# Neutrino oscillations in the Unruh radiation

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The study of the decay of an accelerated proton recently provided a “theoretical proof” of the Unruh effect. On the basis of general covariance of Quantum Field Theory, indeed, it was found that the decay rates in the inertial and comoving frames do coincide only when the thermal nature of the accelerated vacuum is taken into account. Such an analysis was then extended to the case with neutrino mixing. In this Letter, we show that, by further embedding neutrino oscillations in the above framework, general covariance necessarily entails the use of flavor neutrinos as asymptotic states, as well as the occurrence of neutrino oscillations in the Unruh thermal bath.

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Introduction. In the history of Physics, the adoption of principles has revealed to be a formidable investigation tool. Although intimately related to the phenomenological realm from which they stem, once elevated to the status of postulates, physical principles act as lighthouses for the development of a consistent theoretical apparatus. Paradigmatic examples are the principle of conservation of energy, which led for instance to the discovery of the neutrino, and the principle of constancy of speed of light, at the basis of the construction of Special Relativity.

Recently, general covariance has been advocated to exhibit that the internal consistency of quantum field theory (QFT) unavoidably requires the existence of the Unruh effect, also known as Fulling-Davies-Unruh effect [1–3]. Indeed, starting from the statement that acceleration can influence even the proper lifetime of stable particles [4], in a series of remarkable papers [5] it was shown that the tree-level decay rate of an accelerated proton via inverse β-decay is frame-independent only when the thermal nature of the vacuum for a non-inertial observer is considered.

The fact that a theoretical requirement leads to the specific form of the ground state for an accelerated observer should be regarded as a considerable result, especially in view of the perplexities which have been sometimes raised about the physical significance of the Unruh effect [6]. Such a skepticism is enhanced by the lack of direct evidences of this phenomenon, as it also happens for the case of the Hawking radiation [7]. In fact, at present the most likely arena for (indirect) experimental tests of these effects is given by analogue gravity [8].

In the above cited studies on the accelerated proton decay, the emitted neutrino was initially treated as massless and only in Ref. [9] as a massive particle. In these works, however, neutrino mixing was not taken into account. This was done for the first time in Ref. [10], where a discrepancy between the proton decay rate in the inertial and comoving frames was claimed to arise. Subsequently, it was proved that general covariance does indeed hold in the above analysis: this has been achieved by employing either flavor [11] or mass [12] eigenstates for neutrinos, thus leaving an essential ambiguity on the very nature of asymptotic neutrino states.

In this Letter we show that, due to the occurrence of neutrino oscillations in the problem at hand, general covariance leads to the conclusion that flavor eigenstates are fundamental and neutrinos belonging to the Unruh radiation must oscillate. Throughout the paper, we use natural units $\hbar = c = 1$ and the Minkowski metric with the mostly negative signature.

Let us start by setting the framework. Following the approach and the notation of Refs. [5], the proton $|p\rangle$ and the neutron $|n\rangle$ can be viewed as unexcited and excited states of the nucleon, respectively. In addition, we assume to deal with particles that are energetic enough to possess a well-defined trajectory. In these conditions, it is possible to employ the Fermi theory of current-current interaction, where we consider a quantum leptonic and a classical hadronic current $j^{\mu}_{\ell} \to j^{\mu}_{\ell} j^{(c)}_{\mu}$, with $j^{(c)}_{\mu} = q(\tau) u_{\mu} \delta(x) \delta(y) \delta(u - a^{-1})$. Here, $\tau = u/a$ is the nucleon’s proper time (with $v$ being the Rindler time coordinate), $a$ its proper acceleration and $u = a^{-1} = \text{const}$ represents the spatial Rindler coordinate which denotes the world line of the uniformly accelerated particle. The four-velocity of the nucleon $u^\mu$ is given by $u^\mu = (a, 0, 0, 0)$, $u^\mu = (\sqrt{a^2 + 1}, 0, 0, at)$, in Rindler and Minkowski coordinates, respectively. In accordance

