What happens once an accelerating observer has detected a Rindler particle?

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In a seminal paper, Unruh and Wald showed that the detection of a Rindler particle by a linearly uniformly accelerated detector coupled to a Klein-Gordon field in the Minkowski vacuum corresponds to the creation of a Minkowski particle from the inertial viewpoint. In this paper, we revisit the situation studied in that work, but now consider, addition what happens once the detector has been excited somewhere along its world line. From an orthodox point of view, the change in the state of field induced by the measurement is nonlocal and affects both in the left and right Rindler wedges. If one relies on a semiclassical treatment of gravity and takes it to be applicable to this context, this situation seems to open the possibility for designing superluminal communication protocols between two spacelike separated observers confined respectively to the right and left Rindler wedges respectively. We discuss the possible ways in which physics could prevent such measurement-induced, faster-than-light signaling protocols.

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I. INTRODUCTION

The Unruh effect is often succinctly described as the fact that from the point of view of a uniformly accelerated observer the Minkwoski vacuum looks like a thermal state at a temperature proportional to their proper acceleration [1–3]. This is the Unruh temperature, given by $T_{\mathrm{U}} = a/(2\pi)$ in natural units. The study of the Unruh effect continues to be of great interest in theoretical physics, see, e.g., [4] for a study of the problem as return to equilibrium, [5] for the characterization of thermalization time for the Unruh effect, [6] in the context of entanglement harvesting. The intimate relationship between the Unruh effect and Hawking radiation by black holes in equilibrium, which indeed originally motivated Unruh's work [3], is by now well established (see, e.g., [[7], Chap. 5]), and can be seen very explicitly in two-dimensional situations [8]. References [9,10] include thorough discussions of the Unruh effect, including applications and structural properties. Very recently a concrete experimental proposal has been put forward in [11] to detect for the first time the analogous Unruh temperature along uniformly accelerated circular motions. This is an analogous circular Unruh effect, see, e.g., [12–14].

Sitting at the heart of the Unruh effect is the fact that the Minkowski vacuum state restricted to, say, the right Rindler wedge of Minkowski spacetime, $|t| < x$, can be formally represented as thermal mixture of so-called *Fulling-Rindler particles* supported on the right Rindler wedge. These are nothing but particles defined in the Fulling-Rindler quantization in flat spacetime, for which the notion of positive energy is defined with respect to Lorentz boosts, $b^a = a(x\partial_t^a + t\partial_x^a)$, which generate the natural notion of time evolution for linearly accelerated observers.

Following this observation, in 1984 Unruh and Wald wrote a seminal paper [15] where they clarified what occurs when a linearly uniformly accelerated observer detects a Rindler particle: From the point of view of an inertial observer in Minkowski spacetime, the absorption of a Rindler particle—modeled as a two-level detector excitation—corresponds to the emission of a Minkowski particle. The paper [15] is remarkable in that not only did it illustrate the relativity of the notion of a particle, detection and emission, but clarified the fact that working in terms of quantum fields (and taking the notion of particle to have a *contextual* and relative meaning)

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is fully consistent with the basic ideas underlying the equivalence principle.

Furthermore, [15] has served as the starting ground for further developments. For example, the study of *bremsstrahlung* as seen from the point of view of accelerated observers [16,17], the analysis of the decay of accelerated protons, and the finding that such behavior approaches that of accelerated neutrons, as the mass scale characterizing that acceleration—i.e., the corresponding Unruh temperature—increases, and disappears exponentially as that quantity grows beyond the value of the proton-neutron mass gap [18,19]. More recently, it has been studied how linearly uniformly accelerated atoms can produce a squeezed entangled pair of photons, which are predominantly "localized" in opposite Rindler wedges, by becoming excited in [20]. In that same work the analogous phenomenon in black holes spacetimes is studied, which continues to draw parallels between the Unruh and Hawking effects, but furthermore gives stimulating perspectives in the field of analog gravity.

In any case, there are three central issues that are addressed in [15]. The first one is to analyze the unitary evolution of the joint field-detector system when the field is initially in the Minkowski vacuum state and the two-level detector, initially switched off and prepared in the ground state, follows a linearly uniformly accelerated trajectory in the right Rindler wedge. This is done using perturbation theory in the interaction picture up to first order. (Second order contributions were further studied in [21].) The second question is to see what the updated state of the field is, assuming the detector has in fact detected a Rindler particle after some interaction time has elapsed, namely a one-particle state in the Minkowski folium, and to obtain the updated stress-energy tensor. It is found that, since the updated field state is a one-particle state, the energy of the field has increased upon detection. The third question addressed in [15] is whether detecting Rindler particles can be used as a mechanism for extracting an unbounded amount of energy from the field or to send a superluminal signal from the right to the left Rindler wedges. In both cases, the analysis leads to a negative answer.

The motivation of the present work is two-fold:

(i) First, we wish to revisit the three central questions discussed in [15]. Concerning the first one, we note that the calculation in [15] is performed by exploiting an analogy with the situation of a detector interacting with a field in a thermal state that describes a *proper mixture* [22], i.e., such that the actual state of the system is pure, but there is a degree of ignorance as to what the state of the system actually is, which is encoded in weights accounting for a probability distribution of the possible (pure) states the system might be in. This results in a mixed state description of a pure state due to ignorance. On the other hand, the most natural description for the Minkowski vacuum from the point of view of an accelerated observer is that of a thermal state as an *improper mixture* [22] (see footnote 2 for more details), as the left Rindler wedge degrees of freedom must be traced out, yielding a reduced mixed state. In Sec. II we will carry out the calculation from the improper-mixture viewpoint. While the results coincide mathematically, as they should, we think that this treatment is conceptually clearer.

We then proceed to calculate the updated expectation value of the stress-energy tensor once the detector has clicked, which ties in with the second central question in [15]. We obtain expressions in both the right and the left Rindler wedge, adding to the result displayed in [15] for the right Rindler wedge, as we show in Sec. II. Concerning the third central question addressed in [15], on the point of energy extraction, we agree with the no-go argument presented by Unruh and Wald: while the energy is *not conserved* for a single measurement, it is conserved on average for very many successive measurements. We then address the question of superluminal communication, which ties in with the second motivation of this paper.

(ii) Here we shall raise the point that there is a potential issue after a single measurement has been carried out, if one is to trust the semiclassical approximation of quantum gravity "before" and "after" the measurement has been performed and assumes that the state measurement induces a state collapse along a Cauchy surface in spacetime. The point will be that in semiclassical gravity the expectation value of the stress-energy tensor can be used to actually source geometry, see Eq. (35) below. Thus, an abrupt change of this quantity along a Cauchy surface in this setting might be expected to be detectable by an experiment on the gravitational sector. Where and how this abrupt change occurs, i.e., where and how the state of the field can be seen as collapsing upon a measurement of the detector is most likely to play a role on how to prevent this apparent paradox from occurring, but our current understanding of these questions is fuzzy—hence the use of quotation marks around the words *before* and *after*. Thus, it seems to us that the resolution of this puzzling situation is likely connected with the full resolution of *measurement problem* of quantum theory. We will in fact offer at the end what we think is a rather exhaustive list of possibilities for preventing such superluminal signals.

On this point we wish to emphasize that Unruh and Wald point out in their discussion in [[15], Sec. IV] that the presence of a detector (switched on or otherwise) in the right Rindler wedge has no influence on the left Rindler wedge. Their argument is based on the Heisenberg-picture observation that the effects of the detector can only affect the causal future of the coupling region between the detector and the field. In fact, this observation does not even depend on the details of the detector or the field observables, see, e.g., [23] for a precise statement in some generality. The limitation of that argument is that it remains fuzzy as to the what actual measurement the detector does, which is typically described as a projection onto the out-state in the interaction picture. This is a central difference between the above mentioned work and the posture explored in this paper.

We should mention that in the context of nonrelativistic quantum mechanics the no-signaling theorem shows that entanglement between two separated systems cannot be used for superluminal signaling. This result has, at this time, no counterpart in QFT where how to deal with measurement processes, is under development. Furthermore, in this case we will be considering the problem within the semiclassical context for the treatment of gravitation. Moreover the no-signaling theorem involving joint measurements or manipulations made at one "time" on both components of the entangled system. The situation envisioned in this work concerns, as we will see, waiting arbitrarily large times for a detector to get excited and also waiting arbitrarily large times for the manifestation of an effect on the other side.

The organization of this paper is as follows: Concerning the paper's first motivation, in Sec. II we describe the evolution of the field-detector system using a left and right doubled Fock space representation for the field, and we obtain the stress-energy tensor after a Rindler particle has been measured. In doing so, we do not make any assumptions on the details of the coupling of the detector and the field, in particular we do not assume a long-time limit for the interaction, other than assuming that the coupling is week, which allows us to conform ourselves with first-order effects in the coupling. We then discuss the non-conservation of energy upon measurements in Sec. IV in a simplified setting, for the sake of clarity. We then proceed to study our second motivation. The possibility of faster than light signaling, its implications and potential paths for their avoidance appear in Sec. V. Discussions and conclusions appear in Sec. VI.

II. WHAT HAPPENS ONCE AN ACCELERATING OBSERVER HAS DETECTED A RINDLER PARTICLE?

Consider as in [15] a particle detector coupled to a Klein-Gordon field in Minkowski spacetime following a linearly uniformly accelerated trajectory with acceleration a in the right Rindler wedge. In other words, the particle detector follows the integral curve generated by the boost $b^a = a(x\partial_t^a + t\partial_x^a)$. While currently the pointlike Unruh-DeWitt detector [5,24,25] is the most prominent detector model used in studies about the relativistic quantum information and QFT in curved spacetime literature, we shall model our detector as Unruh and Wald have in [15] to stay closer to their original treatment.

The detector is a two-level system with Hilbert space $\mathbb{C}^2$ spanned by energy eigenstates $|\uparrow\rangle$ and $|\downarrow\rangle$. The detector Hamiltonian is $\hat{H}_{\mathrm{D}} \coloneqq \Omega\hat{A}^*\hat{A}$, where $\hat{A}$ and $\hat{A}^*$ are raising and lowering operators, and $\Omega > 0$ is the energy of the excited state, i.e., $\hat{H}_{\mathrm{D}}|\uparrow\rangle = \Omega|\uparrow\rangle$ and $\hat{H}_{\mathrm{D}}|\downarrow\rangle = 0$.

The coupled detector-field theory is described by the interaction Hamiltonian

$$\hat{H} = \hat{H}_{\mathrm{D}} \otimes \mathbb{1} + \mathbb{1}_{\mathrm{D}} \otimes \hat{H}_{\Phi} + \hat{H}_I, \tag{1}$$

where $\hat{H}_{\Phi}$ is the Klein-Gordon Hamiltonian and the interaction Hamiltonian is defined by

$$\hat{H}_I(\tau) = \epsilon(\tau)\int_{\Sigma} e^{2a\xi}\mathrm{d}\xi\mathrm{d}y\mathrm{d}z\Big[\psi(\xi,y,z)\hat{A}(\tau) + \bar{\psi}(\xi,y,z)\hat{A}^*(\tau)\Big] \otimes \hat{\Phi}(\tau,\xi,y,z), \tag{2}$$

where $\hat{\Phi}$ is the Klein-Gordon field, $\psi \in C_0^\infty(\Sigma)$ defines the profile of the spatial extension of the detector and $\epsilon \in C_0^\infty(\mathbb{R})$ is a switching function that controls the interaction between the detector and the Klein-Gordon field along the linearly uniformly accelerated trajectory of the detector. We shall assume that the interaction between the detector and the field is weaker than any other scale in the problem and that it takes place for sufficiently long times.

