}{a} \sinh a \tau, \quad z = \frac{1}{a} \cosh a \tau, \quad x = y = 0.
\]

(3.3)

Substituting (3.3) in the Wightman function (3.2) we obtain
\[ G^+_{\beta}(\tau, \tau') = -\frac{a^2}{16\pi^2} \sum_{n=-\infty}^{+\infty} \frac{1}{\sinh a\Delta\tau/2 + i(n\beta a - \epsilon)e^{a\xi}/2} \frac{1}{\sinh a\Delta\tau/2 + i(n\beta a - \epsilon)e^{-a\xi}/2}, \]

(3.4)

where \( \Delta\tau \equiv \tau - \tau' \), and \( \xi \equiv (\tau + \tau')/2 \). Using the following identity (the prime indicates that \( n = 0 \) is excluded from the sum)

\[ \sum_{n=-\infty}^{+\infty} \frac{1}{(A + iBn)(A + iCn)} = \frac{2}{(C - B)B} \sum_{n=1}^{+\infty} \frac{1}{(A^2/B^2 + n^2)}, \]

(3.5)

in conjunction with \[ 14 \]

\[ \sum_{n=1}^{+\infty} \frac{1}{x^2 + n^2} = \frac{\pi}{2x} \coth \pi x - \frac{1}{2x^2}, \]

(3.6)

we can cast (3.4) in the form

\[ G^+_{\beta}(\tau, \tau') = G^+_\text{vac}(\Delta\tau) + G^+_\text{ther}(\Delta\tau, \xi). \]

(3.7)

The first term in (3.7) corresponds to the pure vacuum contribution \[ 13 \]:

\[ G^+_\text{vac}(\Delta\tau) = -\frac{a^2}{16\pi^2} \sinh^{-2}(a\Delta\tau/2 - i\epsilon), \]

(3.8)

while the second term corresponds to the background thermal bath contribution:

\[ G^+_\text{ther}(\Delta\tau, \xi) = \frac{a^2}{16\pi^2} \sinh^{-2}(a\Delta\tau/2) + \frac{a}{16\pi\beta \sinh a\xi \sinh(a\Delta\tau/2)} \times \left[ \coth \frac{2\pi \sinh(a\Delta\tau/2)}{a\beta e^{a\xi}} - \coth \frac{2\pi \sinh(a\Delta\tau/2)}{a\beta e^{-a\xi}} \right]. \]

(3.9)

The fact that \( G^+_\text{ther} \) depends on \( \xi \) reflects the fact that this is a non-stationary situation. Notice that \( G^+_\text{ther}(\Delta\tau, \xi) \) does not diverge at any point. In particular \( G^+_\text{ther}(\Delta\tau = 0, \xi) = 1/12\beta^2 \). Asymptotically \( G^+_\text{ther}(\Delta\tau, \xi) \) behaves as (see Fig. 3)

\[ G^+_\text{ther}[|\Delta\tau| \gg (a^{-1}, \beta); \ |\xi| \sim e^{-a|\Delta\tau|}, \ G^+_\text{ther}[|\Delta\tau| \neq 0; \ |\xi| \gg (a^{-1}, \beta)] \sim e^{-a|\xi|}. \]

(3.10)

Clearly, \( G^+_\text{ther} \) vanishes in the limit \( \beta \to +\infty \), and thus \( G^+_{\beta \to +\infty}(\tau, \tau') = G^+_\text{vac}(\Delta\tau) \).

Now, we are ready to investigate the total excitation rate, \( \mathcal{P}^\text{exc} = \mathcal{P}^\text{exc}_{\text{vac}} + \mathcal{P}^\text{exc}_{\text{ther}} \), of a detector uniformly accelerated in a background thermal bath. In order to be realistic we
shall consider the detector as being switched on only during a finite period of proper time $|\tau| < T_0/2$, where $T_0 = \text{const} \in \mathbb{R}_+$. Concerning the pure vacuum contribution, $\mathcal{P}_{\text{vac}}^{\text{exc}}$, it could be calculated for some continuous $c(\tau)$ by letting (3.8) into (3.1). Notwithstanding, we will use directly the results of Ref. [15] where the calculations were performed in the detector’s rest frame. The excitation probability for a detector uniformly accelerated in the inertial vacuum, and kept switched on for long enough, $T_0 \gg a^{-1}, \Delta E^{-1}$, is

$$\mathcal{P}_{\text{vac}}^{\text{exc}} \approx \frac{c_0^2}{2\pi \epsilon^2 \Delta E/a} \frac{\Delta E}{a} T_0.$$  

