Cohering and decohering power of massive scalar fields under instantaneous interactions

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Employing a nonperturbative approach based on an instantaneous interaction between a two-level Unruh-DeWitt detector and a massive scalar field, we investigate the ability of the latter to generate or destroy coherence in the detector by deriving the cohering and decohering power of the induced quantum evolution channel. For a field in a coherent state a previously unobserved effect is reported in which the amount of coherence that the field generates displays a revival pattern with respect to the size of the detector. Extending previous results into the nonperturbative regime of an arbitrary coupling strength, it is demonstrated that in the case of a thermal field with a positive mass a detector initialized in a maximally coherent state experiences a smaller degree of decoherence after its interaction with the field, compared to the massless case. In both examples for a suitable choice of detector radius, field energy, and coupling strength it is possible to infer the mass of the field by either measuring the coherence present in the detector in the case of an interaction with a coherent field or the corresponding decoherence of a maximally coherent state in the case of a thermal field. In view of recent advances in the study of Proca metamaterials, these results suggest the possibility of utilizing the theory of massive electromagnetism for the construction of novel applications for use in quantum technologies.

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I. INTRODUCTION

Coherent systems, defined as those systems that can be described by a superposition of different states, form the backbone of the second quantum revolution brought about by the advent of quantum information science and technology [1,2]. Recast into the language of a quantum resource theory [3–8], it was shown that coherence is closely related to entanglement [9–11], another important resource which fuels applications, such as quantum dense coding [12], unhackable cryptography [13], and teleportation [14], for example. More recently an increasing amount of research has been focusing on the importance of coherence in quantum computing by studying its depletion during the execution of algorithms [15–21]. Coherence plays a central part in other physical contexts as well, such as in quantum metrology [22,23], thermodynamics [24–28], and even possibly in biological processes [29,30]. Because of its usefulness as a resource, it is, therefore, of particular interest to study the conditions under which coherence can be extracted or generated from other systems [31–34], as well as to devise methods for its protection [35–37] against the decohering effects of the environment [38–40].

In this paper we will examine the ability of a massive quantum field to generate or destroy coherence in a two-level Unruh-DeWitt (UDW) detector under an instantaneous interaction [41–47] (for an analysis of the amount of coherence present in the modes of the field under a different context see Refs. [48–51]). To accomplish this we will determine the cohering and decohering power of the quantum evolution channel induced by the action of the field on the detector [52–59], taking into consideration the size of the detector, the mass of the field, and its energy. Compared to other approaches that use perturbative methods to study coherence in a relativistic setting [33,34,60–64], a treatment based on an instantaneous interaction permits the exact solution of the final state of the detector for arbitrary coupling strengths. This provides us with the opportunity of uncovering novel effects otherwise hidden in the weak-coupling limit. An example is given in Sec. IV where it is observed that for specific values of the detector's radius, the amount of coherence generated by a coherent field vanishes, an effect which is absent in a perturbative setting [34].

The possibility of measuring the amount of coherence harvested by a detector from a coherent field as a means of probing the mass of axion dark matter has recently been proposed [64]. We show how, under a suitable choice of parameters, it is similarly possible to infer the mass of a scalar field by measuring the cohering power of the coherent field. In this case changes in coherence are easier to detect since they are orders of magnitude larger than what is possible with a weak interaction coupling.

The advantages in considering massive scalar fields become more apparent in Sec. V where the decohering power of a thermal field with inverse temperature β is presented. The ability of the field to preserve part of the coherence stored in a maximally coherent state of the detector is enhanced for increasing values of its mass. This observation is an extension

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of similar perturbative results about the coherent behavior of a pointlike atom immersed in a massive field [63] into the regime of an arbitrary coupling strength for a detector of any size and is in line with other reports on the advantages of massive fields in the processes of entanglement harvesting [65–67] and sensing [68–70]. Since decoherence presents a major hurdle in practical uses of quantum computing, such results may be of interest and could perhaps be leveraged with the use of massive electromagnetic fields in Proca metamaterials [71].

We begin our investigation by giving a short introduction to the resource theory of quantum coherence in Sec. II and the UDW detector model in Sec. III. As is common practice, throughout the paper we employ a natural system of units in which $\hbar = c = k_B = 1$.