The coupling takes place in the right Rindler wedge, where the flat metric can be written in terms of Rindler coordinates

$$t = \frac{1}{a}e^{a\xi}\sinh(a\tau), \qquad x = \frac{1}{a}e^{a\xi}\cosh(a\tau). \tag{3}$$

It takes the form

$$\mathrm{d}s^2 = -e^{2a\xi}(\mathrm{d}\tau^2 - \mathrm{d}\xi^2) + \mathrm{d}y^2 + \mathrm{d}z^2. \tag{4}$$

Furthermore, in the right Rindler wedge, the Klein-Gordon field can be written as

$$\hat{\Phi}(\tau,\xi,y,z) = \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3\kappa(v_{I\vec{\kappa}}(\tau,\xi,y,z)\hat{a}_{\mathrm{R}\vec{\kappa}} + \overline{v_{I\vec{\kappa}}}(\tau,\xi,y,z)\hat{a}^*_{\mathrm{R}\vec{\kappa}}), \tag{5}$$

with $\vec{\kappa} = (\omega,\kappa_y,\kappa_z)$, and where the right Rindler modes can be written in terms of the modified Bessel function of the second kind or MacDonald's function,

$$v_{I\vec{\kappa}}(x) = \sqrt{\frac{\sinh(\frac{\pi\omega}{a})}{4\pi^4 a}}K_{i\omega/a}\left[\frac{\sqrt{\kappa_y^2+\kappa_y^2+m^2}}{a}e^{a\xi}\right] \times e^{-i\omega\tau + i(y\kappa_y + z\kappa_z)} \tag{6}$$

and the formal sharp-momentum annihilation and creation operators are $\hat{a}_{\mathrm{R}\vec{\kappa}} := \hat{a}(\overline{v_{\mathrm{R}\vec{\kappa}}})$ and $\hat{a}^*_{\mathrm{R}\vec{\kappa}} := \hat{a}^*(v_{\mathrm{R}\vec{\kappa}})$ respectively. The annihilation operator annihilates the right Fulling-Rindler vacuum, Ω_{R}, while creation operators create Rindler particles.

A fully analogous description of the quantum theory holds in the left Rindler wedge. Introducing the left Rindler coordinates

$$t = \frac{1}{a} e^{a\tilde{\xi}} \sinh(a\tilde{\tau}), \qquad -x = \frac{1}{a} e^{a\tilde{\xi}} \cosh(a\tilde{\tau}) \tag{7}$$

the field in the left Rindler wedge takes the analogous form

$$\hat{\Phi}(\tilde{\tau}, \tilde{\xi}, y, z) = \int_{\mathbb{R}^+ \times \mathbb{R}^2} \mathrm{d}^3\kappa \Big(v_{I\vec{\kappa}}(\tilde{\tau}, \tilde{\xi}, y, z) \hat{a}_{\mathrm{L}\vec{\kappa}} + \overline{v_{I\vec{\kappa}}}(\tilde{\tau}, \tilde{\xi}, y, z) \hat{a}^*_{\mathrm{L}\vec{\kappa}} \Big), \tag{8}$$

where the left Rindler modes have an identical form to the right modes (6) upon the replacement of τ and ξ by $\tilde{\tau}$ and $\tilde{\xi}$. It is very well known that the Minkowski vacuum restricted to the (right or left) Rindler wedges looks like a thermal mixture of (right or left, resp.) Rindler particles. See Appendix A for details.

In any case, in the case at hand we consider that the state of the system is prepared initially (before the switch-on of ϵ) as the tensor product

$$|s_{-\infty}\rangle = |{\downarrow}\rangle \otimes |\Omega_{\mathrm{M}}\rangle. \tag{9}$$

In the interaction picture, the late-time state of the system (after the switch-off of ϵ) is given by

$$\begin{aligned} |s_{\infty}\rangle &= T\left(\mathrm{e}^{-\mathrm{i}\int_{\mathbb{R}} \mathrm{d}\tau \hat{H}_I}\right) |s_{-\infty}\rangle \\ &= |s_{-\infty}\rangle + \left[-\mathrm{i}\int_{\mathbb{R}} \mathrm{d}\tau \hat{H}_I + O(\epsilon^2)\right] |s_{-\infty}\rangle, \end{aligned} \tag{10}$$

under the assumption that the coupling is weak. To first order in perturbation theory, the late-time state of the system is

$$\begin{aligned} |s_{\infty}\rangle = |{\downarrow}\rangle \otimes |\Omega_{\mathrm{M}}\rangle - i|{\uparrow}\rangle \otimes \int_{\mathbb{R}^+ \times \mathbb{R}^2} \mathrm{d}^3\kappa \int_I \mathrm{dvol}(x) \zeta(x) \\ \times \left(v_{I\vec{\kappa}}(x) \hat{a}_{\mathrm{R}\vec{\kappa}} + \overline{v_{I\vec{\kappa}}(x)} \hat{a}^*_{\mathrm{R}\vec{\kappa}} \right) |\Omega_{\mathrm{M}}\rangle, \end{aligned} \tag{11}$$

where the volume element in the right Rindler wedge is locally $\mathrm{dvol}(x) = e^{2a\xi} \mathrm{d}\tau \mathrm{d}\xi \mathrm{d}y \mathrm{d}z$ and

$$\zeta(x) := e^{i\Omega\tau} \varepsilon(\tau) \overline{\psi(\xi, y, z)}. \tag{12}$$

In [15] the assumption has been made that $\varepsilon(\tau)$ is nearly constant, physically representing a long interaction between the detector and the field, such that switching effects are negligible. In this case, the τ integral can be performed directly and one obtains a factor proportional to $\delta(\Omega - \omega)$ that indicates that only the modes $v_{I\vec{\kappa}}$ with frequency highly localized around Ω will contribute to first order in Eq. (10). This makes perfect sense, after a long interaction time the only modes that get excited from the vacuum state are those whose energy coincides with the frequency gap of the detector.

Let us however digress at this point and not make this approximation. The reason is that in our case we are interested in postmeasurement effects for measurements carried out after a finite time of interaction between the field and the detector. On the assumption that the detector clicks, the updated state for the field becomes

$$\begin{aligned} |f\rangle = -i\mathcal{N} \int_{\mathbb{R}^+ \times \mathbb{R}^2} \mathrm{d}^3\kappa \int_I \mathrm{dvol}(x) \zeta(x) \\ \times \left(v_{I\vec{\kappa}}(x) \hat{a}_{\mathrm{R}\vec{\kappa}} + \overline{v_{I\vec{\kappa}}(x)} \hat{a}^*_{\mathrm{R}\vec{\kappa}} \right) |\Omega_{\mathrm{M}}\rangle, \end{aligned} \tag{13}$$

where the normalization $\mathcal{N}$ is given by

$$\begin{aligned} \mathcal{N} = \Bigg(\int_{\mathbb{R}^+ \times \mathbb{R}^2} \mathrm{d}^3\kappa \int_I \mathrm{dvol}(x) \int_I \mathrm{dvol}(x') \overline{\zeta(x)} \zeta(x') \\ \times \left(\frac{v_{I\vec{\kappa}}(x) \overline{v_{I\vec{\kappa}'}(x')}}{1 - e^{-2\pi\omega/a}} + \frac{\overline{v_{I\vec{\kappa}}(x)} v_{I\vec{\kappa}'}(x')}{e^{2\pi\omega/a} - 1} \right) \Bigg)^{-1/2} \end{aligned} \tag{14}$$

as can be seen in Appendix B.

We are interesting in the change in the expectation value of the stress-energy tensor of the field in the updated state, i.e., we are interested in

$$\Delta T_{ab} := \langle f | \hat{T}_{ab}(x) f \rangle - \langle \Omega_{\mathrm{M}} | \hat{T}_{ab} \Omega_{\mathrm{M}} \rangle \tag{15}$$

in the right and left Rindler wedges. (Note that imposing $\langle \Omega_{\mathrm{M}} | \hat{T}_{ab} \Omega_{\mathrm{M}} \rangle = 0$, we have that $\Delta T_{ab} = \langle f | \hat{T}_{ab}(x) f \rangle$.) To this end, if we use a point-splitting prescription for renormalizing the stress-energy tensor the object of interest is to obtain the two-point function in the left and right Rindler wedges. It follows from the calculations in appendices C and D that the two-point function in the updated state takes the form

$$\langle f | \hat{\Phi}(x) \hat{\Phi}(x') f \rangle = \langle \Omega_{\mathrm{M}} | \hat{\Phi}(x) \hat{\Phi}(x') \Omega_{\mathrm{M}} \rangle + \Delta_{\mathrm{R}}(x, x') \quad \text{in the right Rindler wedge and,} \tag{16}$$

$$\langle f | \hat{\Phi}(x) \hat{\Phi}(x') f \rangle = \langle \Omega_{\mathrm{M}} | \hat{\Phi}(x) \hat{\Phi}(x') \Omega_{\mathrm{M}} \rangle + \Delta_{\mathrm{L}}(x, x') \quad \text{in the left Rindler wedge.} \tag{17}$$

Here, Δ_{R} and Δ_{L} are real, smooth, symmetric bi-functions given by

$$\Delta_{\mathrm{R}}(x,x') = \mathcal{N}^2 \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3\kappa \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3 p \int_I \mathrm{dvol}(y) \int_I \mathrm{dvol}(y') \frac{\overline{\zeta(y)}\zeta(y')}{(1-e^{-2\pi\omega_p/a})(1-e^{-2\pi\omega_\kappa/a})}$$
$$\times \Big(v_{I\vec{p}}(y) v_{I\vec{\kappa}}(x) \overline{v_{I\vec{p}}(x')}\,\overline{v_{I\vec{\kappa}}(y')} + \overline{v_{I\vec{\kappa}}(y)} v_{I\vec{\kappa}}(x) v_{I\vec{p}}(x') \overline{v_{I\vec{p}}(y')} e^{-2\pi\omega_\kappa}$$
$$+ \overline{v_{I\vec{p}}(y)} v_{I\vec{\kappa}}(x) v_{I\vec{p}}(x') \overline{v_{I\vec{\kappa}}(y')} e^{-2\pi\omega_p} + \overline{v_{I\vec{\kappa}}(y)} v_{I\vec{\kappa}}(x) \overline{v_{I\vec{p}}(x')} v_{I\vec{p}}(y') e^{-2\pi\omega_\kappa} e^{-2\pi\omega_p} \Big) + \text{c.c.}, \tag{18}$$

$$\Delta_{\mathrm{L}}(x,x') = \mathcal{N}^2 \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3\kappa \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3 p \int_I \mathrm{dvol}(y) \int_I \mathrm{dvol}(y') \overline{\zeta(y)}\zeta(y') \frac{e^{-\pi\omega_p/a} e^{-\pi\omega_\kappa/a}}{(1-e^{-2\pi\omega_p/a})(1-e^{-2\pi\omega_\kappa/a})}$$
$$\times \Big(v_{I\vec{p}}(y) v_{II\vec{\kappa}}(x) \tilde{v}_{II\vec{p}}(x') \tilde{v}_{I\vec{\kappa}}(y') + v_{I\vec{p}}(y) \tilde{v}_{II\vec{p}}(x) v_{II\vec{\kappa}}(x') \tilde{v}_{I\vec{\kappa}}(y')$$
$$+ v_{I\vec{p}}(y) \overline{v_{II\vec{\kappa}}(x)} \tilde{v}_{II\vec{p}}(x') \overline{\tilde{v}_{I\vec{\kappa}}(y')} + \overline{v_{I\vec{p}}(y)}\, \overline{\tilde{v}_{II\vec{p}}(x)}\, v_{II\vec{\kappa}}(x') \tilde{v}_{I\vec{\kappa}}(y') \Big) + \text{c.c.}, \tag{19}$$

where in Eq. (19) the tilde on the modes denotes a parity operator in the orthogonal direction to the Rindler wedges, i.e., $\tilde{v}_{I\vec{\kappa}} \coloneqq v_{I(\omega_\kappa,-\kappa_\perp)}$ and similarly for $\tilde{v}_{II\vec{\kappa}}$.