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¹ We assume the acceleration to occur along the $z$-direction. With this choice, the Rindler coordinates $(v, x, y, u)$ are related with the Minkowski coordinates $(t, x, y, z)$ as follows: $t = u \sinh v$, $z = u \cosh v$, with $x$ and $y$ untouched.
with Refs. [5, 13], the Hermitian monopole \( \hat{q}(\tau) \) is given by \( \hat{q}(\tau) \equiv e^{iH\tau} \hat{q}_0 e^{-iH\tau} \), where \( H \) is the nucleon Hamiltonian and \( \hat{q}_0 \) is used to reconstruct the effective Fermi constant \( G_F \equiv \langle p | \hat{q}_0 | n \rangle \).

Assuming to deal with a simplified two-flavor model, the interaction of the charged leptons \( \hat{\Psi}_\alpha \) and neutrinos \( \hat{\Psi}_\nu(\alpha = e, \mu) \) with the nucleon current \( j^{(cl)}_\mu \) is described by the Fermi action

\[
\hat{S}_I = \sum_{\alpha = e, \mu} \int d^4x \sqrt{-g} j^{(cl)}(x) \left( \bar{\hat{\Psi}}_{\nu, \alpha} \gamma^\mu \hat{\Psi}_{\nu, \alpha} + \bar{\hat{\Psi}}_{\nu, \alpha} \gamma^\mu \hat{\Psi}_{\nu, \alpha} \right),
\]

where \( g \equiv \text{det}(g_{\mu\nu}) \) and \( \gamma^\mu \) are the gamma matrices in Dirac representation (e.g., see Ref. [14]). In Eq. (1), neutrino fields with definite flavors are related to the ones with definite masses by

\[
\hat{\Psi}_{\nu} = \cos \theta \hat{\Psi}_{\nu_1} + \sin \theta \hat{\Psi}_{\nu_2},
\]

\[
\hat{\Psi}_{\nu_3} = -\sin \theta \hat{\Psi}_{\nu_1} + \cos \theta \hat{\Psi}_{\nu_2},
\]

In what follows, however, we shall focus on the processes involving only the production of a positron. The case in which an anti-muon appears in the final state can be treated analogously and in an independent way.

**Inertial frame.** In the inertial frame, the process we consider is the decay of an accelerated proton via inverse \( \beta \)-decay

\[(i) \quad p \rightarrow n + e^+ + \nu_e, \quad (3)\]

which is pictorially represented in Fig. 1.

The fermion field is expanded as

\[
\hat{\Psi} = \sum_{\sigma = \pm} \int \frac{d^3k}{4\pi^2} \left[ e^{-ik^\mu x_\mu} \bar{u}_\sigma(k) \bar{b}_\sigma(k) + e^{ik^\mu x_\mu} \bar{u}(-\omega) \gamma^\mu \bar{d}_\sigma(k) \right],
\]

where \( \sigma \) is the polarization, \( \omega = \sqrt{k^2 + m^2} \) and \( \bar{u}_\sigma(\omega) \) is the spinor defined as in Ref. [11].

The tree-level transition amplitude for the process (i) reads [11]

\[
\mathcal{A}^{(\nu_e)}_{(i)} = \langle n | \langle e^+, \nu_e | \hat{S}_I | 0 \rangle \langle 0 | p \rangle \]

\[
\mathcal{A}^{(\nu_e)}_{(i)} = \frac{G_F}{24\pi^3} \left[ \cos^2 \theta \mathcal{I}_{\sigma, \bar{\sigma}}(\omega_{\nu_1}, \omega_e) + \sin^2 \theta \mathcal{I}_{\sigma, \bar{\sigma}}(\omega_{\nu_2}, \omega_e) \right],
\]

where for simplicity we have omitted the \( \mathbf{k} \) and \( \sigma \)-dependence of the lepton states, we have assumed equal momenta and polarizations for neutrinos with definite masses and