II. COHERING AND DECOHERING POWER OF QUANTUM CHANNELS

Coherence, i.e., the degree of superposition of a quantum system [3,6,7] is dependent on the choice of basis of the underlying Hilbert space which we use to describe the state ρ of the system. For a state of the form

$$\rho = \sum_{i,j} \rho_{ij} |i\rangle\langle j|, \tag{1}$$

where $\{|i\rangle\}_{i=0}^{d-1}$ is a finite set of basis states spanning the d-dimensional Hilbert space $\mathbb{C}^d$, we say that ρ represents a *coherent state* if there exists, at least, one pair of indices $i \neq j$ such that $\rho_{ij} \neq 0$. A system which is *incoherent* is represented by a diagonal matrix and satisfies

$$\Delta(\rho) = \rho, \tag{2}$$

where

$$\Delta(\rho) = \sum_i \rho_{ii} |i\rangle\langle i| \tag{3}$$

denotes the *dephasing operation* in the chosen basis.

The set of quantum operations acting on a state is similarly divided into those that are capable and those that are incapable of creating coherence. The simplest set of incoherent operations, the so-called *maximally incoherent operations* (MIO) are defined as those completely positive and trace-preserving operations Φ that map the set of incoherent states $\mathcal{I}$ onto a subset of itself,

$$\Phi(\mathcal{I}) \subseteq \mathcal{I}. \tag{4}$$

The ability of a quantum channel Φ to generate coherence out of incoherent states can be determined by calculating its *cohering power* [52–55,59]. Before this can be defined it is necessary first to introduce the notion of a *coherence measure*. This is a non-negative real-valued function C on the set of density matrices with the following properties:

(i) $C(\rho) \geqslant 0$ with equality if and only if $\rho \in \mathcal{I}$.
(ii) $C(\Phi(\rho)) \leqslant C(\rho)$ for every $\Phi \in$ MIO.
(iii) $C(\sum_i p_i \rho_i) \leqslant \sum_i p_i C(\rho_i)$.

The first property requires the measure to be *faithful* so that it can distinguish between coherent and incoherent states. The second property reflects the restrictions of the theory. Since by definition MIOs cannot generate coherent out of incoherent states it makes sense to require the measure to be *monotonic*, the amount of coherence in a state after the action of a (MIO) operation should, therefore, always be less than before. This property is what gives the theory the structure of a *quantum resource* [72]. The final property, which requires the measure to be convex, states that it is not possible to increase the average amount of coherence in a quantum ensemble $\{p_i, \rho_i\}$, where p_i is the probability of obtaining state ρ_i by simply mixing its elements.

Armed with a valid measure of coherence, we are now in a position to define the cohering power of the channel as the maximum amount of coherence obtained by the action of Φ on the set of incoherent states,

$$\mathcal{C}(\Phi) = \max_{\rho \in \mathcal{I}} C(\Phi(\rho)). \tag{5}$$

Because of convexity the maximum on the right-hand side is actually reached by acting Φ on one of the basis states. This simplifies considerably the calculation since the required optimization is now performed over a discrete instead of a continuous set. In this case,

$$\mathcal{C}(\Phi) = \max_{|i\rangle} C(\Phi(|i\rangle\langle i|)). \tag{6}$$

Another property of interest for a quantum channel is the amount of coherence that it destroys when it is applied on a *maximally coherent state*, i.e., a uniform superposition, of the form

$$\psi_d(\boldsymbol{\theta}) = \frac{1}{\sqrt{d}} \sum_{j=0}^{d-1} e^{i\theta_j} |j\rangle. \tag{7}$$

Similar to Eq. (6) we now define the *decohering power* of the channel as the maximum possible difference in the amount of coherence before and after its action on the maximally coherent state,

$$\mathcal{D}(\Phi) = \max_{\boldsymbol{\theta}} [C(\psi_d(\boldsymbol{\theta})) - C(\Phi(\psi_d(\boldsymbol{\theta})))]. \tag{8}$$

In what follows we will employ the oft-used ℓ_1-norm of coherence as our measure. This is given by the sum of the absolute values of the nondiagonal elements of the density matrix,

$$C_{\ell_1}(\rho) = \sum_{i \neq j} |\rho_{ij}|. \tag{9}$$

For the set of maximally coherent states,

$$C_{\ell_1}(\psi_d(\boldsymbol{\theta})) = d - 1 \tag{10}$$

so in this case,

$$\mathcal{D}_{\ell_1}(\Phi) = d - 1 - \min_{\boldsymbol{\theta}} C_{\ell_1}(\Phi(\psi_d(\boldsymbol{\theta}))). \tag{11}$$

III. THE UNRUH-DEWITT DETECTOR MODEL

The UDW detector model is frequently employed as a means of studying the interaction between a two-level system (the detector) and a quantum field [73–75]. The interaction induces transitions between the detector's ground $|g\rangle$ and excited $|e\rangle$ states with an energy gap equal to Ω, which depend on the initial state of the field σ_ϕ as well as on the trajectory of the detector and the structure of the underlying spacetime.