Let us now compute the expectation value of the stress-energy tensor in the state $|f\rangle$. Note that a difference with the calculation in [15] is that we do not treat the state of the system as an improper thermal mixture in the right Rindler wedge, but rather as a pure state defined in the right and left Rindler wedges, which allows us to obtain the renormalized expectation value of the stress energy tensor for points in the right and left Rindler wedges.

Here we use the terminology introduced by d'Espagnat [22] to clarify that the same mathematical object, a density matrix, can be used to represent two very different physical situations: (1) The case in which one is interested in an ensemble of identical quantum systems, each one of which is in one pure and definite quantum state among a list possible such states $\{|i\rangle\}$ and where the fraction of such states in the ensemble is given by a certain classical distribution function $f(i)$, and (2) the case in which a system of interest S is a subsystem of a larger system $S+E$ with the latter in a given pure quantum state, but with our interest focused just on S which can therefore be characterized in terms of the reduced density matrix obtained after tracing over the degrees of freedom of E. For the first case one reserves the name "proper mixture" and says the density matrix represents such proper mixture, and for the second case one reserves the name "improper mixture," and equally indicates the density matrix is to be understood as representing the improper mixture. We note that another situation one might want to consider, is one in which one is dealing with a single quantum system S which is in a pure state, which however is not completely known, and for which one has information about the classical probability $p(i)$ of the system being in each one of the quantum states. For such situation one can often use the characterization provided by case (1) by considering a corresponding imaginary ensemble, in which the fraction is arranged to match the given probability i.e., $f(i) = p(i)$. That situation one also talks by extension about a proper mixture, even though the state of the system is pure, and thus its "properness" (or the fact that we do not express the state as a pure one) is just a result of our ignorance. Finally, as is usual, a density matrix is characterized as thermal if its representation in the energy basis has the standard thermal weights. Thus a thermal density matrix can be proper or improper.

III. THE CHANGE IN THE STRESS-ENERGY TENSOR

In order to compute the stress-energy tensor in the updated state, $\langle f|\hat{T}_{ab}(x)f\rangle$, in the right/left wedge we apply the point-splitting operator

$$\mathcal{T}_{ab} \coloneqq g_b{}^{b'} \nabla_a \nabla_{b'} - \frac{1}{2} g_{ab} g^{cd'} \nabla_c \nabla_{d'} - \frac{1}{2} g_{ab} m^2, \tag{20}$$

where $g_a{}^{a'}$ is the parallel-transport propagator, to $\Delta_{\mathrm{R/L}}(x,x')$ given by Eq. (C12), and then take the coincidence limit as

$$\langle f|\hat{T}_{ab}(x)f\rangle = \lim_{x'\to x} \mathcal{T}_{ab} \Delta_{\mathrm{R/L}}(x,x'). \tag{21}$$

It is useful to locally express the parallel-transport components in Rindler coordinates by using the formula $g_\mu{}^{\mu'}(x,x') = e^I_\mu(x) e^{\mu'}_I(x')$ in terms of the soldering form and its inverse. We then have that in the right Rindler wedge the parallel-transport propagator has components $g_\eta{}^{\eta'}(x,x') = e^{a\xi} e^{-a\xi'}$, $g_\xi{}^{\xi'}(x,x') = e^{a\xi} e^{-a\xi'}$, $g_y{}^{y'} = 1$ and $g_z{}^{z'} = 1$, with all other components vanishing.

We are chiefly concerned with changes in the left Rindler wedge, which is causally disconnected from the detector that clicks. Inserting (19) into Eq. (21) we have that the in the left Rindler wedge stress-energy tensor in the updated state takes a diagonal form and the form it takes can be read directly from (21). For instance, for the energy density of a massless field we have in the left Rindler wedge

$$\begin{aligned}\langle f|\hat{T}_{\eta\eta}(x)f\rangle = \Re\mathcal{N}^2 \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3\kappa \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3 p \int_I \mathrm{dvol}(y) \int_I \mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y') \frac{e^{-\pi\omega_p/a} e^{-\pi\omega_\kappa/a}}{(1-e^{-2\pi\omega_p/a})(1-e^{-2\pi\omega_\kappa/a})} \\ \times \Bigg(v_{I\vec{p}}(y)\partial_\eta v_{II\vec{\kappa}}(x)\partial_\eta \tilde{v}_{II\vec{p}}(x)\tilde{v}_{I\vec{\kappa}}(y') + e^{2a\xi}\sum_{i=1}^{3} g^{ii} v_{I\vec{p}}(y)\partial_i v_{II\vec{\kappa}}(x)\partial_i \tilde{v}_{II\vec{p}}(x)\tilde{v}_{I\vec{\kappa}}(y') \\ + v_{I\vec{p}}(y)\partial_\eta \tilde{v}_{II\vec{p}}(x)\partial_\eta v_{II\vec{\kappa}}(x)\tilde{v}_{I\vec{\kappa}}(y') + e^{2a\xi}\sum_{i}^{3} g^{ii} v_{I\vec{p}}(y)\partial_i \tilde{v}_{II\vec{p}}(x)\partial_i v_{II\vec{\kappa}}(x)\tilde{v}_{I\vec{\kappa}}(y') \\ + v_{I\vec{p}}(y)\partial_\eta \overline{v_{II\vec{\kappa}}(x)}\partial_\eta \tilde{v}_{II\vec{p}}(x)\overline{\tilde{v}_{I\vec{\kappa}}(y')} + e^{2a\xi}\sum_{i}^{3} g^{ii} v_{I\vec{p}}(y)\partial_i \overline{v_{II\vec{\kappa}}(x)}\partial_i \tilde{v}_{II\vec{p}}(x)\overline{\tilde{v}_{I\vec{\kappa}}(y')} \\ + \overline{v_{I\vec{p}}(y)}\partial_\eta \overline{\tilde{v}_{II\vec{p}}(x)}\partial_\eta v_{II\vec{\kappa}}(x)\tilde{v}_{I\vec{\kappa}}(y') + e^{2a\xi}\sum_{i}^{3} g^{ii} \overline{v_{I\vec{p}}(y)}\partial_i \overline{\tilde{v}_{II\vec{p}}(x)}\partial_i v_{II\vec{\kappa}}(x)\tilde{v}_{I\vec{\kappa}}(y')\Bigg),\end{aligned} \tag{22}$$

where $\Re$ denotes the real part and with $g^{11} = e^{-2a\xi}$ and $g^{22} = g^{33} = 1$. In the massive case, one adds the term

$$\frac{1}{2} m^2 e^{2a\xi} \Delta_{\mathrm{L}}(x,x) \tag{23}$$

to the right-hand side of Eq. (22).

Similar expressions can be obtained for the components $\langle f|\hat{T}_{\xi\xi}(x)f\rangle$, $\langle f|\hat{T}_{yy}(x)f\rangle$, and $\langle f|\hat{T}_{zz}(x)f\rangle$ in the left Rindler wedge, and for the four nonvanishing components in the right Rindler wedge.

Instead of spelling out in detail all of the remaining components, we point out that Eq. (22) suffices to a central point of the paper, which is that a measurement that occurs in the right Rindler wedge has nontrivial effects on the causally disconnected left Rindler wedge. It is natural to ask "when" or "where" in spacetime the state collapses after a detector has detected a Rindler particle. This is far from obvious, but a natural assumption seems to be that the state collapses along a Cauchy surface of spacetime (see [26]) intersecting the "detection event" on the right Rindler wedge and extending into the left Rindler wedge.

We should emphasize that regardless of "how big" this change might be, it is a principled statement that by including state collapses with semiclassical gravity has produced an abrupt change in the a region causally disconnected from where the measurement took place, in this case by means of a detector click.

One can see by direct inspection of (22) and (23) that for high accelerations there exists spacetime regions in the left Rindler wedge where the expectation value of the energy density becomes large. For example, for sufficiently small $\tilde{\xi}$, say $\tilde{\xi} \sim 1/a$, the Rindler modes do not exhibit the large-argument suppression of the MacDonald function, but the integrand factors $e^{-\pi\omega/a}/(1-e^{-2\pi\omega/a})$ exhibit the behavior

$$\frac{e^{-\pi\omega/a}}{1-e^{-2\pi\omega/a}} = \frac{1}{2\sinh(\pi\omega/a)} = \frac{a}{2\omega} + O(\omega/a). \tag{24}$$

This observation is consistent with what one would expect in the long interaction time limit case presented in [15], as can be seen from Eq. (3.29) in that paper in the small $\beta = 2\pi/a$ regime.

IV. ENERGETIC CONSIDERATIONS

In this section we revisit some of the energetic considerations discussed in [15] but focusing on the various possible *individual outcomes* of the "detection attempts," rather than on the ensemble averages of detector measurement outcomes, which are the quantities considered in most of the discussion of said reference.

Following [15] and to simplify the discussion we will consider the case of (ensembles of) harmonic oscillators, and entangled pairs of harmonic oscillators instead of quantum fields. There is no loss of conceptual clarity in doing so, but the treatment is mathematically less involved.

Consider a harmonic oscillator with energy eigenstates $\{|n\rangle\}$ with renormalized energies $n\epsilon$ (i.e., we removed the zero point or ground state energy for simplicity of analysis), and a detector with two states $|\downarrow\rangle$ and $|\uparrow\rangle$ with energy levels 0 and ϵ respectively. Let us assume that the initial state of the combined system is

$$|\Psi\rangle = N\left(|0\rangle \otimes |\downarrow\rangle + \frac{1}{\sqrt{\alpha}}|n\rangle \otimes |\uparrow\rangle\right), \tag{25}$$

with $N^2 = 1/(1+\alpha^{-1})$. The expectation value of the energy is $\langle E\rangle_\Psi = \frac{n\epsilon}{1+\alpha}$.

If an observer finds the detector in the unexcited state $|\downarrow\rangle$, the value of the energy becomes $\langle E\rangle_{\mathrm{unex}} = 0$, and the probability for this is $\mathrm{P}_{\mathrm{unex}} = N^2$. On the other hand, if they find the detector in the excited state, the energy becomes $\langle E\rangle_{\mathrm{ex}} = n\epsilon$ and the probability for this is $\mathrm{P}_{\mathrm{ex}} = N^2/\alpha$. The average result is of course $\langle E\rangle = \frac{n\epsilon}{1+\alpha}$, which is the same as $\langle E\rangle_\Psi$. We note however that in each specific case (i.e., for each specific outcome of the observation) the actual value of the energy differs from $\langle E\rangle_\Psi$, the expectation value of the initial state of the combined system. This is of course not surprising, given that the state $|\Psi\rangle$ is not an eigenstate of the total energy operator.

As has been argued in [27], the relevant issue regarding energy conservation is not its conservation "on average," but its conservation on each single individual instance of any experiment. In this sense the possibility of preparing a state such as $|\Psi\rangle$ and to make the observations described above already set serious doubts about the general validity of anything like a law of "energy conservation" in the quantum setting.