\[
\mathcal{I}_{\sigma, \bar{\sigma}}(\omega_{\nu_1}, \omega_e) = \int_{-\infty}^{+\infty} d\tau \mu_{\nu} \left[ u_{\nu}^{(+\omega_{\nu_1})} \gamma^\mu_{\nu} \bar{u}_{\nu}^{(-\omega_e)} \right]
\]

\[
\times e^{i \left[ \Delta m \tau + a^{-1}(\omega_{\nu_1} + \omega_e) \sinh \alpha \tau - a^{-1}(k^e + k^\nu) \cosh \alpha \tau \right]},
\]

with \( j = 1, 2 \).

**Fig. 1:** Decay process in the inertial frame without (i) and with (ii) neutrino oscillation. Time flows in the vertical direction.

Note that the asymptotic flavor neutrino state \( |\nu_e\rangle \) in Eq. (5) has been expressed in terms of the corresponding mass states \( |\nu_i\rangle (i = 1, 2) \) by means of Pontecorvo mixing transformations [15]

\[
|\nu_e\rangle = \cos \theta |\nu_1\rangle + \sin \theta |\nu_2\rangle, \quad (7a)
\]

\[
|\nu_\mu\rangle = -\sin \theta |\nu_1\rangle + \cos \theta |\nu_2\rangle. \quad (7b)
\]

The differential transition probability is defined as

\[
\frac{d^3P^{(\nu_e)}_{(i)}}{d^3k_e d^3k_\nu} = \sum_{\sigma, \bar{\sigma}} |\mathcal{A}^{(\nu_e)}_{(i)}|^{2}, \quad (8)
\]

and the transition rate is given by \( \Gamma = P/T \), with

\[
T = \int_{-\infty}^{+\infty} d\tau
\]

being the nucleon total proper time.

In Ref. [11], it has been shown that the transition rate for the process Eq. (8) is

\[
\Gamma_{in}^{(\nu_e)} = \cos^4 \theta \Gamma_1 + \sin^4 \theta \Gamma_2 + \cos^2 \theta \sin^2 \theta \Gamma_{12}. \quad (10)
\]

Here, we have introduced the shorthand notation

\[
\Gamma_{j} = \frac{1}{T} \sum_{\sigma, \bar{\sigma}} \frac{G_F^2}{28\pi^6} \int d^3k_\nu \int d^3k_e |\mathcal{I}_{\sigma, \bar{\sigma}}(\omega_{\nu_j}, \omega_e)|^2, \quad (11)
\]

where \( j = 1, 2 \) and

\[
\Gamma_{12} = \frac{1}{T} \sum_{\sigma, \bar{\sigma}} \frac{G_F^2}{28\pi^6} \int d^3k_\nu \int d^3k_e
\]

\[
\times \left[ \mathcal{I}_{\sigma, \bar{\sigma}}(\omega_{\nu_1}, \omega_e) \mathcal{I}_{\sigma, \bar{\sigma}}^*(\omega_{\nu_2}, \omega_e) + \text{c.c.} \right]. \quad (12)
\]

The aim of the calculations contained in Ref. [11] is to exhibit the equality between Eq. (10) and its counterpart in the accelerated frame (see below), which guarantees the validity of the principle of general covariance.
At this point, it must be emphasized that in the above calculation an infinite proper time interval is considered, which allows for the electron neutrino produced in the proton decay to oscillate. Thus, we should take into account not only the process contemplated in Eq. (3), but also the following one:

(ii) \[ p \to n + e^+ + \nu_{\mu}. \] (13)

The above relation must be intended in the sense of Fig. 1: although it is true that the lepton charge must necessarily be conserved in the vertex (at tree-level), as soon as the outgoing neutrino is produced, flavor oscillations will inevitably take place. We remark that this process has not been included in the analysis of Ref. [11], without however affecting the validity of the results there contained.

