THERMAL AND DYNAMICAL PARTICLE CREATION
IN A CURVED GEOMETRY

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Abstract
A generalization of Termo Field Dynamics to a curved geometry is proposed. In particular a neutral scalar field minimally coupled to gravity is considered as matter content in a Robertson-Walker metric. A non linear amplification in the particle creation is obtained, due to the altogether action of thermal and geometric effects. As a consequence the frequencies in the system look like red shifted with respect to the case where the thermal creation is not taken into account.

1. Introduction
A realistic description of the dynamics of our universe must take into account all the possible interactions. But that is practically impossible. An approach that we can used is to separate the universe in system and enviroment. Then we supouse degrees of freedom with different hierarchy, some of them considered as quantum degrees of freedom and the other ones constitute the classical background. The interaction between both kinds of degrees of freedom can be produced through the changes in the spacetime which are driven by the dynamical Einstein equations and by spontaneous thermal fluctuations of the enviroment which acts as a thermal reservoir. In the cosmological semiclassical approach (see ref. [1]) the role of the classical background is played by the gravitational field. This background can also play the role of thermal reservoir, as is shown in ref. [2] and the conformal quantum thermal fluctuation of the gravitational field can be interpreted as the tilde fields of the formalism of Thermo Field Dynamics (TFD) [3]. However the tilde fields “inhabit” also in a curved geometry, then they are also affected by the changes in
the metric. So we will consider a matter field in a thermal bath, due to the radiation background, and also to the interaction with the gravitational field. Therefore we can use *Quantum Field Theory in Curved Space Time at Finite Temperature*, or in a more realistic way, instead of finite temperature, a nonequilibrium description. In order to do that it is necessary to introduce some kind of interaction between the system and the thermal bath. A simple and elegant model of that interaction is proposed in the TFD formalism [3]. The fundamental hypothesis is: *the increase of the energy in the system can be produced by the excitation of an additional quanta or by annihilation of holes of particles from the reservoir*. This is spontaneous process, as is stressed in ref. [4]. The other mechanism that increase the energy of the system is due to the creation of particles from the interaction between the scalar and gravitational fields (see ref. [1]). This interaction is introduced indirectly by means of the curved spacetime (CST). The evolution of the particle creation due to the last process is driven by the field equation. When the in-out situation is considered the Bogoliubov transformation that connect both Cauchy surfaces produces squeezed states as is shown in references [5] and [6]. In that sense the isotropic universe, taking this work into account, is another example, as the ones given in ref. [2], that relate the dynamics to the squeezed states.

The TFD hypothesis, as is shown in ref. [3] is implemented introducing an extended Fock space that includes the tilde states $|\tilde{n}>$ of the reservoir with the states $|n>$ of the system. Then we have the extended Fock space given by:

$$\{|n, \tilde{n}>, \{\tilde{n}>, \{n>\right\}$$

In particular we use for the vacuum the notation $|0> = |0, 0> = |0> |0>$, with the operators $a_k$, $\tilde{a}_k$, $a_k^\dagger$ and $\tilde{a}_k^\dagger$, such that

$$a_k|0> = 0, \quad \tilde{a}_k|0> = 0$$

$$a_k^\dagger|0> = |1_k>, \quad \tilde{a}_k^\dagger|0> = |\tilde{1}_k>, \quad \text{etc.}$$

Moreover the operators satisfy the commutation relations

$$[a_k, a_{k'}^\dagger] = \delta_{k,k'}, \quad [\tilde{a}_k, \tilde{a}_{k'}^\dagger] = \delta_{k,k'}$$

In the formulation given in ref.[7] the tilde field is a fictitious one, it does not represent any physical magnitude. In ref. [8] Israel found a parallelism between this approach and the problems where an event horizon is present, like for example the
Rindler observer and the black hole radiation. Then the role of tilde fields is played by physical but hidden modes. In ref. [2] the conformal fluctuation of the metric was proposed as tilde field. The real role of the tilde field surely must be played by a more complicated combination of the different fields present in the background. The use of one particular candidate can be interpreted as the choice a kind of coarse graining in the model of the system in a thermal bath. In the present work we can obtain interesting conclusions without specifying the nature of the tilde fields. In order to calculate statistical mean values of the physical observables, in TFD (see ref. [3]), the thermal vacuum state is introduced. If $T = T(\beta)$ is the temperature of the bath, the thermal state is defined by the transformation

$$|0, \beta > = U[a_k, \tilde{a}_k, a_k^\dagger, \tilde{a}_k^\dagger]|0>$$

(1.1)