A. Proper mixture

Consider now the case of a system in thermal equilibrium at temperature T. We take the system in question once again to be a simple harmonic oscillator. This corresponds, as is traditionally treated on statistical mechanic textbooks, to an ensemble (a canonical ensemble) of identical particles and might be described by the proper mixture:

$$\rho = N\Sigma_{n=0}^{\infty} e^{-\beta n\epsilon}|n\rangle\langle n|, \tag{26}$$

where $N = 1 - e^{-\beta\epsilon}$ is a normalization constant ensuring $\mathrm{Tr}\rho = 1$. The mean energy is then $\langle E\rangle_T = \frac{\epsilon e^{-\beta\epsilon}}{1-e^{-\beta\epsilon}}$. Let us now consider a two level detector (as in the previous discussion) initially in the un-excited state (and vanishing energy) and make it interact with our thermal ensemble (for this purpose we in fact consider an ensemble of identical detectors). The initial situation will thus be described by the density matrix:

$$\rho\otimes|\downarrow\rangle\langle\downarrow| = [(1-e^{-\beta\epsilon})\Sigma_{n=0}^{\infty} e^{-\beta n\epsilon}|n\rangle\langle n|]\otimes|\downarrow\rangle\langle\downarrow|. \tag{27}$$

After letting the system interact for a suitably long time we will find a result analogous to that encountered in the Eq. (3.25) of [15]. That is:

$$\rho_{\mathrm{late}} = N_{\mathrm{late}}\Sigma_{n=0}^{\infty} e^{-\beta n\epsilon}[|n\rangle\otimes|\downarrow\rangle - i\gamma\sqrt{n}|n-1\rangle\otimes|\uparrow\rangle + \ldots] \times [\langle n|\otimes\langle\downarrow| + i\gamma\sqrt{n}\langle n-1|\otimes\langle\uparrow| + \ldots] \tag{28}$$

where γ is a small parameter representing the strength and time duration of the interaction and N_{late} is a normalization constant ensuring $\mathrm{Tr}\rho_{\mathrm{late}} = 1$.

B. Improper mixture

Consider now the situation in which we are informed that the detector is excited. To do this we simply compute a partial trace after applying the corresponding projector $|\uparrow\rangle\langle\uparrow|$. The resulting density matrix up, to first order in the expansion is

$$\rho_{\mathrm{post}} = N_{\mathrm{post}}\Sigma_{n=1}^{\infty} e^{-\beta n\epsilon} n|n-1\rangle\langle n-1|, \tag{29}$$

where $N_{\mathrm{post}} = (\frac{e^{-\beta\epsilon}}{(1-e^{-\beta\epsilon})^2})^{-1}$ is still another normalization constant ensuring $\mathrm{Tr}\rho_{\mathrm{post}} = 1$.

We note that the absence of the first term ought be considered to be the result of both the measurement of the excitation level of the exited detector and the post selection within the ensemble in which the elements containing the unexcited detector state were removed. In other words, we have that in going from Eq. (28) to Eq. (29) actually modifies the set of systems composing the ensemble itself.

The next observation is that the resulting ensemble, represented by the density matrix (29) is not longer thermal, which in turn implies that the expectation value of the energy

$$\mathrm{Tr}(\hat{H}\rho_{\mathrm{post}}) = N_{\mathrm{post}}\Sigma_{m=0}^{\infty} e^{-\beta(m+1)\epsilon}(m+1)m\epsilon = \frac{2\epsilon e^{-\beta\epsilon}}{(1-e^{-\beta\epsilon})}. \tag{30}$$

differs slightly from that of a truly thermal ensemble ($\langle E\rangle_T = \frac{\epsilon e^{-\beta\epsilon}}{1-e^{-\beta\epsilon}}$). We note again that even after adding the energy of excitation of the detector ϵ the expectation value of the energy has changed as the result of the measurement (and postselection).

Finally let us consider the case of a single pair of entangled harmonic oscillators with "thermal weights." That is the case of a pure state, which upon consideration of the reduced density matrices for each oscillator would result in a thermal density matrix, but of course one of an improper nature. Let us add a detector arranged to interact just with the harmonic oscillator (II) and which is initially prepared in the unexcited state, so the state of the complete system is

$$|\Psi\rangle_{\mathrm{initial}}^{(T)} = A_{\mathrm{initial}}\Sigma_{n=0}^{\infty} e^{-\beta n\epsilon}|n\rangle^{(I)}\otimes|n\rangle^{(II)}\otimes|\downarrow\rangle, \tag{31}$$

where $A_{\mathrm{initial}} = \sqrt{1-e^{-2\beta\epsilon}}$. After letting the system interact for a suitably long time, the state of the system will be

$$|\Psi\rangle_{\mathrm{late}}^{(T)} = A_{\mathrm{late}}\Sigma_{n=0}^{\infty} e^{-\beta n\epsilon}|n\rangle^{(I)} \otimes [|n\rangle^{(II)}\otimes|\downarrow\rangle - i\gamma\sqrt{n}|n-1\rangle^{(II)}\otimes|\uparrow\rangle + \ldots]. \tag{32}$$

If we project on the subspace corresponding to the excited detector (ignoring the irrelevant factor i and ensuring that the final state is normalized) we find:

$$|\Psi\rangle^{(T)}_{\mathrm{pos-tsel}} = A_{\mathrm{post-sel}} \Sigma_{n=0}^{\infty} e^{-\beta n\epsilon} |n\rangle^{(I)} \otimes [\sqrt{n}|n-1\rangle^{(II)} + \ldots], \tag{33}$$

where in this case $A_{\mathrm{post-sel}} \approx (\frac{e^{-2\beta\epsilon}}{(1-e^{-2\beta\epsilon})^2})^{-1/2}$.

Now the expectation values of the energy of the harmonic oscillator I is $\langle E\rangle_I = \frac{\epsilon(1+e^{-2\beta\epsilon})}{(1-e^{-2\beta\epsilon})}$, which is different from that of a purely thermal state, while the expectation value of the harmonic oscillator II is $\langle E\rangle_{II} = \frac{2\epsilon e^{-2\beta\epsilon}}{(1-e^{-2\beta\epsilon})}$. Note that if we add the energy of the excited detector, ϵ, we have that $\langle E\rangle_{II} + \epsilon = \langle E\rangle_I$. On the other hand the expectation energies of both cases have definitely changed as a result of the "projection."

At this point we might find nothing seriously puzzling because we have been dealing with systems that are not initially in energy eigenstates. The change in the case of the harmonic oscillator II in the above example, is of course a bit puzzling due to the fact that this oscillator has not been made to directly interact with a detector to bring about the "projection," however this is nothing more than the usual change occurring as a result of quantum entanglement with a second system which has been subjected to a measurement.

It is worth noting that, in dealing with mixed states it is only when the expectation of the energy momentum is extracted from an improper mixture (and thus, indirectly from a pure state) that it make sense to use it in semiclassical gravity. If this is done with a proper mixture what we would obtain is indeed an average value of that quantity over some ensemble (such as in the case of the many realizations involved in stochastic gravity) and then self consistency will be in doubt.[1]

The situation considered in Sec. IV is, however a bit more troublesome as it seems to have the potential for a serious violation of our fundamental ideas about relativity, in particular, the potential to offer a path for superluminal communication.

V. FASTER THAN LIGHT SIGNALING?

For the remainder of this section, and as we have specified in the introduction, we make the assumption that a measurement of the detector leads to a state collapse along a Cauchy surface. We emphasise that the superluminal puzzle that will be discussed does not rely on a particular choice of a Cauchy surface, but rather occurs for any such choices. We next offer our motivations for taking this point of view.

Let us stress that when adopting a purely "operational interpretation" of quantum states, in which they merely encode "information" about a system, as is done in [23,28], the framework would lie outside the scope of our assumptions. On the other hand, there are good reasons to assume that quantum states contain more information than mere subjective knowledge about the system, and that measurements are more than mere updates of that state of knowledge, most notably the aforementioned PBR theorem [29].

Moreover, it is interesting to examine the consequences of a Cauchy collapse measurement modeling, since it is assumed, implicitly or explicitly, in many discussions on quantum field theory. Let us explain: When describing the quantum theory for a physical system, the most straightforward procedure is to consider all degrees of freedom of the system, and work with the quantum state of the complete system as vectors (more precisely rays) on that Hilbert space. When one is interested in just a few degrees of freedom, one might trace over the other degrees of freedom, obtaining a reduced density matrix, corresponding to an improper mixture according to d' Espagnat's terminology discussed above. In the case of quantum fields in globally hyperbolic spacetimes, quantum states are generically thought of as associated to Cauchy surfaces. Indeed, global hyperbolicity, and the foliation by Cauchy surfaces is used explicitly in the construction of QFT in curved spacetimes, see for instance the first paragraph of Sec. 4. 2 in [7]. That states are associated with Cauchy surfaces in many QFT treaties of the more "particle physics" flavor becomes quite evident when working in the interaction picture, which, in curved spacetimes, becomes generalized by the Tomonaga-Schwinger formalism: Given a globally hyperbolic spacetime and Ψ_Σ a state on the Cauchy surface Σ, the state on the surface Σ' obtained from Σ by deforming it to enclose between the two the infinitesimal four volume δV around the event x is given by

$$i\delta|\Psi_\Sigma\rangle = \lim_{\delta V\to 0} \frac{|\Psi'_\Sigma\rangle - |\Psi_\Sigma\rangle}{\delta V} = \mathcal{H}(x)\delta V|\Psi_\Sigma\rangle \tag{34}$$

where $\mathcal{H}(x)$ is the interaction Hamiltonian density.

Often the used wording indicates one is associating states to spacelike surfaces, without specifying that such surfaces are implicitly taken to be Cauchy spacelike surfaces. That is illustrated by the fact that the states one talks about are pure states. As previously noted, in discussions concerning for instance *the Unruh effect* and related topics, when considering states associated with spacelike surfaces that fail to be Cauchy, the state should correspond to an improper mixture, indicating that one is dealing with a partial characterization of a pure state associated with a full Cauchy surface.

[1]In fact, we do not know what it means to make averages over collections of space-times and the nonlinearity of GR clearly casts serious doubts that, say, Einstein's equations would be preserved under any kind of averaging one might want to consider.

The association of states of quantum fields with Cauchy surfaces is quite evident, for instance, in the analysis of the black hole information problem (see for instance: [30] or [31]). In this case, when one restricts consideration to surfaces that are not really Cauchy ones, one finds, as previously noted, states that are improper mixtures.[2]

In view of these considerations, we are motivated to consider sates of quantum fields as associated with Cauchy surfaces, and thus, to take the updating of states as a result of measurements (or spontaneous collapses) to take place on Cauchy surfaces. As we will see, however, this rather natural posture seems to be at the core of the problem we shall discuss in the present section. We note that the association of states of quantum fields with Cauchy surfaces is also explicitly referred to in previous works involving discussions on superluminal communication and measurements, e.g., [34–38]. In this sense, this work adds to this existing literature.

At the root of the possibility of even asking oneself whether superluminal signaling is possible, is the, by now rather well understood fact, that our world contains some nonlocal features. This is reflected in the violations of Bell's inequalities [39], which have been experimentally confirmed in multiple experiments [40–44].[3]

Nevertheless, there are a number of arguments, and a general widespread conviction that such non-locality cannot be exploited to communicate superluminally. Indeed, the nonlocality present in the situation examined by Bell does not allow for superluminal communication. Faster-than-light communication would force us to revise the very foundations of special relativity—and of physics as a whole. The widespread posture in the physics community is that somehow nature contains features that prevent the nonlocal nature of the quantum state (which, as noted, cannot be taken as being purely epistemic, as per the PBR theorem [29]) from being used for superluminal signaling. In any case, "textbook" quantum mechanics do not seem to offer paths that allow for faster-than-light communication.[4] This is captured by the set of results known as the no-signaling theorem.

[2]As discussed in [32,33] that outcome is quite natural and involves—as argued in both Wald's and Maudlin's works—no paradox, unless one assumes that quantum gravity effects not only cure the singularity, but remove the need to add anything, like a boundary, to the spacetime manifold, and that, somehow, at least some sense of global hyperbolicity is restored.