Coupling the *monopole* operator of the detector,

$$\hat{\mu}(t) = e^{i\Omega t}|e\rangle\langle g| + e^{-i\Omega t}|g\rangle\langle e|, \tag{12}$$

to the field operator $\hat{\varphi}(t, \mathbf{x})$, evaluated at the detector's position $\mathbf{x}$ at time t, defines the UDW interaction Hamiltonian,

$$\hat{H}_{\rm int}(t) = \chi(t)\hat{\mu}(t) \otimes \hat{\varphi}(t, \mathbf{x}), \tag{13}$$

where the real-valued *switching function* function $\chi(t)$ describes the strength of the interaction at each instant in time. For a massive scalar field with mass m, the field operator in flat Miknowski spacetime is given by

$$\hat{\varphi}(t, \mathbf{x}) = \int \frac{d^3\mathbf{k}}{\sqrt{(2\pi)^3 2\omega(\mathbf{k})}} (\hat{a}_{\mathbf{k}} e^{i[\mathbf{k}\cdot\mathbf{x} - \omega(\mathbf{k})t]} + \text{H.c.}), \tag{14}$$

where $\hat{a}_{\mathbf{k}}$ and $\hat{a}^{\dagger}_{\mathbf{k}}$ denote the annihilation and creation operators, respectively, of a field mode with momentum $\mathbf{k}$ and energy $\omega(\mathbf{k}) = \sqrt{|\mathbf{k}|^2 + m^2}$ that satisfy the canonical commutation relations,

$$[\hat{a}_{\mathbf{k}}, \hat{a}_{\mathbf{k}'}] = [\hat{a}^{\dagger}_{\mathbf{k}}, \hat{a}^{\dagger}_{\mathbf{k}'}] = 0, \quad [\hat{a}_{\mathbf{k}}, \hat{a}^{\dagger}_{\mathbf{k}'}] = \delta(\mathbf{k} - \mathbf{k}'). \tag{15}$$

The UDW Hamiltonian describes a pointlike interaction in which the field interacts with the detector at a single point in space each time. By averaging over a region in a neighborhood of the detector's position it is possible to extend the model in order to take into account the finite size of the detector. For a detector at rest at position $\mathbf{x}$ [76], Eq. (13) is then replaced by

$$\hat{H}_{\rm int}(t) = \chi(t)\hat{\mu}(t) \otimes \int f(\mathbf{x} - \mathbf{x}')\hat{\varphi}(t, \mathbf{x}')d^3\mathbf{x}'. \tag{16}$$

The real-valued *smearing function* $f(\mathbf{x})$ with dimensions $(length)^{-3}$ reflects the shape and size of the detector [77–80] with a mean effective radius R equal to

$$R = \int |\mathbf{x}| f(\mathbf{x}) d^3\mathbf{x}. \tag{17}$$

By taking the pointlike limit $f(\mathbf{x}) = \delta(\mathbf{x})$, (i.e., $R \to 0$), Eq. (13) is immediately recovered. Setting,

$$F(\mathbf{k}) = \int f(\mathbf{x}) e^{-i\mathbf{k}\cdot\mathbf{x}} d^3\mathbf{x} \tag{18}$$

for the Fourier transform of the smearing function, we can rewrite Eq. (16) as

$$\hat{H}_{\rm int}(t) = \chi(t)\hat{\mu}(t) \otimes \hat{\varphi}_f(t, \mathbf{x}), \tag{19}$$

with a "smeared" field operator of the form

$$\hat{\varphi}_f(t, \mathbf{x}) = \int \frac{d^3\mathbf{k}}{\sqrt{(2\pi)^3 2\omega(\mathbf{k})}} [F(\mathbf{k})\hat{a}_{\mathbf{k}} e^{i[\mathbf{k}\cdot\mathbf{x} - \omega(\mathbf{k})t]} + \text{H.c.}]. \tag{20}$$