The transition amplitude for the process (ii) is indeed non-vanishing and is given by

\[
\mathcal{A}_{(ii)}^{(\nu_e)} = \langle n| \langle e^+, \nu_{\mu}| \hat{S}_t | 0 \rangle \otimes | p \rangle = -\frac{G_F}{2\sqrt{\pi}} \cos \theta \sin \theta \left[ \mathcal{I}_{\sigma_e, \sigma_e} (\omega_{\pi e}, \omega_e) - \mathcal{I}_{\sigma_e, \sigma_e} (\omega_{\nu e}, \omega_e) \right].
\] (14)

In terms of \( \Gamma \), the quantity \( \mathcal{A}_{(ii)}^{(\nu_e)} \) of Eq. (14) associated to the process (ii) leads to the following transition rate:

\[
\Gamma_{in}^{(\nu_e)} = \cos^2 \theta \sin^2 \theta (\Gamma_1 + \Gamma_2 - \Gamma_{12}),
\] (15)

where all the contributions in the r.h.s. of the above expression have already been introduced in Eqs. (11) and (12). We notice that the above transition rate is proportional to \( \sin^2 2\theta \), thus showing that it is originated by interference.

Finally, observe that

\[
\Gamma_{in} \equiv \Gamma_{in}^{(\nu_e)} + \Gamma_{in}^{(\nu_{\mu})} = \cos^2 \theta \Gamma_1 + \sin^2 \theta \Gamma_2.
\] (16)

**Accelerated frame.** From the point of view of an observer comoving with the proton, the only possible way to render the particle’s decay viable is to suppose the existence of a thermal bath of electrons, neutrinos and the corresponding antiparticles. In such conditions, recalling that here we consider the electron decay channel only, three processes can be taken into account to match the inertial process Eq. (3):

(iii) \[ p^+ + e^- \to n + \nu_e, \] (17a)

(iv) \[ p^+ + \bar{\nu}_e \to n + e^+, \] (17b)

(v) \[ p^+ + e^- + \bar{\nu}_e \to n. \] (17c)

In order to calculate the proton’s decay rate in the accelerated frame, we need the expansion for the fermion fields in Rindler coordinates, that is

\[
\hat{\Psi} = \sum_{\sigma = \pm} \int_0^{+\infty} d\omega \int \frac{d^2k}{(2\pi)^2} \left[ e^{i(-\omega_{\nu}/a + k_{\alpha} x^\alpha)} \tilde{u}_\sigma(\omega) \tilde{b}_\sigma \right. \\
+ e^{i(-\omega_{\nu}/a + k_{\alpha} x^\alpha)} \left. \tilde{u}_\sigma(-\omega) \right| \tilde{d}_\sigma \right],
\] (18)

where \( k_{\alpha} x^\alpha = k_x x^x + k_y y^y \), \( \omega \) is the Rindler frequency which does not satisfy any dispersion relation, \( w = (\omega, k^x, k^y) \) and \( \tilde{u}_\sigma(\omega) \) is defined as in Ref. [11].

By way of illustration, let us consider the scattering (iii) of Eq. (17a) (see also Fig. 2); similar calculations can be carried out for the processes in Eqs. (17b), (17c). We obtain [11]

\[
\mathcal{A}_{(iii)}^{(\nu_e)} \equiv \langle n| \langle \nu_e | \hat{S}_t | e^- \rangle \otimes | p \rangle = \frac{G_F}{(2\pi)^2} \left[ \cos^2 \theta j^{(1)}_{\sigma_e, \sigma_e}(\omega_v, \omega_e) + \sin^2 \theta j^{(2)}_{\sigma_e, \sigma_e}(\omega_v, \omega_e) \right],
\] (19)

where we have assumed equal frequencies and polarizations for different neutrinos with definite masses, whereas the quantity denoted with \( J \) is equal to

\[
J^{(i)}_{\sigma_e, \sigma_e}(\omega_v, \omega_e) = \delta(\omega_v - \omega_e - \Delta m) \tilde{u}^{(i, \omega_v)}_{\sigma_e} \tilde{u}^{(i, \omega_v)}_{\sigma_e}. \] (20)

with \( i = 1, 2 \). The spinor component related to the neutrino field contains the information on the mass of \( \nu_i \), and by means of the current hypothesis this is the only difference between the functions \( J^{(i)} \) with different indexes.