[3]Some even simpler theoretical settings, such as the GHZ construction [45] (which have not been experimentally realized due to technical difficulties), are expected to provide further, and even more transparent evidence, for nonlocality. See for instance the discussion about the GHZ scheme in [46].

[4]This feature is shared by most alternatives being considered in connection with the quest to resolving the conceptual difficulties afflicting quantum theory related to "the measurement problem." In the present work, we will not enter—except for a few comments—into the discussion involving those issues, a matter that will be left for future studies.

Things become more complicated in the context of quantum field theory, as the question of what quantities (represented by, say, self-adjoint operators), can be measured or not is one where there is no complete and generally accepted answer, see [47] and a possible resolution [48]. The situation we study here involves yet another aspect that seems problematic.

Under the Cauchy collapse assumption, the change that we have noted in the expectation value of the stress-energy tensor in the left Rindler wedge (region II) due to measurements taking place in the right Rindler wedge (region I)—if detectable—however seems to provide a path for superluminal communication, which we describe in the following *gedankenexperiment*:

Gedankenexperiment: Suppose a linearly uniformly accelerated observer in the right Rindler wedge, say Alice, is equipped with a highly-efficient detector, for which it is highly probable to detect a Rindler particle. This enables Alice to affect the expectation value of the stress-energy tensor in the left Rindler wedge by choosing to turn her detector on or not. A causally-disconnected observer, say Bob, in the left Rindler wedge, who can probe, either the state of the field, or more specifically, the expectation value of stress-energy tensor (for example by probing the gravitational field with a torsion balance), would then be able to infer whether Alice has or has not turned on her detector in the right Rindler wedge. This seems to be a path for achieving superluminal communication between Alice and Bob as all events in the world line of one are spacelike separated from all events of that of the other.

A simple protocol for superluminal signaling relying on the above *Gedankenexperiment* can be thought of as follows.

Alice and Bob are given instructions that Alice could send to Bob a signal (for instance that she has decided on something and the answer is "yes") by turning her detector on. She would *not* turn her detector at all if the answer is "no". Bob will then monitor the value of the expansion of nearby geodesics by, say, looking at the freely falling particles he is continuing releasing. Bob must be careful to ensure that nothing he does generates in his surroundings any energy momentum tensor that could mimic that associated with the change in the state resulting from a detection of a quantum field by Alice's detector. If, at any time in his world line, he detects the corresponding geodesic expansion he would know Alice's decision is "yes". The point is that no matter when, along his world line, that would happen, the communication would be superluminal, as Alice and Bob are never in causal contact. This is, of course, a rather poor quality signaling protocol because Bob could eventually know if Alice's decision is "yes", but he would never know if Alice's decision is "no". This can be remedied by having Alice, use two different kind of detectors (with, say, vastly different energies of excitation), which could produce two different changes in the energy momentum tensor at Bob's location. The detection of one of the two changes would indicate if Alice's answer is "yes" or "no", thus corresponding to a complete 1- bit signaling.

The efficacy of this protocol depends on the magnitude of Alice's and Bob's absolute accelerations a_{A} and a_{B} and the change on the state of the quantum field induced by the excitation of Alice's detector, which we can modify by selecting, say, the detector's internal energy gap, Ω. Ignoring switching effects, reasonably good performance for the detector could be achieved by setting the energy gap near the Unruh temperature (in natural units), since for long-time interactions the transition probability of the detector should behave as a Planckian distribution [5], while at the same time, Ω could be made arbitrarily large. On the other hand, the distance D_{A} along a surface orthogonal to ξ from Alice's world line and the bifurcating surface of the Killing Horizon is determined by a_{A}, and decreases as the later increases. So for larger D_{A} the protocol becomes poorer. Optimization of the protocol performance with respect to a_{B} will have to balance the fact that the analogous distance from Bob to the bifurcating horizon decreases with increasing a_{B} and the detailed functioning of the device he uses in detecting the changes in the expectation value of the energy momentum tensor of the field.

It is worth pointing out that, as the signaling protocol relies on Bob's measurements of the expectation value of stress-energy tensor in region II, one must consider what such measurement implies. On one hand, and as is generally the case when considering expectation values, their measurement involves making several projective measurements of the quantity of interest, and then taking suitable averages of the corresponding results. In the case at hand, one could imagine Bob carrying out the multiple measurements over a period of time, but, of course, as the situation is not stationary (the effects of interest would not be the same at all points along Bobs world-line), one could be concerned that here one faces a serious impediment. This concern seems to us to be misguided, as there is no reason[5] that Bob could not "renormalize" the results of each experiment in an appropriately differential manner, according to the point on his world line at which each piece of data is taken,[6] and only then, proceed to make the corresponding average. Moreover, Bob could, instead, carry out the measurements all at once by using a large number of the appropriate devices. Here, one might worry about the effect one device could have on the others, or the effect that each device could have on the state of the quantum field itself. We think these concerns could be addressed, for instance, by choosing to rely on the so called weak measurement schemes in which the effect of each probe on the state of the system under measurement can be made arbitrarily small [49]. So, it is not clear to us that such considerations would be enough to prevent the measurement in question with enough accuracy, for Bob to distinguish whether the changes induced by Alice actions have taken place or not. Finally, and in what was the original consideration behind this manuscript, one could consider the setting in a semi-classical gravity context in which it would be the expectation value of the energy momentum tensor what dictates how the spacetime is curved, a context in which, as noted earlier, Bob could directly observe, say, the impact by Alices decision, on the Ricci tensor, a quantity that, at the classical level, can be measured directly via the geodesic deviation among nearby geodesics. In the end, while all these questions are clearly relevant, the central point is that a serious problem would arise even if the tiniest amount of information is transmitted superluminally. Namely, even if the only thing the protocol manages to achieve is to modify Bobs probability assessment (which in the absence of any new information was $1/2$) regarding Alices choice to turn her detector on or not, we would face a serious conceptual conundrum.

There seems to be no trivial or evident obstacle for the protocol to work at some level of reliability, and, as noted, any nonzero value of that would represent the opening of a door for superluminal communication. Indeed, there exist currently serious proposals to measure the gravitational field with very high precision, e.g., [50] what would be opening the path for an experiment relying on the semi-classical path. We will discuss further down the manuscript what seem to be the most natural possibilities by which nature might prevent, in principle, the working of such protocol.

A. Possible obstacles faced by the superluminal signaling protocol

The possibility of Alice signaling to Bob in a superluminal way should be considered quite problematic in view of the implications that would have for our understanding of the world. In fact, the situation is aggravated by the fact that, even if the detector does not detect a particle, there are higher-order effects (in the coupling constant) that induce changes in the stress-energy tensor on the left Rindler wedge [51].[7]

Thus, it seems imperative to consider the caveats that might provide a path to avoid such a problematic communication protocol. Before we start, let us make some

[5]Assuming, for instance, that he is aware of Alices world line, and the only thing he does not know beforehand is whether or not she would decide to turn her detector on.

[6]Bob is supposed to know what Alice's world line would be, and only lacks knowledge on whether she would decide to turn on her detector or not.

[7]Although it is not clear if the effect when no particle is absorbed is an artifact of perturbation theory.

observations regarding semiclassical gravity. Semiclassical gravity is a formalism in which the gravitational field is taken as dynamical but treated in classical terms as the metric of a spacetime (M, g_{ab}) in general relativity, while the matter is treated in the language of quantum fields on such spacetime. The metric is taken to satisfy the semiclassical version of Einstein's field equations

$$G_{ab}(x) = 8\pi G\langle T^{\mathrm{ren}}_{ab}(x)\rangle, \tag{35}$$

with $\langle T^{\mathrm{ren}}_{ab}\rangle$, the expectation value of the matter's quantum stress-energy tensor in a suitable quantum state, while the quantum matter obeys the dynamics of QFT. See, e.g., [52] for a review.

In incorporating state collapses in the semiclassical gravity context, we remark that there is an issue, which we have touched on briefly above. The point is that once something like state reduction is considered (be it in the Copenhagen approach or in any of the spontaneous collapse theories models, e.g., [53–64]), one has the ambiguity of which state should be employed in computing the right-hand side of Eq. (35). In other words, it seems natural that, if we are interested in the value of the left-hand side of Eq. (35) at a point x, the right-hand side should be computed using a state associated with a Cauchy surface that passes through x, but, of course, there are infinite such surfaces, and the proposal to consider Eq. (35), even as an approximation, must be completed with a detailed prescription in this regard, in order to, at least, have a well-defined proposal [65–69]. We should emphasize, however, that, no matter what Cauchy surface is used, a portion of the Cauchy surface must extend to the left Rindler wedge. Thus, assuming that the collapse occurs along any arbitrary Cauchy surface, the left Rindler wedge must include a "precollapse" and a "postcollapse" region with different spacetime geometry—according to semiclassical gravity.

Let us now offer, and briefly discuss, what we think is the list of serious set of options to be considered, that could help in avoiding the conclusion of superluminal signaling:

(i) *Strong enough departures from semiclassical gravity, at least in region II (this would concern only the semiclassical detection path discussed above)*:
(ii) *Existence of effects that are indistinguishable from observations of the change in the stress-energy in region II*:
(iii) *A fundamental undetectability of the change of the state in region II, and, in particular, that of expectation value of the stress-energy tensor in said region*:
(iv) *Some fundamental impediment of the construction of the set up*.

Let us now briefly discuss the options (i)–(iv) considered above.

(i) First, and in order to address the issue, we must clarify what is meant by the words "strong enough." We take that to indicate that as a result of some unknown aspect of physics (originating, say, in aspects of quantum gravity), the validity of Eq. (35) with the right-hand side in the updated state, cf. the discussion of Sec. III, would be violated to such a large degree that, for any design of Bob's measuring instrument, the expected result will be modified by a factor of at least the same order of magnitude. As we have pointed out in Sec. III, it is possible to make the expectation value of the stress-energy tensor at points in region II[8] arbitrarily large in the post-collapse state by increasing the detector gap. However, it must be acknowledged that, if Ω becomes large enough, our modeling of the state collapse in semiclassical gravity along a Cauchy surface could be questioned on the basis that a large violation of conservation of the stress-energy tensor along a Cauchy surface is too strong a deviation from the semiclassical regime. Thus, one might argue that, although one could trust semiclassical gravity before and after the state collapse, there is no way to model the system "during" the state collapse in semiclassical terms, and a more refined understanding of how to model state collapses in semiclassical gravity could be required. We note, however, that, even if that is the case, such conclusion will not prevent the superluminal signaling, as all we need is for Bob to detect the effect at any time.

On the other hand, it seems rather unlikely that arbitrarily large departures of semiclassical gravity will be associated with such a simple, and rather common, situation. The full analysis of the question thus requires consideration of the means by which Bob might detect the corresponding change in the spacetime metric in region II.[9] As it seems clear that a modification of the spacetime curvature might, in principle, be measured by the study of the geodesic deviation equation, and, in particular, the expansion of the congruence of geodesics in that region, it appears that the question will have to be connected at least to some fundamental limitation on the validity of the notion of test point particles following geodesics of

[8] It is worth noting that our faster-than-light signaling protocol does not really require the validity of semiclassical gravity during the measurements. The only real requirement is that Bob be able to measure every once in a while the expectation value of the energy momentum tensor by any means. As such, this result seems to contradict the so called no signaling theorems. However, the particular experimental situation we are considering does not fall under the hypothesis of such theorems which concerns finite time measurements. We will discuss this issue further in this paper.