A. Evolution under an instantaneous interaction

In order to obtain the final state of the detector after the interaction has been switched off, we must first evolve the combined system of detector and field with the unitary operator $\hat{U}$ generated by the time integral of the interaction Hamiltonian,

$$\hat{U} = \mathcal{T} \exp\left(-i \int_{-\infty}^{+\infty} \hat{H}_{\rm int}(t) dt\right), \tag{21}$$

where $\mathcal{T}$ denotes the time-ordering operator. Tracing out the field degrees of freedom induces a quantum evolution channel on the initial-state ρ of the detector defined by

$$\Phi(\rho) = \text{tr}_{\varphi}[\hat{U}(\rho \otimes \sigma_{\varphi})\hat{U}^{\dagger}]. \tag{22}$$

Under a δ-switching function centered around t_0,

$$\chi(t) = \lambda\delta(t - t_0), \tag{23}$$

with λ as a coupling constant with the same dimensions as the length, it is possible to drop the time ordering in (21) [41–47]. In this case,

$$\hat{U} = \exp[-i\lambda\hat{\mu}_0 \otimes \hat{\varphi}_{f_0}], \tag{24}$$

where $\hat{\mu}_0 = \hat{\mu}(t_0)$ and $\hat{\varphi}_{f_0} = \hat{\varphi}_f(t_0, \mathbf{x})$. With a little bit of algebra, it is easy to show that since $\hat{\mu}_0^2 = I$, the evolution operator can be rewritten as

$$\hat{U} = \frac{I - \hat{\mu}_0}{2} \otimes \exp(i\lambda\hat{\varphi}_{f_0}) + \frac{I + \hat{\mu}_0}{2} \otimes \exp(-i\lambda\hat{\varphi}_{f_0}). \tag{25}$$

Inserting Eq. (25) into Eq. (22), we find that the action of the channel on the detector,

$$\Phi(\rho) = (1 - |z|)B(\rho) + |z|V\rho V^{\dagger} \tag{26}$$

is equal to a convex combination of a bit flip channel [1],

$$B(\rho) = \frac{\rho + \hat{\mu}_0\rho\hat{\mu}_0}{2}, \tag{27}$$

and a unitary rotation,

$$V = \sqrt{\frac{|z| + \text{Re}\, z}{2|z|}}\hat{I} - i\sqrt{\frac{|z| - \text{Re}\, z}{2|z|}}\hat{\mu}_0, \tag{28}$$

where

$$z = \text{tr}_{\varphi}[e^{i2\lambda\hat{\varphi}_{f_0}}\sigma_{\varphi}]. \tag{29}$$

B. Cohering and decohering power of scalar fields

According to Eqs (6) and (9), the ℓ_1-cohering power of the channel induced by the UDW interaction of the detector with the massive field is equal to the maximum amount of coherence obtained by acting Φ on either the ground or excited state of the detector. In both cases, this amount is the same and equal to

$$\mathcal{C}_{\ell_1}(\Phi) = \left|\left\langle \sin\left(2\lambda\hat{\varphi}_{f_0}\right)\right\rangle\right|, \tag{30}$$

where $\langle\hat{X}\rangle = \text{tr}_{\varphi}(\hat{X}\sigma_{\varphi})$ denotes the expectation value of field operator $\hat{X}$. Equation (30) is, in fact, equal to the maximum possible amount of coherence that can be obtained by acting Φ on any state of the detector (for details consult the Appendix).

To obtain the ℓ_1-decohering power requires a little more effort. Replacing the maximally coherent state,

$$\psi_2(\theta) = \frac{1}{\sqrt{2}}(|g\rangle + e^{i\theta}|e\rangle), \tag{31}$$

in Eq. (26) we see that the coherence of the final state of the detector is equal to

$$C_{\ell_1}(\Phi(\psi_2(\theta))) = \sqrt{\cos^2(\theta - \Omega t_0) + (\text{Re}\ z)^2 \sin^2(\theta - \Omega t_0)}. \tag{32}$$

For a maximally coherent state with $\theta = \Omega t_0$ the amount of coherence before and after the interaction has taken place is frozen [81]. For this choice of phase, the state is a fixed point of the evolution channel. This observation holds, in general, and is independent of details, such as the mass of the field, its initial state or the size of the detector.