The sum of the transition rates for the three processes in Eqs. (17) yields [11]

\[
\Gamma^{(\nu_e)}_{acc} = \cos^4 \theta \tilde{\Gamma}_1 + \sin^4 \theta \tilde{\Gamma}_2 + \cos^2 \theta \sin^2 \theta \tilde{\Gamma}_{12},
\] (21)

where

\[
\tilde{\Gamma}_j \equiv \mathcal{N} \int_{-\infty}^{+\infty} d\omega \mathcal{R}_j(\omega), \quad j = 1, 2,
\] (22)

and \( \mathcal{N} = 2\pi^{-7} \pi^{-2} G_F e^{-\pi \Delta m/a} \). The functions \( \tilde{\Gamma}_{12} \) and \( \mathcal{R}_j \) in Eqs. (21) and (22), respectively, are defined by
ν
ν
δm
to each other only up to a first-order expansion in the

that \( \Gamma \)

the corresponding decay rate \( \Gamma \)

involutions as potential candidates for the non-inertial counter-

eral covariance, we now seek the corresponding processes

considered, which in the inertial frame is represented by

additional contribution to the proton decay rate has to be

act flavor neutrino states, defined as eigenstates of flavor

Pontecorvo states Eq. (7) can be identified with the ex-

ings in the Unruh thermal bath are considered in the

FIG. 2: Decay processes in the accelerated frame. Oscillations of neutrinos in the Unruh thermal bath are considered in the

last two diagrams. Time flows in the vertical direction.

\[
\mathcal{R}_j(\omega) = \int d^2k_\nu d^2k_e l_{\nu_j} l_e \left| K_{\frac{\omega}{2} + \frac{1}{2}} \left( \frac{l_{\nu_j}}{a} \right) K_{\frac{\omega}{2} - \frac{1}{2}} \left( \frac{l_e}{a} \right) \right|^2 + m_{\nu_j} m_e \text{Re} \left[ \int d^2k_\nu d^2k_e K^2_{\frac{\omega}{2} + \frac{1}{2}} \left( \frac{l_{\nu_j}}{a} \right) K^2_{\frac{\omega}{2} - \frac{1}{2}} \left( \frac{l_e}{a} \right) \right], \tag{23}
\]

and

\[
\bar{\Gamma}_{12} = \frac{N}{\sqrt{l_{\nu_1} l_{\nu_2}}} \int d\omega d^2k_\nu d^2k_\nu \left\{ l_{\nu_1} K_{\frac{\omega}{2} + \frac{1}{2}} \left( \frac{l_{\nu_1}}{a} \right) \right\}^2 \left( k^2_\nu + m_{\nu_1} m_{\nu_2} + l_{\nu_1} l_{\nu_2} \right) \text{Re} \left[ K_{\frac{\omega}{2} + \frac{1}{2}} \left( \frac{l_{\nu_1}}{a} \right) K_{\frac{\omega}{2} - \frac{1}{2}} \left( \frac{l_{\nu_2}}{a} \right) \right] + m_e (l_{\nu_1} m_{\nu_2} + l_{\nu_2} m_{\nu_1}) \text{Re} \left[ K^2_{\frac{\omega}{2} + \frac{1}{2}} \left( \frac{l_{\nu_1}}{a} \right) K^2_{\frac{\omega}{2} - \frac{1}{2}} \left( \frac{l_{\nu_2}}{a} \right) \right] \right\}, \tag{24}
\]

where \( l_{\nu_j} \equiv \sqrt{m_{\nu_j}^2 + (k^x)^2 + (k^y)^2} \) and we have used

the shorthand notation \( \bar{\omega} \equiv \omega - \Delta m \).