[9] We will not study here in any detail the detection method to be employed by Bob, but just point out that interesting ideas to measure small gravitational effects have been used in say searches for deviation of the universality of free fall, and also note a recent proposal to look for Plank scale dark matter [70].

the underlying metric. There are, of course, some limitations on that arising from simple quantum considerations about the description of the so called free point particles, which, in fact negate, the notion of well defined trajectories.

Moreover, as discussed in [71], the standard quantum mechanical minimal de-localization of a quantum particle (characterized, for instance, by its Compton wavelength) implies such particle ought not to be considered as pointlike, and as is well known, in general, even at the classical level, extended objects fail to follow geodesics. It is unclear at this point if these kind of considerations will be enough to dismiss our example, given the fact that, as noted, the effect could be made as large as one wants, and the time available for Bob to make the measurement is arbitrarily large.

Finally, one can raise the issue of whether one should not trust semiclassical gravity in any situation in which the quantum fluctuations (i.e., the quantum uncertainties) in the energy momentum tensor are larger in magnitude than the expectation value of the energy momentum tensor. Considering that such a restriction would imply that semiclassical gravity cannot be used in the case of Minkowski vacuum, it seems to us that the use of such a generic prohibition would rule out the use of semiclassical gravity in, essentially, all situations. That, it seems to us, would be a very drastic and unwarranted conclusion.

(ii) Here, we must consider the existence of other effects that might not be effectively distinguished from the changes resulting from the collapse. Those might be intrinsically associated with the gravitational effects of either the measuring devices that Bob would be introducing in order to detect the changes in the spacetime metric in region II. They also might be associated with the mere existence of (the non-vanishing stress-energy tensor corresponding to) Alice and her detectors, whose own gravitational effects we have been neglecting so far, although, classically, these effects will, in principle, propagate causally and should not affect region II. It is likely that deviations from Minkowski spacetime reflect on the *nonlocal* state of the field from a semiclassical-gravity viewpoint in such a way that the quantum state of the field cannot be, in principle, the Minkowski vacuum state to begin with. Moreover, the stress-energy of Alice and her apparatus must be considered in the constraints of the theory, and the spacetime will have a non-vanishing ADM mass, reflected in the asymptotic behavior in all spacetime directions.

Another possibility is to consider that the model used for Alice's detector, although standard in the literature, is a local model (in Region I) of the detector. Treating Alice's detector as a quantum field yields a detector model with support in both Regions I and II. In the Heisenberg picture, we know that the existence of a field-like detector for Alice will only affect the dynamics in the causal future of the coupling region between Alice's detector (or probe) and the quantum field considered as a system. This can also be formalized in the algebraic QFT language [23]. However, as emphasized in [23], a local and covariant measurement scheme is only able to describe the *probe-system measurement chain* under the assumption that somewhere, someone knows how to measure something. Thus, it seems that the issue of sending a faster-than-light signal by the "act of measurement" cannot be resolved by simply changing the detector model. Sorkin has pointed out, in his "impossible measurements" protocol [47] that the existence of spatially extended detectors would make faster than light signaling almost unavoidable. See, however [48].

(iii) This possibility could result from a variety of reasons. For instance, it might be that the appropriate way to evaluate expectation value of the stress-energy tensor, at any point, is taken on states associated with 3-surfaces which would not incorporate the change in the state resulting from a collapse or a measurement on region I. One such option would require to take the expectation value of the stress-energy tensor at x to be computed using the state corresponding to something like the surface $\partial J^-(x)$ (the boundary of the causal past of x), and that, at the same time, the detailed theory characterising the measurement (i.e., something like the spontaneous collapse theories considered in [72–74]) is such that the state associated with $\partial J^-(x)$ is unaffected. A scheme of that kind would ensure that the measurement of Alice would have no consequences at spacelike-separated events. Such proposal is not without difficulties, one of which is the fact that, in general, $\partial J^-(x)$ is not a Cauchy surface, and another is that it is not smooth at x. The first difficulty might be resolved if one has some good reason to consider that there is a certain initial state associated with an initial Cauchy surface Σ_{in}, which might be considered as the initially "*prepared*" state, or perhaps *the initial state of the universe* or something like that, and then to consider computing the expectation values of quantities of interest at x, using the state associated with the surface $(I^+(\Sigma_{\text{in}}) \cap \partial J^-(x)) \cup (\Sigma_{\text{in}} - J^-(x))$, while the second problem might be dealt with an adjustment of the recipe based on the taking of appropriate limits of a succession of smooth Cauchy surfaces that have as a limit (in a suitable sense), the surface indicated above. The issue is, however, quite delicate as illustrated by the discussions in [65], the pursuit of which lies outside the scope of the present work.

An argument in favor of the impossibility of detecting the signal is the following. One can assume that stress-energy conservation must hold, *on average*, i.e., for successive measurements, which prevents one from extracting arbitrarily large amounts of energy from the quantum field by detector measurements, as argued in [15]. Thus, an arbitrarily efficient detector must produce arbitrarily small changes for the stress-energy tensor when it detects a Rindler particle. The issue now becomes whether there exists an "engineering window" where the efficiency of the detector and the amplitude of the signal can be compensated, such that a faster-than-light signal is measurable by Bob, and that such a signal can be sent with sufficient certainty by Alice. On this point, it follows, from the inequality $\Delta E \Delta T \geq \hbar$, that, in order to detect such a small signal, Bob requires a large amount of time. On the other hand, in principle, Bob has an infinite time to measure an arbitrarily small signal, for he can orbit along a Killing vector boost orbit in the left Rindler wedge, and one could argue that such engineering window can always be found. However, any physical signal sent by Alice must decay as it approaches $\mathcal{I}^+$. Thus, it is expected that the signal will become weaker as Bob's proper time elapses, and this might render the signal "effectively undetectable ".

(iv) Finally, there could be aspects of the set up that would make it simply unfeasible. One possibility we should consider is the following. In order for Bob to be sure that the change in the energy momentum he observes corresponds indeed to a signal that was sent by Alice, he must be sure that similar signals are not reaching him coming from elsewhere. That is, he must be quite sure that the state of the quantum field prior to Alice's manipulation of her detector is the Minkowski vacuum, and, as he must be ready to receive Alice's signal at any point in his world line, he must be sure that such characterization of the state of the quantum field must be the appropriate one up to arbitrarily distant regions of "space". So, there must be a preparation of the initial state of the quantum field on an extremely large spatial region occurring well before the whole protocol is even started (and if we want to ensure Bob has, in effect, an arbitrarily long time to make the detection, it seems the preparation of the state ought to encompass a full Cauchy surface. However, as already noted, the simultaneous (in some reference frame) measurements of observables associated with such extended spatial regions—say, as a means of state preparation—are known to be rather problematic in their own right as argued by Sorkin [47] (again, see however [48]). Furthermore, such a global preparation seems to contain non-local elements that are similar to those occurring in Sorkin's Impossible measurements example, and might be subjected to similar questionings. Moreover, one might argue that, as we need to include Alice, Bob and their measuring devices then, the state of the field could not possibly be, strictly speaking, the Minkowski vacuum. It would be surprising if such concerns could not be overcome, even in principle. Needless is to say that a rather general argument along such lines is not available at this time as far as we know.

In any case, such arguments must be considered with care, because, in principle, even a very unreliable communication protocol which allows faster than light communications would be quite problematic. We think these ideas deserve much deeper exploration.

VI. FINAL REMARKS

We have revisited the issue of detection of a Rindler particle by an accelerator detector confined to the right Rindler wedge, focusing attention on the implications of the reduction of the state associated with the actual detection or what is often referred to as the measurement part of the process. We have noted that the concomitant modification of the state of the quantum field, imply changes in the (expectation value of the renormalized) stress-energy tensor in both Rindler wedges, under the assumption that a measurement induces a state collapse along a Cauchy surface.

Concerning the right Rindler wedge, it is interesting that the resulting state and stress-energy tensor expectation values become non-thermal which is, in a sense, easy to understand as result of the disruption brought by the interaction with the detector. Of course, one expects that in the long run further interaction between the field and detector or among the field modes themselves will bring the system to a new state of thermal equilibrium.

Concerning the left Rindler wedge, however, the change in the stress-energy tensor expectation value is more problematic, as it seems to open a possibility for faster than light communication, under the assumptions here contemplated. We have noted some of what we see as the most natural possibilities to avoid such conclusion, but further studies of these issues are required in order to get to more definite and solid conclusions.

For the moment, we can only stress that the fuzziness in the theoretical characterization of the act of measuring in quantum theory can bring about complications even in apparently innocuous circumstances, such as in the context of the detection of a Rindler particle in Minkowski spacetime. Despite being overlooked in many practical applications in physics, the literature on the measurement problem is quite large, and the positions taken in its face are quite varied. We do not intend to discuss them in detail. It suffices for us to mention that the paths to overcome the measurement problem have been classified by Maudlin [75] as follows: It is internally inconsistent to hold simultaneously the following three propositions about a quantum theory:

(i) The description of a physical system as provided by the quantum state or wave function is complete.

(ii) The evolution of the quantum state is always dictated by the Schrödinger equation (or its relativistic generalizations).

(iii) Individual experiments produce definite (even if often unpredictable) results.

Thus, one must negate at least one of (i)–(iii) above. Negating (i) leads down, in general, the path of the so-called *hidden variable theories*, of which the de Broglie-Bohm theory is the best known example [76]. Negating (ii) implies that the collapse or reduction of the wave function plays a central role, as in the Copenhaguen textbook interpretation, as well as in so-called *spontaneous collapse* or *dynamical state reduction* theories, such as GRW, CSL, Penrose-Diósi, etc [72–74,77–80]. The negation of (iii) leads to *many-worlds*- or *many-minds*-type interpretations [81,82].

In any case, we should restate that the superluminar tension relies on the application of two central hypotheses in this work. The first one is that measurements induce the collapse of the wave function, and that such collapse occurs on a Cauchy surface of spacetime. The second one is that the problem studied in (nearly) flat spacetime is within the regime of applicability of semiclassical gravity. It is well possible that either one of the two hypotheses fail, or that they cannot be taken to hold together. If this is the case, it would seem that there exist apparently innocuous situations, such as the one here studied, for which a correct theoretical description of the evolution of the system must resort to quantum gravity. This seems to be in agreement with conclusions drawn from [83,84]. However if one takes semiclassical gravity to fail in this case, it seems likely that other situations that are often discussed using semiclassical language, for instance Hawking radiation and the back-reaction in black holes (even at early times), might require reconsideration as well. On the other hand, if one takes the Cauchy collapse postulate to be inappropriate, it is unclear how to assign an interpretation to the quantum state of a system that is in harmony with the PBR theorem (i.e., that is not purely epistemic). Another possibility is that semiclassical gravity and measurements can coexist (in the semiclassical regime of quantum gravity) in a fashion that is yet to be better understood, and beyond the "effective" treatment considered in [85].

Finally, it is our hope to draw the attention of the community to these delicate issues that afflict our understanding of quantum theory in general and of QFT in particular, especially in gravitational contexts.