(3.11)

Notice the linear dependence with $T_0$ (see Fig. 3), which is exactly what one should expect due to the Unruh effect [13]. Here $c_0 = \text{const}$ is the coupling constant between the field and the monopole while the detector is kept switched on. In the regime considered above, the detailed form of $c(\tau)$ is not important. The only restriction is that $c(\tau) \in C^0$, since discontinuities in $c(\tau)$ would result in ultraviolet divergences [15].

The thermal correction, $\mathcal{P}_{\text{ther}}^{\text{exc}}$, on the pure vacuum term (3.11) will be obtained by introducing (3.9) in (3.1)

$$\mathcal{P}_{\text{ther}}^{\text{exc}} = c_0^2 \int_{-T_0/2}^{+T_0/2} d\tau \int_{-T_0/2}^{+T_0/2} d\tau' e^{-i\Delta E\Delta \tau} G_{\text{ther}}^+(\Delta \tau, \xi),$$  

(3.12)

where we have already considered the fact that the detector is kept switched on only during a finite amount of proper time $T_0$. The integrals above were solved numerically for $\Delta E = 9.7 \times 10^{14} \text{s}^{-1}$, $a = 9.7 \times 10^{14} \text{s}^{-1}$, $\beta = 2.5 \times 10^{-14} \text{s}$, and plotted as a function of $T_0$ in Fig. 3. These values of $\Delta E$, $a$, and $\beta$ have a clear physical motivation. Electrons in particle accelerators have their spin coupled to the magnetic field. It induces a splitting of the “spin up”, and “spin down” levels. The energy gap associated with such a splitting is $\Delta E = 2|\mu||B|$, where $\mu \approx e/2m_e$ is the electron’s magnetic moment (it is assumed the gyromagnetic factor to be $g = 2$), and $B$ is the magnetic field. Following [13] we consider the depolarization of an accelerated electron as representing the excitation of the detector, since both ones share the common feature of being two–level systems. More detailed calculations are not supposed to change the order of magnitude of the results obtained, and consequently our conclusions.
At LEP conditions an electron has a typical Lorentz factor of $\gamma \equiv (1 - v^2)^{-1/2} = 10^5$, and proper acceleration of $a = 2.9 \times 10^{23} m/s^2$, which corresponds to $a = 9.7 \times 10^{14} s^{-1}$ in natural units. The energy gap between the two spin levels is $\Delta E \approx a = 9.7 \times 10^{14} s^{-1}$\cite{[6]}. The lab-time for building up the polarization is about 2 hours, which corresponds to $7.2 \times 10^{-2} s$ in the electron’s proper time. Finally, the background thermal bath has a temperature corresponding to $\beta = (300 K)^{-1} \approx 2.5 \times 10^{-14} s$. From Fig. 3 it is clear that the pure vacuum contribution, $P_{\text{exc vac}}$, increases with $T_0$ much faster than the background thermal contribution, $P_{\text{exc ther}}$. This is a consequence of the fact that $G_{\text{ther}}^+$ decreases exponentially for large $|T_0|$ (Fig. 2), except along the $\xi$-axis. Eventually, Fig. 3 just reflects the fact that the faster the detector moves the less it interacts with the background thermal bath as discussed in Section I.

Corrections on $P_{\text{exc vac}}$ due to the background thermal contribution, $P_{\text{exc ther}}$, are only important for particles accelerated during short periods of time. It might be the case in some regimes of heavy-ion collisions. Barshay and Troost suggested that the large transverse acceleration $a$ of projectile and target which occurs in high-energy hadronic collisions is connected with the thermal emission of particles at Unruh temperature of about $100 MeV \approx 10^{12} K$\cite{[16]}. In acceleration regimes in which the Unruh temperature is smaller, background thermal contributions may be important. In this vein it might be useful to expand $G_{\text{ther}}^+(\Delta \tau, \xi)$ in terms of $\beta$ factors. For this purpose, we have used Eq.(1.411.8) of [14] in (3.9) obtaining

$$G_{\text{ther}}^+(\Delta \tau, \xi) = \frac{1}{12 \beta^2} - \sum_{n=2}^{+\infty} (\frac{4\pi)^2n^{-2}}{(2n)!a^{2n-2}\beta^{2n} \sinh a^2 \eta} \sinh[(2n-1)\frac{a\Delta \tau}{2}] \sinh[(2n-1)a^2 \eta],$$