In Ref. [11] it was shown that, by taking into account

the three interactions in Eq. (17) of the proton with the

particles of the thermal bath, the quantity \( \Gamma_{\text{acc}} \) matches

the corresponding decay rate \( \Gamma_{\text{in}} \) evaluated in the

inertial frame. In particular, it is possible to exhibit [11]

that \( \Gamma_i = \bar{\Gamma}_i \) for \( i = 1, 2 \), whereas \( \Gamma_{12} \) and \( \Gamma_{12} \) are equal to

each other only up to a first-order expansion in the

parameter \( \delta m \equiv m_{\nu_2} - m_{\nu_1} \). In such an approximation, Pontecorvo states Eq. (7) can be identified with the exact

flavor neutrino states, defined as eigenstates of flavor charges [16] [17].

On the other hand, we have seen above that an ad-

ditional contribution to the proton decay rate has to be

considered, which in the inertial frame is represented by

the process (ii) in Fig. 1. Guided by the principle of gen-

eral covariance, we now seek the corresponding processes

in the comoving frame which should lead to the same

result.

To this aim, we consider the three following contribu-

tions as potential candidates for the non-inertial counter-

part of the decay Eq. (13):

\[
\begin{align*}
\text{(vi)} & \quad p^+ + e^- \rightarrow n + \nu_\mu, \tag{25a} \\
\text{(vii)} & \quad p^+ + \bar{\nu}_\mu \rightarrow n + e^+, \tag{25b} \\
\text{(viii)} & \quad p^+ + e^- + \bar{\nu}_\mu \rightarrow n, \tag{25c}
\end{align*}
\]

which are depicted in Fig. 2 in the last three diagrams. Note that, while the process (vi) is of the same type of
(ii) since it entails an oscillation of the emitted (electron)
neutrino, the processes (vii) and (viii) are essentially due
to the oscillation of an (muon) antineutrino that is al-
ready present in the Unruh thermal bath.

In order to legitimate the validity of our assumption, we need to perform the same calculations that lead to the
decay rate of Eq. (21). The outcome of this procedure

turns out to be

\[
\Gamma_{\text{acc}} = \cos^2 \theta \sin^2 \theta \left( \bar{\Gamma}_1 + \bar{\Gamma}_2 - \bar{\Gamma}_{12} \right). \tag{26}
\]

By virtue of the aforesaid observations contained in detail

in Ref. [11] which allow us to state that \( \Gamma_{\text{in}} = \Gamma_{\text{acc}} \), it

is possible to infer that such an equivalence holds also

between the decay rates of Eqs. (13) and (26).

Moreover, if we compute the total comoving decay rate

which includes neutrino oscillations, we deduce that

\[
\Gamma_{\text{acc}} = \Gamma_{\text{acc}} + \Gamma_{\text{acc}} = \cos^2 \theta \bar{\Gamma}_1 + \sin^2 \theta \bar{\Gamma}_2. \tag{27}
\]
By comparing this with the total inertial decay rate of Eq. (16), we find that
\[ \Gamma_{in} = \Gamma_{acc}, \] (28)
which means that such a result does not depend on the quantities \( \Gamma_{12} \) and \( \Gamma_{e2} \), whose treatment would require additional computational effort [11].

Remarkably, Eq. (28) not only involves a generalization of the analysis of the accelerated proton decay to the case in which the produced neutrino oscillates, but it also unambiguously corroborates our guess of selecting the processes in Eqs. (25) as the counterpart for the decay (ii) in the inertial frame. Hence, the requirement of the principle of general covariance clearly results in the necessity of having an Unruh thermal bath containing flavor neutrinos which do oscillate.

**Conclusions.** In this Letter, we have extended the formalism firstly introduced in Refs. [5] and then developed in Refs. [13, 14] to the case in which neutrino oscillations are taken into account.