The operator $U$ allows us to obtain the thermal vacuum as a mixture of states at zero temperature. Then the fundamental hypothesis of TFD can be imposed by the expression:

$$a|0, \beta >= \lambda \tilde{a}^\dagger|0, \beta >$$

(1.2)

In general for a system out of equilibrium, the operators are time dependent and related with the latest ones by transformations as follows (see ref. [9]):

$$a(t) = S^{-1}(t) a S(t)$$

$$\tilde{a}^\dagger(t) = S^{-1}(t) a^\dagger S(t)$$

Where the $\dagger$ symbol means that in general $(a^\dagger)^\dagger \neq a$ because $S$ is not necessarily unitary. The eqs (1.2) turn therefore to

$$a(t)|0, \beta(t) >= \lambda(t) \tilde{a}^\dagger(t)|0, \beta(t) >$$

(1.3a)

$$<0, \beta(t)|a^\dagger(t) = \gamma(t) <0, \beta(t)|\tilde{a}(t)$$

(1.3b)

We can use in the following the notation of ref. [9], i.e. : $\lambda = F$, $\gamma = fF^{-1}$. In TFD the statistical mean value of some physical observable is the thermal vacuum expectation value of the operator associated to the observable (see ref. [6]). In particular the mean value of created particles will be

$$n_{t\beta}(t) = <0, \beta|a^\dagger(t)a(t)|0, \beta >$$

(1.4)
The eqs (1.3) motivate the definition of the operators $a(\beta)$ and $a^\dagger(\beta)$ such that

$$a(\beta)|0,\beta> = 0 \quad \text{(1.5a)}$$

$$<0,\beta|a^\dagger(\beta) = 0 \quad \text{(1.5b)}$$

with $a(\beta(t)) \propto (a(t) - F\tilde{a}^\ddagger(t))$ and $a^\dagger(\beta(t)) \propto (a^\ddagger(t) - fF^{-1})$. Following refs [2] and [4] we can choose $f$ and the proportionality constant such that:

$$
\begin{pmatrix}
    a(\beta(t)) \\
    \tilde{a}^\dagger(\beta(t))
\end{pmatrix} = (1 + n_{t\beta})^{1/2} \begin{pmatrix}
    1 & -F \\
    -fF^{-1} & 1
\end{pmatrix} \begin{pmatrix}
    a(t) \\
    \tilde{a}^\ddagger(t)
\end{pmatrix}
\quad \text{(1.6)}
$$

$f$, $F$ and $n_{t\beta}$ are also time dependent functions. In order to calculate $n_{t\beta}$ by means of eq.(1.4), we need perform the inverse operation to transformation (1.6), i.e.:

$$
\begin{pmatrix}
    a(t) \\
    \tilde{a}^\ddagger(t)
\end{pmatrix} = (1 + n_{t\beta})^{-1/2}(1 - f)^{-1} \begin{pmatrix}
    1 & F \\
    fF^{-1} & 1
\end{pmatrix} \begin{pmatrix}
    a(\beta(t)) \\
    \tilde{a}^\dagger(\beta(t))
\end{pmatrix}
\quad \text{(1.7)}
$$

2. Quadrivectorial notation

It will be convenient to introduce a quadrivectorial notation, in order to include also the $a^\dagger$ and $\tilde{a}$ operators. Then we can define (see ref. [10]):

$$A_k := \begin{pmatrix}
    a_k \\
    a_k^\dagger \\
    \tilde{a}_k \\
    \tilde{a}_k^\dagger
\end{pmatrix}
\quad \text{(2.1)}$$

Let us also introduce the $4 \times 4$ matrix $\Upsilon$ so that the transformation can be written as

$$A_k(\beta(t)) = \Upsilon_k(\beta(t))A_k(t) \quad \text{(2.2a)}$$

$$A_k(t) = \Upsilon_k^{-1}(\beta(t))A_k(\beta(t)) \quad \text{(2.2b)}$$

where

$$\Upsilon_k(\beta(t)) = (1 + n_{t\beta})^{1/2} \begin{pmatrix}
    \text{I} & \text{L} \\
    \text{L} & \text{I}
\end{pmatrix}
\quad \text{(2.3)}$$
\[
\mathbf{L} = \begin{pmatrix}
0 & -F \\
-fF^{-1} & 0
\end{pmatrix}
\]
\[
\mathbf{Y}_k^{-1} = (1 + nt\beta)^{-1/2}(1 - f)^{-1}\begin{pmatrix}
\mathbf{I} & -\mathbf{L} \\
-\mathbf{L} & \mathbf{I}
\end{pmatrix}
\]  
(2.4)

Moreover \( \mathbf{I} \) is the \( 2 \times 2 \) identity matrix.