It is straightforward now to show that the minimum in Eq. (32) is obtained by setting $\theta = \frac{\pi}{2} + \Omega t_0$. With the help of Eq. (11) we, therefore, find that the ℓ_1-decohering power of the field-induced channel is equal to

$$\mathcal{D}_{\ell_1}(\Phi) = 1 - \left|\left\langle \cos\left(2\lambda\hat{\varphi}_{f_0}\right)\right\rangle\right|. \tag{33}$$

We now proceed to study the cohering and decohering power of a field in a coherent and a thermal state, respectively.

IV. COHERING POWER OF COHERENT SCALAR FIELDS

A coherent state $|a\rangle$ of the field is described by a complex valued *coherent amplitude distribution* $a(\mathbf{k})$ such that the action of the annihilation operator $\hat{a}_{\mathbf{k}}$ on the state is equal to [41,82,83]

$$\hat{a}_{\mathbf{k}}|a\rangle = a(\mathbf{k})|a\rangle. \tag{34}$$

Let us now decompose the field into two parts,

$$\hat{\varphi}_{f_0} = \hat{a} + \hat{a}^\dagger, \tag{35}$$

each containing only annihilation or creation operators, respectively,

$$\hat{a} = \int \frac{d^3\mathbf{k}}{\sqrt{(2\pi)^3 2\omega(\mathbf{k})}} F(\mathbf{k})\hat{a}_{\mathbf{k}} e^{i[\mathbf{k}\cdot\mathbf{x}-\omega(\mathbf{k})t_0]}, \tag{36a}$$

$$\hat{a}^\dagger = \int \frac{d^3\mathbf{k}}{\sqrt{(2\pi)^3 2\omega(\mathbf{k})}} F^*(\mathbf{k})\hat{a}^\dagger_{\mathbf{k}} e^{-i[\mathbf{k}\cdot\mathbf{x}-\omega(\mathbf{k})t_0]}. \tag{36b}$$

By employing the Baker-Campbell-Hausdorff formula,

$$e^{\hat{X}+\hat{Y}} = e^{\hat{X}} e^{\hat{Y}} e^{-(1/2)[\hat{X},\hat{Y}]}, \tag{37}$$

which holds true when both $[\hat{X}, [\hat{X}, \hat{Y}]] = 0$ and $[\hat{Y}, [\hat{X}, \hat{Y}]] = 0$, it can be shown that

$$\begin{aligned}\langle e^{i2\lambda\hat{\varphi}_{f_0}}\rangle_a &= e^{-2\lambda^2[\hat{a},\hat{a}^\dagger]}\langle e^{i2\lambda\hat{a}^\dagger} e^{i2\lambda\hat{a}}\rangle_a. \\ &= e^{-2\lambda^2[\hat{a},\hat{a}^\dagger]} e^{i4\lambda\,\mathrm{Re}\,\langle\hat{a}\rangle_a}\end{aligned} \tag{38}$$

where

$$[\hat{a}, \hat{a}^\dagger] = \frac{1}{(2\pi)^3}\int \frac{|F(\mathbf{k})|^2}{2\omega(\mathbf{k})} d^3\mathbf{k}, \tag{39}$$

and a subscript in the expectation value of the field operator is included in order to indicate its dependence on the coherent amplitude distribution. From Eq. (30) it follows that the ℓ_1-cohering power of a coherent scalar field is equal to

$$\mathcal{C}_{\ell_1}(\Phi) = e^{-2\lambda^2[\hat{a},\hat{a}^\dagger]}|\sin(4\lambda\,\mathrm{Re}\langle\hat{a}\rangle_a)|. \tag{40}$$