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APPENDIX A: THE MINKOWSKI VACUUM IN THE RINDLER WEDGES

The initial state is taken to be the Minkowski vacuum, whose Wightman function takes the following form in the right Rindler wedge

$$\langle \Omega_{\mathrm{M}} | \hat{\Phi}(x) \hat{\Phi}(x') \Omega_{\mathrm{M}} \rangle = \int_{\mathbb{R}^+ \times \mathbb{R}^2} \mathrm{d}^3\kappa \int_{\mathbb{R}^+ \times \mathbb{R}^2} \mathrm{d}^3\kappa' \langle \Omega_{\mathrm{M}} | \Big(v_{I\vec{\kappa}}(x) \hat{a}_{\mathrm{R}\vec{\kappa}} + \overline{v_{I\vec{\kappa}}(x)} \hat{a}^*_{\mathrm{R}\vec{\kappa}} \Big) \Big(v_{I\vec{\kappa}'}(x') \hat{a}_{\mathrm{R}\vec{\kappa}'} + \overline{v_{I\vec{\kappa}'}(x')} \hat{a}^*_{\mathrm{R}\vec{\kappa}'} \Big) \Omega_{\mathrm{M}} \rangle. \quad \text{(A1)}$$

Using the relations [[9], Eqs. (2.125–(2.127)],

$$\langle \Omega_{\mathrm{M}} | \hat{a}^*_{\mathrm{R}\vec{\kappa}} \hat{a}_{\mathrm{R}\vec{\kappa}'} \Omega_{\mathrm{M}} \rangle = \langle \Omega_{\mathrm{M}} | \hat{a}^*_{\mathrm{L}\vec{\kappa}} \hat{a}_{\mathrm{L}\vec{\kappa}'} \Omega_{\mathrm{M}} \rangle = (e^{2\pi\omega/a} - 1)^{-1} \delta^3(\vec{\kappa} - \vec{\kappa}'), \quad \text{(A2a)}$$

$$\langle \Omega_{\mathrm{M}} | \hat{a}_{\mathrm{R}\vec{\kappa}} \hat{a}^*_{\mathrm{R}\vec{\kappa}'} \Omega_{\mathrm{M}} \rangle = \langle \Omega_{\mathrm{M}} | \hat{a}_{\mathrm{L}\vec{\kappa}} \hat{a}^*_{\mathrm{L}\vec{\kappa}'} \Omega_{\mathrm{M}} \rangle = (1 - e^{-2\pi\omega/a})^{-1} \delta^3(\vec{\kappa} - \vec{\kappa}'), \quad \text{(A2b)}$$

$$\langle \Omega_{\mathrm{M}} | \hat{a}_{\mathrm{L}\vec{\kappa}} \hat{a}_{\mathrm{R}\vec{\kappa}'} \Omega_{\mathrm{M}} \rangle = \langle \Omega_{\mathrm{M}} | \hat{a}^*_{\mathrm{L}\vec{\kappa}} \hat{a}^*_{\mathrm{R}\vec{\kappa}'} \Omega_{\mathrm{M}} \rangle = (e^{\pi\omega/a} - e^{-\pi\omega/a})^{-1} \delta^3(\vec{\kappa} - \vec{\kappa}'), \quad \text{(A2c)}$$

(and zero otherwise) we obtain that

$$\langle\Omega_{\rm M}|\hat{\Phi}(x)\hat{\Phi}(x')\Omega_{\rm M}\rangle=\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\left(\frac{v_{I\vec{\kappa}}(x)\overline{v_{I\vec{\kappa}}(x')}}{1-e^{-2\pi\omega/a}}+\frac{\overline{v_{I\vec{\kappa}}(x)}v_{I\vec{\kappa}}(x')}{e^{2\pi\omega/a}-1}\right). \tag{A3}$$

Likewise, in the left Rindler wedge the Wightman function takes the form

$$\langle\Omega_{\rm M}|\hat{\Phi}(x)\hat{\Phi}(x')\Omega_{\rm M}\rangle=\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\left(\frac{v_{II\vec{\kappa}}(x)\overline{v_{II\vec{\kappa}}(x')}}{1-e^{-2\pi\omega/a}}+\frac{\overline{v_{II\vec{\kappa}}(x)}v_{II\vec{\kappa}}(x')}{e^{2\pi\omega/a}-1}\right), \tag{A4}$$

where $v_{II\vec{\kappa}}$ are left Rindler modes.

APPENDIX B: NORMALIZATION OF THE UPDATED STATE

We have seen in Sec. II that if the detector clicks the updated state of the field becomes

$$|f\rangle=-i\mathcal{N}\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_I\mathrm{dvol}(x)\zeta(x)\left(v_{I\vec{\kappa}}(x)\hat{a}_{\mathrm{R}\vec{\kappa}}+\overline{v_{I\vec{\kappa}}(x)}\hat{a}^*_{\mathrm{R}\vec{\kappa}}\right)|\Omega_{\rm M}\rangle. \tag{B1}$$

In this appendix, we see that the normalization constant is given by Eq. (14). We compute

$$\begin{aligned}\mathcal{N}^{-2}\langle f|f\rangle=&\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa'\int_I\mathrm{dvol}(x)\int_I\mathrm{dvol}(x')\overline{\zeta(x)}\zeta(x')\\&\times\langle\Omega_{\rm M}|\left(v_{I\vec{\kappa}}(x)\hat{a}_{\mathrm{R}\vec{\kappa}}+\overline{v_{I\vec{\kappa}}(x)}\hat{a}^*_{\mathrm{R}\vec{\kappa}}\right)\left(v_{I\vec{\kappa}'}(x')\hat{a}_{\mathrm{R}\vec{\kappa}'}+\overline{v_{I\vec{\kappa}'}(x')}\hat{a}^*_{\mathrm{R}\vec{\kappa}'}\right)\Omega_{\rm M}\rangle.\end{aligned} \tag{B2}$$

Using the relations [[9], Eqs. (2.125)–(2.126)], we obtain that

$$\mathcal{N}^{-2}\langle f|f\rangle=\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_I\mathrm{dvol}(x)\int_I\mathrm{dvol}(x')\overline{\zeta(x)}\zeta(x')\left(\frac{v_{I\vec{\kappa}}(x)\overline{v_{I\vec{\kappa}}(x')}}{1-e^{-2\pi\omega/a}}+\frac{\overline{v_{I\vec{\kappa}}(x)}v_{I\vec{\kappa}}(x')}{e^{2\pi\omega/a}-1}\right). \tag{B3}$$

APPENDIX C: THE TWO-POINT FUNCTION IN THE UPDATED STATE IN THE RIGHT RINDLER WEDGE

The updated stress-energy tensor can be computed from the two-point function in the updated state. In this appendix, we show that the two-point function in the updated state in the right Rindler wedge is given by

$$\langle f|\hat{\Phi}(x)\hat{\Phi}(x')f\rangle=\langle\Omega_{\rm M}|\hat{\Phi}(x)\hat{\Phi}(x')\Omega_{\rm M}\rangle+\Delta_{\rm R}(x,x'). \tag{C1}$$

where $\Delta_{\rm R}(x,x')$ is given by Eq. (C12) below.

The two-point function in the right Rindler wedge reads

$$\begin{aligned}\langle f|\hat{\Phi}(x)\hat{\Phi}(x')f\rangle=\mathcal{N}^2&\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa'\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p'\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\\&\times\langle\Omega_{\rm M}|\left(v_{I\vec{p}}(y)\hat{a}_{\mathrm{R}\vec{p}}+\overline{v_{I\vec{p}}(y)}\hat{a}^*_{\mathrm{R}\vec{p}}\right)(v_{I\vec{\kappa}}(x)\hat{a}_{\mathrm{R}\vec{\kappa}}+\overline{v_{I\vec{\kappa}}}(x)\hat{a}^*_{\mathrm{R}\vec{\kappa}})\\&\times(v_{I\vec{\kappa}'}(x')\hat{a}_{\mathrm{R}\vec{\kappa}'}+\overline{v_{I\vec{\kappa}'}}(x')\hat{a}^*_{\mathrm{R}\vec{\kappa}'})\left(v_{I\vec{p}'}(y')\hat{a}_{\mathrm{R}\vec{p}'}+\overline{v_{I\vec{p}'}(y')}\hat{a}^*_{\mathrm{R}\vec{p}'}\right)\Omega_{\rm M}\rangle.\end{aligned} \tag{C2}$$

Using [[9], Eqs. (2.122)–(2.124)] one can obtain the relations

$$\begin{aligned}&\langle\Omega_{\rm M}|\hat{a}_{\mathrm{R}\vec{\kappa_1}}\hat{a}_{\mathrm{R}\vec{\kappa_2}}\hat{a}_{\mathrm{R}\vec{\kappa_3}}\hat{a}_{\mathrm{R}\vec{\kappa_4}}\Omega_{\rm M}\rangle=0,\\&\langle\Omega_{\rm M}|\hat{a}_{\mathrm{R}\vec{\kappa_1}}\hat{a}_{\mathrm{R}\vec{\kappa_2}}\hat{a}_{\mathrm{R}\vec{\kappa_3}}\hat{a}^*_{\mathrm{R}\vec{\kappa_4}}\Omega_{\rm M}\rangle=0,\end{aligned} \tag{C3a}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3b}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3c}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{R}\vec{\kappa}_2}\hat{a}_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3d}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = \frac{1}{1-e^{-2\pi\omega_1/a}}\frac{1}{1-e^{-2\pi\omega_2/a}}(\delta(\vec{\kappa}_1-\vec{\kappa}_3)\delta(\vec{\kappa}_2-\vec{\kappa}_4)+\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3)), \tag{C3e}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = \frac{1}{1-e^{-2\pi\omega_1/a}}\frac{1}{1-e^{-2\pi\omega_3/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_2)\delta(\vec{\kappa}_3-\vec{\kappa}_4) \tag{C3f}$$

$$+\frac{1}{1-e^{-2\pi\omega_1/a}}\frac{e^{-2\pi\omega_3/a}}{1-e^{-2\pi\omega_3/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3), \tag{C3g}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{R}\vec{\kappa}_2}\hat{a}_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle \tag{C3h}$$

$$=\frac{e^{-2\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{1}{1-e^{-2\pi\omega_4/a}}(\delta(\vec{\kappa}_1-\vec{\kappa}_2)\delta(\vec{\kappa}_3-\vec{\kappa}_4)+\delta(\vec{\kappa}_1-\vec{\kappa}_3)\delta(\vec{\kappa}_2-\vec{\kappa}_4)), \tag{C3i}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle \tag{C3j}$$

$$=\frac{1}{1-e^{-2\pi\omega_1/a}}\frac{\mathrm{e}^{-2\pi\omega_4/a}}{1-e^{-2\pi\omega_4/a}}(\delta(\vec{\kappa}_1-\vec{\kappa}_2)\delta(\vec{\kappa}_3-\vec{\kappa}_4)+\delta(\vec{\kappa}_1-\vec{\kappa}_3)\delta(\vec{\kappa}_2-\vec{\kappa}_4)), \tag{C3k}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = \frac{e^{-2\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{\mathrm{e}^{-2\pi\omega_3/a}}{1-e^{-2\pi\omega_3/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_2)\delta(\vec{\kappa}_3-\vec{\kappa}_4) \tag{C3l}$$

$$+\frac{\mathrm{e}^{-2\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{1}{1-e^{-2\pi\omega_2/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3), \tag{C3m}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = \frac{e^{-2\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{\mathrm{e}^{-2\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}}(\delta(\vec{\kappa}_1-\vec{\kappa}_3)\delta(\vec{\kappa}_2-\vec{\kappa}_4)+\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3)), \tag{C3n}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3o}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3p}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3q}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{C3r}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{R}\vec{\kappa}_2}\hat{a}^*_{\mathrm{R}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0. \tag{C3s}$$

Inserting Eq. (C3) into (C2) we obtain that the two point function can be obtained as a sum of six contributions, each coming from one of the non-trivial expressions in Eq. (C3), i.e., it takes the form

$$\langle f|\hat{\Phi}(x)\hat{\Phi}(x')f\rangle(x,x') = \sum_{n=1}^{6} G_{\mathrm{R}n}(x,x'). \tag{C4}$$