(3.13)

where $B_n$ are the Bernoulli numbers, and $\beta > 2a^{-1}|e^{a\xi} \sinh(a\Delta \tau/2)|$. In the next section we will analyze a detector moving circularly, and discuss our results in connection with the proposal of testing the Unruh effect in storage rings by looking at the polarization distribution of the accelerated electrons.
IV. CIRCULARLY MOVING DETECTORS IN A BACKGROUND THERMAL BATH

The world line of a detector describing a circular motion with radius \( R \), and constant speed \( v \) is

\[
t = \gamma \tau, \quad x = R \cos \omega \gamma \tau, \quad y = R \sin \omega \gamma \tau, \quad z = 0,
\]

where \( \omega = v/R \). The proper acceleration of the detector is \( a \equiv \sqrt{a_{\mu}a^{\mu}} = v^2\gamma^2/R \), where \( a_{\mu} = u^{\nu} \nabla_{\nu} u_{\mu} \).

Substituting (4.1) in (3.2), and using (3.5) and (3.6) we decompose the relevant Green function again in a pure vacuum part, and in a background thermal part as

\[
D_\beta^+(\Delta \tau) = D_{\text{vac}}^+(\Delta \tau) + D_{\text{ther}}^+(\Delta \tau),
\]

where

\[
D_{\text{vac}}^+(\Delta \tau) = (4\pi^2)^{-1} \left[ -\gamma^2(\Delta \tau - i\epsilon)^2 + 4v^4\gamma^4 \sin^2(a\Delta \tau/2v\gamma)/a^2 \right]^{-1},
\]

and

\[
D_{\text{ther}}^+(\Delta \tau) = -\frac{1}{4\pi^2\gamma^2\Delta \tau^2} \left\{ \left[ \frac{4v^4 \sin^2(A\Delta \tau/2v)}{A^2\Delta \tau^2} - 1 \right]^{-1} + \frac{\pi A\Theta \Delta \tau^2}{4v^2 \sin(A\Delta \tau/2v)} \right\}
\times \left[ \coth \left( \pi \Theta \Delta \tau \left( 1 - \frac{2v^2}{A\Delta \tau} \sin \frac{A\Delta \tau}{2v} \right) \right) - (v \rightarrow -v) \right].
\]

Here \( A \equiv a/\gamma \), and \( \Theta \equiv \gamma/\beta \) play the role of an effective proper acceleration, and an effective background temperature respectively. \( D_{\text{ther}}^+(\Delta \tau) \) is finite everywhere. In particular, \( \lim_{\Delta \tau \rightarrow 0} D_{\text{ther}}^+(\Delta \tau) = 1/12\beta^2 \). Asymptotically, \( D_{\text{ther}}^+(\Delta \tau \gg 1) \sim (4\pi^2\gamma^2\Delta \tau^2)^{-1} \). The fact that \( D_{\beta}^+ \) does not depend on \( \xi \) reflects the fact that this situation is stationary.

In order to calculate the average vacuum excitation rate, \( dP_{\text{vac}}^{\text{exc}}/dT \), for ultra-relativistic detectors it is convenient to express (4.3) as

\[
D_{\text{vac}}^+(\Delta \tau) = (4\pi^2)^{-1} \left[ -(\Delta \tau - i\epsilon)^2 - a^2\Delta \tau^4/12 + O(\gamma^{-1}) \right].
\]
Substituting (4.5) in (3.1) we obtain for \( \gamma >> 1 \)

\[
\frac{dP_{\text{vac}}^{\text{exc}}}{dT} \approx \frac{c_0^2 \alpha e^{-\sqrt{12} \Delta E/a}}{4\pi \sqrt{12}}.
\]  

Thus at LEP we expect \( dP_{\text{vac}}^{\text{exc}} /dT \approx 7.0 \times 10^{11} c_0^2 \).