Assuming a static detector with a Gaussian smearing function,

$$f(\mathbf{x}) = \frac{\exp\left[-\frac{4|\mathbf{x}|^2}{\pi R^2}\right]}{(\pi R/2)^3}, \tag{41}$$

and a corresponding Fourier transform of the form

$$F(\mathbf{k}) = \exp\left[-\frac{\pi|\mathbf{k}|^2R^2}{16}\right], \tag{42}$$

the commutator between $\hat{a}$ and $\hat{a}^\dagger$ is a function of the effective radius of the detector and the mass of the field,

$$\begin{aligned}[\hat{a}, \hat{a}^\dagger] &= \frac{1}{4\pi^2}\int_0^\infty \frac{k^2 e^{-\frac{\pi k^2R^2}{8}}}{\sqrt{k^2+m^2}} dk \\ &= \frac{m^2}{16\pi^{3/2}} U\left(\frac{3}{2}, 2, \frac{\pi m^2R^2}{8}\right),\end{aligned} \tag{43}$$

where

$$U(a, b, z) = \frac{2}{\Gamma(a)}\int_0^\infty e^{-zt^2} t^{2a-1}(1+t^2)^{b-a-1} dt \tag{44}$$

denotes *Tricomi's confluent hypergeometric function* [84] with $\Gamma(a)$ the Gamma function.

In order to further simplify calculations we will assume from now on that the interaction between the detector and the field takes place at time $t_0 = 0$ with the detector positioned at the origin of the coordinate system $\mathbf{x} = 0$, and that the coherent amplitude distribution of the field is given by a skewed-Gaussian function,

$$a(\mathbf{k}) = \sqrt{\frac{|\mathbf{k}|}{\omega(\mathbf{k})}} \frac{\exp\left(-\frac{2|\mathbf{k}|^2}{\pi E^2}\right)}{(\pi E/2)^{3/2}}, \tag{45}$$

with mean energy E equal to the expectation value of the field Hamiltonian,

$$\hat{H}_\phi = \int \omega(\mathbf{k})\hat{a}^\dagger_{\mathbf{k}}\hat{a}_{\mathbf{k}} d^3\mathbf{k}. \tag{46}$$

With these assumptions in hand, the expectation value of the real part of Eq. (36a) which now depends on the mean effective radius of the detector, the mass, and the mean energy of the field is equal to

$$\begin{aligned}\mathrm{Re}\,\langle\hat{a}\rangle_a &= \sqrt{\frac{8}{\pi^4E^3}}\int_0^\infty \frac{k^{5/2}}{\sqrt{k^2+m^2}} \exp\left[-\frac{k^2}{2\sigma^2}\right] dk, \\ &= m\sqrt{\frac{2m^3}{\pi^4E^3}}\Gamma\left(\frac{7}{4}\right) U\left(\frac{7}{4}, \frac{9}{4}, \frac{m^2}{2\sigma^2}\right),\end{aligned} \tag{47}$$

where for ease of notation, we introduce the following quantity with units of inverse square mass:

$$\frac{1}{\sigma^2} = \frac{4}{\pi E^2} + \frac{\pi R^2}{8}. \tag{48}$$

In Figs. 1(a)–1(c), we present for a detector with a fixed radius the ℓ_1-cohering power of the field as a function of its mean energy and the strength of its interaction with the detector. As is to be expected, in the trivial cases where there is no interaction with the field ($\lambda = 0$) and for an incoherent field in the vacuum state ($E = 0$), the cohering power is zero. We observe that a massless field is generally better at generating coherence since the maximum amount that can be obtained decreases for larger field masses. When the Compton wavelength of the field is much smaller than the radius of the