Subtracting Eq. (A3) from this expression we obtain that

$$\Delta_{\mathrm{R}}(x,x') = \langle f|\hat{\Phi}(x)\hat{\Phi}(x')f\rangle - \langle\Omega_{\mathrm{M}}|\hat{\Phi}(x)\hat{\Phi}(x')\Omega_{\mathrm{M}}\rangle = \sum_{n=1}^{6}\Delta_{\mathrm{R}n}(x,x') \tag{C5}$$

with

$$\Delta_{\mathrm{R1}}(x,x') \coloneqq \mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{v_{I\vec{p}}(y)v_{I\vec{\kappa}}(x)\overline{v_{I\vec{p}}}(x')\overline{v_{I\vec{\kappa}}(y')}}{(1-e^{-2\pi\omega_p/a})(1-e^{-2\pi\omega_\kappa/a})}, \tag{C6}$$

$$\Delta_{\mathrm{R2}}(x,x') \coloneqq \mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{v_{I\vec{p}}(y)\overline{v_{I\vec{p}}}(x)v_{I\vec{\kappa}}(x')\overline{v_{I\vec{\kappa}}(y')}}{(1-e^{-2\pi\omega_p/a})(1-e^{-2\pi\omega_\kappa/a})}, \tag{C7}$$

$$\begin{aligned}\Delta_{\mathrm{R3}}(x,x') \coloneqq{}& \mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{\overline{v_{I\vec{\kappa}}(y)}v_{I\vec{\kappa}}(x)v_{I\vec{p}}(x')\overline{v_{I\vec{p}}(y')}e^{-2\pi\omega_\kappa}}{(1-e^{-2\pi\omega_\kappa})(1-e^{-2\pi\omega_p})}\\ &+\mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{\overline{v_{I\vec{\kappa}}(y)}v_{I\vec{p}}(x)v_{I\vec{\kappa}}(x')\overline{v_{I\vec{p}}(y')}e^{-2\pi\omega_\kappa}}{(1-e^{-2\pi\omega_\kappa})(1-e^{-2\pi\omega_p})},\end{aligned} \tag{C8}$$

$$\begin{aligned}\Delta_{\mathrm{R4}}(x,x') \coloneqq{}& \mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{v_{I\vec{\kappa}}(y)\overline{v_{I\vec{\kappa}}(x)}\,\overline{v_{I\vec{p}}(x')}\,v_{I\vec{p}}(y')e^{-2\pi\omega_p}}{(1-e^{-2\pi\omega_\kappa})(1-e^{-2\pi\omega_p})}\\ &+\mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{v_{I\vec{\kappa}}(y)\overline{v_{I\vec{p}}(x)}\,\overline{v_{I\vec{\kappa}}(x')}\,v_{I\vec{p}}(y')e^{-2\pi\omega_p}}{(1-e^{-2\pi\omega_\kappa})(1-e^{-2\pi\omega_p})},\end{aligned} \tag{C9}$$

$$\Delta_{\mathrm{R5}}(x,x') \coloneqq \mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{\overline{v_{I\vec{\kappa}}(y)}v_{I\vec{\kappa}}(x)\overline{v_{I\vec{p}}(x')}v_{I\vec{p}}(y')e^{-2\pi\omega_\kappa}e^{-2\pi\omega_p}}{(1-e^{-2\pi\omega_\kappa})(1-e^{-2\pi\omega_p})}, \tag{C10}$$

$$\Delta_{\mathrm{R6}}(x,x') \coloneqq \mathcal{N}^2\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3\kappa\int_{\mathbb{R}^+\times\mathbb{R}^2}\mathrm{d}^3p\int_I\mathrm{dvol}(y)\int_I\mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y')\frac{\overline{v_{I\vec{\kappa}}(y)}\,\overline{v_{I\vec{p}}(x)}\,v_{I\vec{\kappa}}(x')v_{I\vec{p}}(y')e^{-2\pi\omega_\kappa}e^{-2\pi\omega_p}}{(1-e^{-2\pi\omega_\kappa})(1-e^{-2\pi\omega_p})}. \tag{C11}$$

Collecting, we can write

$$\Delta_{\mathrm{R}}(x,x') = \Delta_{\mathrm{R1}}(x,x') + \Delta_{\mathrm{R3}}(x,x') + \Delta_{\mathrm{R5}}(x,x') + \mathrm{c.c.}, \tag{C12}$$

which yields Eq. (18).

APPENDIX D: THE TWO-POINT FUNCTION IN THE UPDATED STATE IN THE LEFT RINDLER WEDGE

The updated stress-energy tensor can be computed from the two-point function in the updated state. In this appendix, we show that the two-point function in the updated state in the right Rindler wedge is given by

$$\langle f|\hat{\Phi}(x)\hat{\Phi}(x')f\rangle = \langle\Omega_{\mathrm{M}}|\hat{\Phi}(x)\hat{\Phi}(x')\Omega_{\mathrm{M}}\rangle + \Delta_{\mathrm{L}}(x,x'), \tag{D1}$$

as in Eq. (19).

The two-point function in the left Rindler wedge reads

$$\begin{aligned}\langle f|\hat{\Phi}(x)\hat{\Phi}(x')f\rangle = \mathcal{N}^2 \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3\kappa \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3\kappa' \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3 p \int_{\mathbb{R}^+\times\mathbb{R}^2} \mathrm{d}^3 p' \int_I \mathrm{dvol}(y) \int_I \mathrm{dvol}(y')\overline{\zeta(y)}\zeta(y') \\ \times \langle\Omega_{\mathrm{M}}|\left(v_{I\vec{p}}(y)\hat{a}_{\mathrm{R}\vec{p}} + \overline{v_{I\vec{p}}(y)}\hat{a}^*_{\mathrm{R}\vec{p}}\right)\left(v_{II\vec{\kappa}}(x)\hat{a}_{\mathrm{L}\vec{\kappa}} + \overline{v_{II\vec{\kappa}}}(x)\hat{a}^*_{\mathrm{L}\vec{\kappa}}\right) \\ \times \left(v_{II\vec{\kappa}'}(x')\hat{a}_{\mathrm{L}\vec{\kappa}'} + \overline{v_{II\vec{\kappa}'}}(x')\hat{a}^*_{\mathrm{L}\vec{\kappa}'}\right)(v_{I\vec{p}'}(y')\hat{a}_{\mathrm{R}\vec{p}'} + \overline{v_{I\vec{p}'}(y')}\hat{a}^*_{\mathrm{R}\vec{p}'})\Omega_{\mathrm{M}}\rangle. \end{aligned} \tag{D2}$$

Using [[9], Eqs. (2.122)–(2.124)] one can obtain the relations

$$\begin{aligned}\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}}\tilde{\delta}(\vec{\kappa}_1-\vec{\kappa}_3)\tilde{\delta}(\vec{\kappa}_2-\vec{\kappa}_4) \\ &\quad + \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-\pi\omega_3/a}}{1-e^{-2\pi\omega_3/a}}\tilde{\delta}(\vec{\kappa}_1-\vec{\kappa}_2)\tilde{\delta}(\vec{\kappa}_3-\vec{\kappa}_4), \\ \langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= 0,\end{aligned} \tag{D3a}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}^*_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{D3b}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{D3c}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{D3d}$$

$$\begin{aligned}\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}^*_{\mathrm{L}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= \frac{1}{1-e^{-2\pi\omega_1/a}}\frac{1}{1-e^{-2\pi\omega_2/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3) \\ &\quad + \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-\pi\omega_3/a}}{1-e^{-2\pi\omega_3/a}}\tilde{\delta}(\vec{\kappa}_1-\vec{\kappa}_2)\tilde{\delta}(\vec{\kappa}_3-\vec{\kappa}_4),\end{aligned} \tag{D3e}$$

$$\begin{aligned}\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}}\tilde{\delta}(\vec{\kappa}_1-\vec{\kappa}_3)\tilde{\delta}(\vec{\kappa}_2-\vec{\kappa}_4) \\ &\quad + \frac{1}{1-e^{-2\pi\omega_1/a}}\frac{e^{-2\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3)\end{aligned} \tag{D3f}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{D3g}$$

$$\langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{L}\vec{\kappa}_2}\hat{a}^*_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle = 0, \tag{D3h}$$

$$\begin{aligned}\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}_{\mathrm{L}\vec{\kappa}_2}\hat{a}^*_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= \frac{e^{-2\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{1}{1-e^{-2\pi\omega_2/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3) \\ &\quad + \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}}\tilde{\delta}(\vec{\kappa}_1-\vec{\kappa}_3)\tilde{\delta}(\vec{\kappa}_2-\vec{\kappa}_4)\end{aligned} \tag{D3i}$$

$$\begin{aligned}\langle\Omega_{\mathrm{M}}|\hat{a}^*_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{L}\vec{\kappa}_2}\hat{a}_{\mathrm{L}\vec{\kappa}_3}\hat{a}_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= \frac{e^{-2\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-2\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}}\delta(\vec{\kappa}_1-\vec{\kappa}_4)\delta(\vec{\kappa}_2-\vec{\kappa}_3) \\ &\quad + \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}}\frac{e^{-\pi\omega_3/a}}{1-e^{-2\pi\omega_2/a}}\tilde{\delta}(\vec{\kappa}_1-\vec{\kappa}_2)\tilde{\delta}(\vec{\kappa}_3-\vec{\kappa}_4), \\ \langle\Omega_{\mathrm{M}}|\hat{a}_{\mathrm{R}\vec{\kappa}_1}\hat{a}^*_{\mathrm{L}\vec{\kappa}_2}\hat{a}^*_{\mathrm{L}\vec{\kappa}_3}\hat{a}^*_{\mathrm{R}\vec{\kappa}_4}\Omega_{\mathrm{M}}\rangle &= 0,\end{aligned} \tag{D3j}$$

$$\langle \Omega_{\mathrm{M}} | \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_1} \hat{a}_{\mathrm{L}\vec{\kappa}_2} \hat{a}^{*}_{\mathrm{L}\vec{\kappa}_3} \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_4} \Omega_{\mathrm{M}} \rangle = 0, \tag{D3k}$$

$$\langle \Omega_{\mathrm{M}} | \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_1} \hat{a}^{*}_{\mathrm{L}\vec{\kappa}_2} \hat{a}_{\mathrm{L}\vec{\kappa}_3} \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_4} \Omega_{\mathrm{M}} \rangle = 0, \tag{D3l}$$

$$\langle \Omega_{\mathrm{M}} | \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_1} \hat{a}^{*}_{\mathrm{L}\vec{\kappa}_2} \hat{a}^{*}_{\mathrm{L}\vec{\kappa}_3} \hat{a}_{\mathrm{R}\vec{\kappa}_4} \Omega_{\mathrm{M}} \rangle = 0, \tag{D3m}$$

$$\begin{aligned}\langle \Omega_{\mathrm{M}} | \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_1} \hat{a}^{*}_{\mathrm{L}\vec{\kappa}_2} \hat{a}^{*}_{\mathrm{L}\vec{\kappa}_3} \hat{a}^{*}_{\mathrm{R}\vec{\kappa}_4} \Omega_{\mathrm{M}} \rangle &= \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}} \frac{e^{-\pi\omega_2/a}}{1-e^{-2\pi\omega_2/a}} \tilde{\delta}(\vec{\kappa}_1 - \kappa_3) \tilde{\delta}(\vec{\kappa}_2 - \vec{\kappa}_4) \\ &\quad + \frac{e^{-\pi\omega_1/a}}{1-e^{-2\pi\omega_1/a}} \frac{e^{-\pi\omega_3/a}}{1-e^{-2\pi\omega_3/a}} \tilde{\delta}(\vec{\kappa}_1 - \kappa_2) \tilde{\delta}(\vec{\kappa}_3 - \vec{\kappa}_4).\end{aligned} \tag{D3n}$$

Following steps analogous to those in Appendix C we obtain Eq. (19).

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