In these works, the principle of general covariance has been exploited in the analysis of accelerated proton decay to probe the internal consistency of the theory and to shed light on new aspects of this problem. In such a perspective, here we have further studied this subject by considering flavor oscillations for the emitted neutrinos. By enforcing general covariance, we have then found that the Unruh radiation “seen” by the accelerated proton must involve oscillating neutrinos. This is a novel feature which emerges in a natural way in the present approach.

A further interesting observation that can be deduced from our analysis is related to the identities Eqs. (16) and (27) that are true for the inertial and the comoving frame, respectively. For this purpose, we recall that the equality has to be regarded both in the inertial and the comoving frames. Such an equation constitutes a consistency check for the correctness of the calculations in Refs. [11] and [14]. The physical meaning of Eq. (29) can be understood by considering the charges for mixed neutrino fields as derived from Noether’s theorem [17]. Indeed, by denoting with
\[ Q_i = \int d^3x \Psi\bar{\nu}_{\alpha}(x)\Psi_{\nu_i}(x), \quad i = 1, 2, \]
the (time-dependent) flavor charges, one can see that [17]
\[ Q = \sum_i Q_i = \sum_{\alpha} Q_{\alpha}(t), \]
where \( Q \) represents the total charge. The above relation can be interpreted as the conservation of the total lepton number. On the one hand, this can be viewed as the sum of two separately conserved family lepton numbers, when no mixing is present; on the other hand, the same conserved number is obtained by the sum of non-conserved flavor charges, which indeed are associated to neutrino oscillations.

Apart from its relevance in the context of neutrino mixing and oscillations, we stress that the Unruh effect is an excellent benchmark for both testing well-established predictions and pointing out novel effects in fundamental physics, as it combines such wide domains as general relativity, quantum mechanics and thermodynamics. For instance, in Refs. [18] it has been shown that the Unruh spectrum may exhibit exotic non-thermal corrections even within the standard QFT, thus emphasizing how such a framework provides an active forge of still unexplored scenarios. A similar non-thermal behavior has been obtained in Refs. [19], where Planck scale effects on the Unruh bath have been derived in the context of the Generalized Uncertainty Principle [20]. Further features of the Unruh effect may be addressed in connection with entanglement properties for accelerated observers, whose implications have been investigated also in the context of black hole physics [21]. In particular, in Ref. [22], it has been proved that entanglement turns out to be an observer-dependent quantity in non-inertial frames, due to the Unruh radiation. The question thus arises as to how this setting is modified in the presence of mixed neutrinos, particularly in view of the discussion of Refs. [23].

As a final remark, we stress that in the current Letter we have made use of the simplest framework of neutrino mixing among two generations. The extension to three flavors is in principle straightforward, and represents one of the future directions of our investigation. We envisage how this setting is modified in the presence of mixed neutrinos, particularly in view of the discussion of Refs. [23].

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References

1. S. A. Fulling, Phys. Rev. D 7, 2850 (1973).
2. P. C. W. Davies, J. Phys. A 8, 609 (1975).
3. W. G. Unruh, Phys. Rev. D 14, 870 (1976).
4. R. Muller, Phys. Rev. D 56, 953 (1997).
5. G. E. A. Matsas and D. A. T. Vanzella, Phys. Rev. D 59, 094004 (1999); D. A. T. Vanzella and G. E. A. Matsas, Phys. Rev. D 63, 014010 (2000); D. A. T. Vanzella and G. E. A. Matsas, Phys. Rev. Lett. 87, 151301 (2001).
[6] V. A. Belinskii, B. M. Karnakov, V. D. Mur and N. B. Narozhny, Pis’ma Zh. Eksp. Teor. Fiz. 65, 861 [JETP Lett. 65, 902 (1997)]; A. M. Fedotov, V. D. Mur, N. B. Narozhny, V. A. Belinskii and B. M. Karnakov, Phys. Lett. A 254, 126 (1999); N. B. Narozhny, A. M. Fedotov, B. M. Karnakov, V. D. Mur and V. A. Belinskii, Phys. Rev. D 65, 025004 (2001); N. B. Narozhny, A. M. Fedotov, B. M. Karnakov, V. D. Mur and V. A. Belinskii, Phys. Rev. D 70, 048702 (2004).