From the application of the transformation (1.7) in eq. (1.4), we obtain

\[nt\beta(1 + nt\beta) = f(1 - f)^{-2}\]

then the relation between \( nt\beta \) and \( f \) is:

\[nt\beta = \frac{f}{1 - f}\]  
(2.5)

It is interesting to note two things:

a) \( nt\beta \) does not depend on \( F \).

b) The Planckian spectrum corresponds to a particular \( f \) function; 

\[f = \exp(-\beta\epsilon)\]

with \( \epsilon \) the energy by mode.

3. **Dynamical and thermal effects in CST**

The change in time of the operators is driven by the dynamical equations. These equations, in our case, has the form of Klein-Gordon ones in a curved space-time. The main difference with flat space-time is that the coefficients in a normal mode expansion are time dependent. In order to obtain the time dependence of the annihilation-creation operators we will use the method developed by Parker [11] about which we will present a brief review. In our case the system and reservoir is considered in a curved geometry. The system is for us a massless and neutral scalar field minimally coupled to the gravitational field, given by a Robertson-Walker metric.

The action for the matter field is

\[S = \frac{1}{2} \int \sqrt{-g} d^4 x \partial_\mu \varphi \partial^\mu \varphi\]  
(3.1)

and the metric

\[ds^2 = dt^2 - a(t)^2 (dx^2 + dy^2 + dz^2)\]  
(3.2)

then the variation of the action gives the field equation

\[\nabla_\mu \partial^\mu \varphi = 0\]  
(3.3)
(with $\mu = 0, 1, 2, 3$).

In order to make the calculation easier, we can use a discretization in the same form as in [11], i.e. we can introduce the periodic boundary condition $\varphi(x + nL, t) = \varphi(x, t)$, where $n$ is a vector with integer Cartesian components and $L$ a length which goes to infinity at the end of the calculation. Then the integral $\int d^3k$ can be replaced by $(2\pi/L)^3 \sum_k$. In order to expand a general solution of the field equation we can introduce the set of functions $\{\phi_k(x)\} \cup \{\phi^*_k(x)\}$ defined by

$$\phi_k(x) = \frac{1}{(L a(t))^{3/2} \sqrt{2W}} \exp i(k \cdot x - \int_{t_0}^t W(k, t') dt'),$$  \hspace{1cm} (3.4)$$

where $W$, a real function of $k = |k|$ and $t$, has to be determined by the field equations and the boundary conditions. Then we can expand the field as

$$\varphi(x, t) = \sum_k [a_k(t) \phi_k(x) + a_k^\dagger(t) \phi^*_k(x)] = \sum_k[a_k\psi_k(x) + a_k^\dagger \psi^*_k(x)]$$  \hspace{1cm} (3.5)$$

with $a_k := a_k(t = t_1)$ when $t \geq t_1 \geq t_0$ and with (as in ref. [11])

$$\psi_k(x) = h(k, t) \exp i k \cdot x$$  \hspace{1cm} (3.6)$$

$$h(k, t) = \frac{1}{(L a(t))^{3/2} \sqrt{2W}} [\alpha(k, t)^* e^{-i \int_{t_0}^t W dt'} + \beta(k, t)^* e^{i \int_{t_0}^t W dt'}]$$  \hspace{1cm} (3.7)$$

$\alpha$ and $\beta$ are known as Bogoliubov coefficients. In order the normalization condition, holds for the base of field equation solutions, those coefficients must satisfy

$$|\alpha|^2 - |\beta|^2 = 1$$  \hspace{1cm} (3.8)$$

Eq. (3.8) is equivalent to the parametrization of the coefficients in the form:

$$\alpha(k, t) = e^{-i \gamma_{\alpha}(k,t)} \cosh \theta(k, t)$$  \hspace{1cm} (3.9a)$$

$$\beta(k, t) = e^{i \gamma_{\beta}(k,t)} \sinh \theta(k, t)$$  \hspace{1cm} (3.9b)$$