In order to compare this result with the average thermal contribution, \( dP_{\text{ther}}^{\text{exc}} /dT \), we substitute (4.4) in (3.1) obtaining

\[
\frac{dP_{\text{ther}}^{\text{exc}}}{dT} = c_0^2 \int_{-\infty}^{+\infty} d(\Delta \tau) \ e^{-i\Delta E \Delta \tau} D_{\text{ther}}^{+}(\Delta \tau).
\]  

Evaluating numerically this expression with LEP values we obtain \( dP_{\text{ther}}^{\text{exc}} /dT \approx 3 \times 10^8 c_0^2 \).

Thus, after the equilibrium in the polarization distribution is reached the pure vacuum contribution is expected to be about 3 orders of magnitude larger than the background thermal contribution. Before concluding, we notice that for “quasi-inertial” detectors, i.e. \( A << \Theta, \Delta E \), (4.7) can be approximated by the inertial formula (2.6). In particular, in the limit \( A \to 0 \), \( dP_{\text{ther}}^{\text{exc}} /dT \) turns out to be exactly (2.6). At LEP conditions we have \( A = 9.7 \times 10^9 s^{-1} \), \( \Theta = 4.0 \times 10^{18} s^{-1} \), and \( \Delta E = 9.7 \times 10^{14} s^{-1} \). Thus one could have used directly (2.6) to estimate the average thermal contribution obtaining \( dP_{\text{ther}}^{\text{exc}} /dT \approx 2.9 \times 10^8 c_0^2 \).

V. CONCLUSION

We have derived the response of inertial and accelerated detectors in a background thermal bath. The faster the detector moves, the less important will be the background thermal bath. This is so because time dilatation induces the detector to interact only with the low frequency modes present in the bath. Although the thermal bath is rich of low frequency modes, the phase space volume element suppresses infrared contributions in the excitation probability. Bell and Leinaas suggested that the depolarization of electrons in storage rings could be explained through the Unruh effect, i.e. due to the appearance of a thermal bath in the electron’s rest frame of about 1200K. In their analysis it is assumed that the electrons are accelerated in the inertial vacuum. We have estimated whether considering the fact that
the electrons are actually accelerated in a *background thermal bath* of about 300 K would add or not any substantial contribution in the depolarization rate. We obtain under LEP conditions the interesting result that although the Unruh thermal bath is only about 4 times hotter than the background thermal bath, the term $P_{\text{ther}}^{\text{exc}}$ due to the external bath is various orders of magnitude smaller than the pure vacuum contribution $P_{\text{vac}}^{\text{exc}}$. Concerning the proposal of testing the Unruh effect in storage rings, it corroborates the usual assumption of considering the electrons as being accelerated in the inertial vacuum [6,10]. Notwithstanding, according to our results, we expect that background thermal corrections will be relevant in situations where non ultra-relativistic particles are accelerated moderately for short periods of time.

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FIGURES

FIG. 1. The excitation probability of an inertial detector moving with speed $v$ in a thermal bath characterized by a temperature $\beta^{-1}$ is plotted. The faster the detector moves the less it interacts with the thermal bath. Here we have used $\beta = 2.5 \times 10^{-14} s$, and $\Delta E = 9.7 \times 10^{14} s^{-1}$.

FIG. 2. $G^+_\text{ther}(\Delta \tau, \xi)$ is finite everywhere, and asymptotically it decreases exponentially except along the $\xi$-axis on which it is constant. Notice that $G^+_\text{ther}(\Delta \tau, \xi)$ is completely symmetric in the other quadrants. Here we have used $\beta = 2.5 \times 10^{-14} s$, and $a = 9.7 \times 10^{14} s^{-1}$.

FIG. 3. $P^\text{exc}_\text{vac}$ and $P^\text{exc}_\text{ther}$ are plotted as a function of $T_0$. $P^\text{exc}_\text{vac}$, which is represented with a full line, increases linearly for large $T_0$, while $P^\text{exc}_\text{ther}$, which is represented with a dotted line, increases much slower. This is a consequence of the fact that $G^+_\text{ther}$ decreases exponentially for large $T_0$, except along the “monorail” $\Delta \tau = 0$. Eventually, it reflects the fact that the faster the detector moves with respect to the background thermal bath, the less the detector interacts with it. Here we have used $\beta = 2.5 \times 10^{-14} s$, $a = 9.7 \times 10^{14} s^{-1}$, and $\Delta E = 9.7 \times 10^{14} s^{-1}$.
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