[7] S. W. Hawking, Nature 248, 30 (1974); S. W. Hawking, Commun. Math. Phys. 43, 199 (1975) Erratum: [Commun. Math. Phys. 46, 206 (1976)].

[8] C. Barceló, S. Liberati and M. Visser, Living Rev. Rel. 14, 3 (2011); J. Drori, Y. Rosenberg, D. Bermudez, Y. Silberberg and U. Leonhardt, Phys. Rev. Lett. 122, 010404 (2019).

[9] H. Suzuki and K. Yamada, Phys. Rev. D 67, 065002 (2003).

[10] D. V. Ahluwalia, L. Labun and G. Torrieri, Eur. Phys. J. A 52, 189 (2016).

[11] M. Blasone, G. Lambiase, G. G. Luciano and L. Petruziello, Phys. Rev. D 97, 105008 (2018).

[12] G. Cozzella, S. A. Fulling, A. G. S. Landulfo, G. E. A. Matsas and D. A. T. Vanzella, Phys. Rev. D 97, 105022 (2018).

[13] N. D. Birrell and P. C. W. Davies, Quantum Fields in Curved Space, (Cambridge University Press, Cambridge, 1982).

[14] C. Itzykson and J. B. Zuber, Quantum Field Theory (McGrawHill, New York, 1980).

[15] S. M. Bilenky and B. Pontecorvo, Phys. Rept. 41, 225 (1978).

[16] M. Blasone and G. Vitiello, Annals Phys. 244, 283 (1995).

[17] M. Blasone, P. Jizba and G. Vitiello, Phys. Lett. B 517, 471 (2001).

[18] R. Carballo-Rubio, L. J. Garay, E. Martín-Martínez and J. De Ramón, [arXiv:1804.00685 [quant-ph]]; J. Marino, A. Noto and R. Passante, Phys. Rev. Lett. 113, no. 2, 020403 (2014); M. Blasone, G. Lambiase and G. G. Luciano, Phys. Rev. D 96, 025023 (2017); M. Blasone, G. Lambiase and G. G. Luciano, J. Phys. Conf. Ser. 956, 012021 (2018).

[19] F. Scardigli, M. Blasone, G. Luciano and R. Casadio, Eur. Phys. J. C 78, 728 (2018); G. G. Luciano and L. Petruziello, [arXiv:1902.07059 [hep-th]].

[20] H. S. Snyder, Phys. Rev. 71, 38 (1947); D. Amati, M. Ciafaloni and G. Veneziano, Phys. Lett. B 197, 81 (1987); M. Maggiore, Phys. Lett. B 319, 83 (1993); A. Kempf, G. Mangano and R. B. Mann, Phys. Rev. D 52, 1108 (1995); F. Scardigli, Phys. Lett. B 452, 39 (1999).

[21] M. R. Hwang, D. K. Park and E. Jung, Phys. Rev. A 83, 012111 (2011); M. Shamirzai, B. N. Esfahani and M. Soltani, Int. J. Theor. Phys. 51, 787 (2012). S. Khan, N. A. Khan and M. K. Khan, Commun. Theor. Phys. 61, 281 (2014).

[22] I. Fuentes-Schuller and R. B. Mann, Phys. Rev. Lett. 95, 120404 (2005).

[23] M. Blasone, F. Dell’Anno, S. De Siena and F. Illuminati, EPL 85, 50002 (2009); M. Blasone, F. Dell’Anno, S. De Siena, M. Di Mauro and F. Illuminati, Phys. Rev. D 77, 096002 (2008).