Replacing eqs (3.5)-(3.9) in the field eq. (3.3) we have the system of eqs

$$\displaystyle (1 + \cos \Gamma \tanh \theta) M + 2W \gamma_{\alpha} = 0$$  \hspace{1cm} (3.10a)$$

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\[ M \sin \Gamma + 2W \dot{\theta} = 0 \quad (3.10b) \]

where we define \( \Gamma = \gamma_\alpha + \gamma_\beta - 2 \int_{t_0}^{t} W dt' \), and

\[ M = -\frac{1}{2}(\dot{W}/W) + \frac{1}{4}(\dot{W}/W)^2 - \frac{9}{4}(\dot{a}/a)^2 - \frac{3}{2}(\dot{a}/a) + \frac{k^2}{a^2} - W^2 \quad (3.11) \]

From eqs (3.5)-(3.11) we can obtain the Bogoliubov transformation between the operators given by

\[
a_k(t) = e^{i\gamma_\alpha(k,t)} \cosh \theta(k,t) a_k + e^{i\gamma_\beta(k,t)} \sinh \theta(k,t) a_{-k}^\dagger
\]

\[
a_{-k}^\dagger(t) = e^{-i\gamma_\beta(k,t)} \sinh \theta(k,t) a_k + e^{-i\gamma_\alpha(k,t)} \cosh \theta(k,t) a_{-k}^\dagger
\]

Then we can define the mean value of created particles between the Cauchy surfaces labeled by the times 0 and \( t \), in the form

\[
n_{t,0}(t) = \langle 0|a_{k}^\dagger(t)a_k(t)|0 \rangle
\]

Using the inverse transformation of eqs (3.12) it is easier to see that

\[
n_{0,t}(t) = \langle 0, t|a_{k}^\dagger a_k|0, t \rangle
\]

Moreover from eqs (3.10) and (3.13) we obtain that \( n \) can be written in the form \( n_{0,t} = (g - 1)^{-1} \) with \( g \) a function of the particle model used, i.e.:

\[
g = \frac{\cos^2 \Gamma}{(1 + 2\frac{W}{M} \dot{\gamma}_\alpha)^2}
\]

(3.15)

The function \( g \) plays a role analogous to that of \( f^{-1} \) from eq. (2.5). For some particle model eq. (3.15) permits us to obtain Planckian spectrum.

In the quadrivectorial notation the transformation given by eqs (3.12) is

\[ A_k(t) = \Omega_k(t) A_k \quad (3.16) \]

with

\[
\Omega = \left( \begin{array}{cc} U & 0 \\ 0 & U \end{array} \right)
\]
\[
U = \begin{pmatrix}
  e^{i\gamma_\alpha(\mathbf{k},t)} \cosh \theta(\mathbf{k},t) & e^{i\gamma_\beta(\mathbf{k},t)} \sinh \theta(\mathbf{k},t) P \\
  e^{-i\gamma_\beta(\mathbf{k},t)} \sinh \theta(\mathbf{k},t) P & e^{-i\gamma_\alpha(\mathbf{k},t)} \cosh \theta(\mathbf{k},t)
\end{pmatrix}
\]

\(P\) is a parity operator, such that

\[
P a_k = a_{-k}
\]

and \(0\) is the 2x2 zero matrix. The transformation given by the matrix \(\Omega\) is unitary (see ref. [11]). Then in our case \(\dagger = \dagger\) because the dissipation is not introduced. Moreover as the dynamical change of the operators is supposed independent of the thermal one, so we have

\[
A(t, \beta(t)) = \Upsilon(\beta(t)) \Omega(t) A := \Lambda(t) A
\]

(3.17)

therefore

\[
\Lambda = (1 + n_{0\beta})^{1/2}(1 + n_{t0})^{1/2}
\begin{pmatrix}
  R & T \\
  T & R
\end{pmatrix}
\]

with

\[
R = \begin{pmatrix}
  e^{i\gamma_\alpha} & e^{i\gamma_\beta} \tanh \theta P \\
  e^{-i\gamma_\beta} \tanh \theta P & e^{-i\gamma_\alpha}
\end{pmatrix},
\]

\[
T = \begin{pmatrix}
  -Fe^{-i\gamma_\beta} \tanh \theta P & -Fe^{-i\gamma_\alpha} \\
  -fF^{-1} e^{i\gamma_\alpha} & -fF^{-1} e^{i\gamma_\beta} \tanh \theta P
\end{pmatrix}
\]