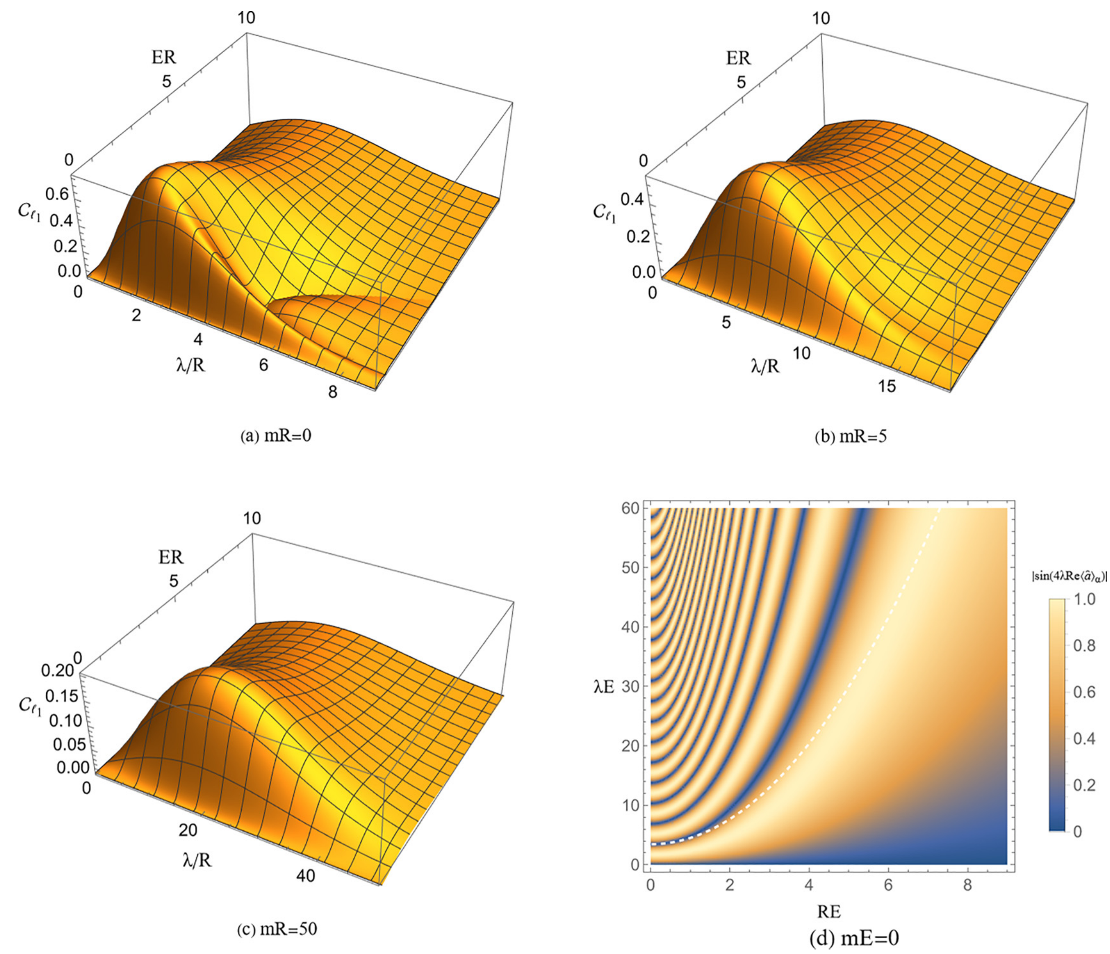

FIG. 1. (a)–(c) ℓ_1-cohering power of a massive scalar field in a coherent state interacting instantaneously with a static Unruh-DeWitt detector with a fixed radius R as a function of the energy of field E and the coupling strength λ for three different values of the field's mass. (d) Oscillatory behavior of ℓ_1-cohering power for a massless field with a fixed energy E. The upper bound for the coupling strength such that no coherence revival can occur is given by the dashed line.

detector,

$$R \gg \frac{2\pi}{m}, \tag{49}$$

the phase of the oscillating term in Eq. (40) becomes very small and tends to zero in the asymptotic limit $mR \to \infty$. In this case $\langle e^{i2\lambda\hat{\varphi}_{f_0}}\rangle_a = 1$, so the field has no effect on the detector since the induced quantum channel is equal to the identity operator $\Phi(\rho) = \rho$ which has zero cohering power.

In the other limit of a very strong coupling or equivalently a pointlike detector,

$$\lambda \gg R, \tag{50}$$

the exponential term in Eq. (40) now takes small values and vanishes when either $\lambda \to \infty$ or $R \to 0$. In this case $\langle e^{i2\lambda\hat{\varphi}_{f_0}}\rangle_a = 0$ and the induced quantum channel acts as the bit flip channel $\Phi(\rho) = B(\rho)$ again with a zero cohering power [52].

Due to its oscillatory behavior, the cohering power of the field displays a revival pattern with respect to the radius of the detector. Because of the exponential damping in Eq. (40) this pattern is hard to spot by looking directly at the cohering power but is easily discernible once this factor is ignored as in Fig. 1(d), for example. If one wishes to utilize a scalar field to generate coherence in a detector one, therefore, needs to take into account its size. By demanding that $4\lambda\,\mathrm{Re}\,\langle\hat{a}\rangle_a \leqslant \pi$, and noting that for a fixed mean-field energy and detector radius the integral on the right-hand side of Eq. (47) is maximized when $m = 0$, it can easily be shown that the region of the parameter space for which no revival pattern can occur is given by

$$\frac{\lambda E}{\pi} \leqslant \frac{(2\pi)^{\frac{3}{4}}}{\Gamma(1/4)}\left[1 + \frac{\pi^2 E^2 R^2}{32}\right]^{\frac{5}{4}}. \tag{51}$$