By means of the last transformation we can reobtain \(n_{t,\beta}\) in the form

\[
n_{t,\beta} = <0|a^{\dagger}(t, \beta)a(t, \beta)|0>
\]

(3.18)

It is easy to prove that also the following equality is valid:

\[
<0, t, \beta|a^{\dagger}a|0, t, \beta> = <0|a^{\dagger}(t, \beta)a(t, \beta)|0>
\]

(3.19)

Performing the calculation of \(n_{t,\beta}\), using the transformation (3.17), results (as it is shown in ref. [10]):

\[
n_{t\beta} = n_{t0} + n_{0\beta} + 2n_{t0}n_{0\beta}
\]

(3.20)

As is mentioned in ref.[10], the increment in the number of created particles, due to the temperature is equal to the one produced in curved space at zero temperature, when there is a particle distribution in the initial state, as is shown in ref.[11]. That is due to the analogous role of the operators \(\tilde{a}_k\) to the \(a_{-k}\) in curved space.
(see ref.[12]). A similar amplification effect in the particle creation was obtained in ref.[13] with a massive scalar field interacting with a massless scalar field in a Robertson-Walker metric.

4. **Thermal equilibrium**

We will see now what the effect in the thermal distribution of the field modes is, due to the presence of the thermal reservoir. In order to do that the equilibrium situation is considered. Then let us extremize the Helmholtz free energy

$$F = E - \frac{1}{\beta} K$$ (4.1)

Where $E$ is the statistical mean energy at time $t$ and temperature $T = T(\beta)$, given by

$$E = \langle 0, t, \beta | \hat{H}(t) | 0, t, \beta \rangle$$ (4.2)

whith $\hat{H}$ the hamiltonian operator. $K$ is the entropy, which is obtained as

$$K = \langle 0, t, \beta | \hat{K}(t) | 0, t, \beta \rangle$$ (4.3)

where the $\hat{K}$ operator was introduced in ref.[7] in the form

$$\hat{K} = - \sum_k \{ a^+_k a_k \log n_{t\beta}(k) - a_k a^+_k \log (1 + n_{t\beta}(k)) \}$$ (4.4)

Therefore we obtain the entropy of a Bose gas (see ref.[14]), given by

$$K = - \sum_k \{ n_{t\beta}(k) \log n_{t\beta}(k) - (1 + n_{t\beta}(k))\log (1 + n_{t\beta}(k)) \}$$ (4.5)

with $n_{t\beta}$ as in eq.(1.4).

We can calculate the energy, as in ref.[9], from the tensorial energy-momentum operator

$$\hat{T}_{\mu \nu} = \sum_{kk'} \{ a_k a_{k'} D_{\mu \nu} [\psi_k(x), \psi_{k'}(x)] + a_k a^+_k D_{\mu \nu} [\psi_k(x), \psi^*_{k'}(x)] + a^+_k a_{k'} D_{\mu \nu} [\psi^*_{k'}(x), \psi_{k'}(x)] + a^+_k a_{k'} D_{\mu \nu} [\psi^*_{k'}(x), \psi^*_{k'}(x)] + \{ k \leftrightarrow k' \} \}$$ (4.6)

with

$$D_{\mu \nu}[\varphi, \psi] = \frac{1}{2} \partial_\mu \varphi \partial_\nu \psi - \frac{1}{4} g_{\mu \nu} \partial^\sigma \varphi \partial_\sigma \psi$$
Then we can calculate the hamiltonian by

$$\hat{H} = \int a^3 d^3 x \hat{T}_{00}$$

where $\int d^3 x = L^3$ in the discrete approximation.

Instead of eq.(4.2) we can calculate the energy, in an equivalent and more easier way, as $E = \langle 0, \beta \mid \hat{H} \mid a_k(t), a_k^\dagger (t) \rangle \langle 0, \beta \rangle$, with

$$\hat{H} = \sum_{kk'} \{ a_k(t)a_{k'}(t)F_1 + a_k^\dagger (t)a_{k'}^\dagger (t)F_1^* + a_k(t)a_{k'}^\dagger (t)F_2 + a_k^\dagger (t)a_{k'}(t)F_2^* \} \quad (4.7)$$

where

$$F_1 = \frac{1}{2} \dot{\phi}_k \dot{\phi}_{k'} + \frac{1}{2} \left( \frac{k^2}{a^2} \right) \phi_k \phi_{k'} \quad (4.8a)$$

$$F_2 = \frac{1}{2} \dot{\phi}_k \dot{\phi}_k^* + \frac{1}{2} \left( \frac{k^2}{a^2} \right) \phi_k \phi_{k'}^* \quad (4.8b)$$

with $\phi = \phi[W]$ given by eq. (3.4) and $h$ the Hubble coefficient that for R-W universes is the relative velocity of the scale factor, i. e., $h = \dot{a}/a$. If we suppose that $W$ satisfies the Cauchy data that diagonalized the hamiltonian at $t$ time (see ref.[15]), the expression (4.7) turns to

$$\hat{H} = \sum_{kk'} W \left\{ \frac{1}{2} + a_{k'}^\dagger (t)a_k(t) \right\} \quad (4.9)$$

Taking into account eqs. (1.4) and (4.1) we have

$$F = \sum_k \{ W \left( \frac{1}{2} + n_{t\beta} \right) + \frac{1}{\beta} \left[ n_{t\beta}(k) log n_{t\beta}(k) - (1 + n_{t\beta}(k)) log (1 + n_{t\beta}(k)) \right] \} \quad (4.10)$$

In equilibrium

$$\delta F = 0 \quad (4.11)$$

Then the equilibrium condition is

$$\frac{\partial F}{\partial n_{t\beta}} = 0 \quad (4.12)$$

That gives us the usual Planckian distribution

$$n_{t\beta} = \frac{1}{e^{\beta W} - 1} \quad (4.13)$$
If we call $W_0$ the frequency of the normal modes in thermal equilibrium and curved spacetime, when the thermal effect, in the particle creation, is not considered, we can write in equilibrium $n_{t0} = 1/(e^{\beta W_0} - 1)$. Moreover from eq.(3.20) we have

$$n_{t\beta} = n_{t0}(1 + \frac{n_{0\beta}}{n_{t0}} + 2n_{0\beta})$$

For high temperature, when the approximation $W\beta << 1$ (idem for $W_0$) is valid, taking until the second term in the exponential expansions, gives

$$W \simeq \frac{W_0}{1 + n_{0\beta}/n_{t0} + 2n_{0\beta}} \quad (4.14)$$

In the denominator of eq. (4.14) all the terms are positive, therefore the frequency $W$ looks like red shifted with respect to $W_0$.

5. Conclusions and Comments

The particle creation due to the thermal bath and the one coming from the interaction with the geometry, both mechanisms of particle creation as we can see from eq. (3.20), act symmetrically, however the nature of both processes is different. One of them is a spontaneous process due to the thermal fluctuations of the reservoir, while the other one is produced by the dynamical change of geometry. The symmetry is due to the fact that in the two situations the vacuum state changes by means of Bogoliubov transformations, because still for the fields in the curved geometry the field equation is Klein-Gordon like. The similarity of the two transformations was analized in ref. [16]. There we have seen that the isotropy of the spacetime produces a mirror symmetry between the modes of the field and one of these symmetrical spaces can be associated with the tilde field of TFD. In a more complex case, when there is interaction with other fields, additional amplification of the particle production can be present (see ref. [13]), producing a “cascade” effect.

From eq.(4.14) we can see that the observed frequency of the system with reservoir, with respect to the one at zero temperature looks like red shifted due to the effect of thermal creation of particles (or modes of the field).

The calculation is identical if the field is massive, but the interpretation can be different, because the thermal effect changes the redistribution of modes of the field by the interaction with the reservoir. Then the particles obtain quanta of energy of the medium, but no new particle appears. In the massless case the simple creation of modes can be interpreted as new massless particles. In a flat spacetime it is necessary to introduce some symmetry breaking mechanism, in order to produce massive particles from the vacuum (see ref. [17]). However in curved spacetime
massive particles are created without an explicit breaking of symmetry. This can be interpreted using the analogy between the vacua in the curved spacetime and the ones related to accelerated observers in flat spacetime (it is known in the literature [1],[18], as Rindler observer). As it is shown in ref. [19] the Rindler observer has a Planckian spectrum in his reference system while in the Minkowski vacuum particles are not observed. It is similar to the behaviour of the electromagnetic field with respect to different reference systems. Actually we have treated with a toy model of electromagnetic the field, represented in our case by the massless scalar field.

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