In Fig. 2, we demonstrate the dependence of the cohering power on the field mass for different values of the coupling constant for a field with a fixed mean energy E and a detector with a mean effective radius equal to $R = 1/E$. Replacing π with $\pi/2$ on the denominator of the left-hand side of Eq. (51), it follows that when $\lambda E \lesssim 2.4$ the cohering power of the field is in a one-to-one correspondence with its mass. The value of the cohering power in this case could, therefore, be employed as a method of probing the value of the field's mass.

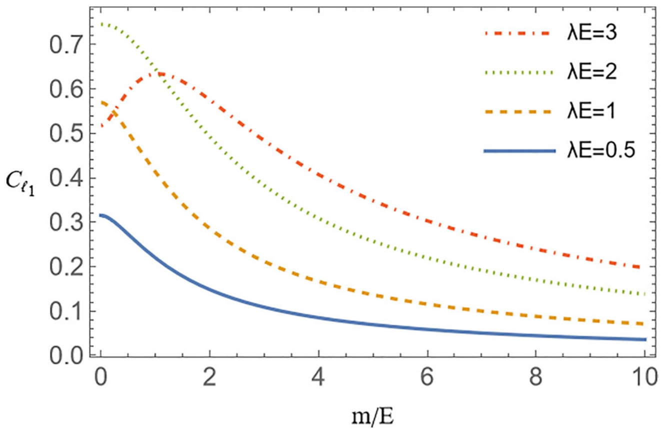

FIG. 2. Dependence of ℓ_1-cohering power of the field as a function of mass for a fixed mean-field energy E and a detector with mean radius equal to $R = 1/E$.

V. DECOHERING POWER OF THERMAL FIELDS

For a thermal field at an inverse temperature β,

$$\sigma_\phi = \frac{e^{-\beta \hat{H}_\phi}}{Z}, \tag{52}$$

with partition function $Z = \mathrm{tr}_\varphi e^{-\beta \hat{H}_\phi}$, let $\langle \hat{X} \rangle_\beta$ denote the dependence of the expectation value of field operator $\hat{X}$ on the temperature. Employing the same decomposition as in Eq. (35), it can be shown that in this case,

$$\langle e^{i2\lambda \hat{\varphi}_{f_0}} \rangle_\beta = e^{-2\lambda^2 \langle \hat{\varphi}^2_{f_0} \rangle_\beta}. \tag{53}$$

To see this we must first rewrite the left-hand side following the same steps that led to the derivation of Eq. (38),

$$\langle e^{i2\lambda \hat{\varphi}_{f_0}} \rangle_\beta = e^{-2\lambda^2 [\hat{a}, \hat{a}^\dagger]} \langle e^{i2\lambda \hat{a}^\dagger} e^{i2\lambda \hat{a}} \rangle_\beta. \tag{54}$$

To compute the expectation value on the right-hand side, we now Taylor expand $e^{2i\lambda \hat{a}}$ and $e^{2i\lambda \hat{a}^\dagger}$ to obtain

$$\langle e^{i2\lambda \hat{a}^\dagger} e^{i2\lambda \hat{a}} \rangle_\beta = \sum_{m,m'=0}^{\infty} \frac{(i2\lambda)^{m+m'}}{(m!)(m'!)} \langle (\hat{a}^\dagger)^m (\hat{a})^{m'} \rangle_\beta. \tag{55}$$

Because the field is diagonal in the energy basis we only need consider terms where $m = m'$ since any other term will be equal to zero. We will now show that

$$\langle (\hat{a}^\dagger)^m (\hat{a})^m \rangle_\beta = m! (\langle \hat{a}^\dagger \hat{a} \rangle_\beta)^m. \tag{56}$$

With the help of the following identity:

$$e^{\hat{X}} \hat{Y} e^{-\hat{X}} = \hat{Y} + [\hat{X}, \hat{Y}] + \frac{1}{2!}[\hat{X}, [\hat{X}, \hat{Y}]] + \frac{1}{3!}[\hat{X}, [\hat{X}, [\hat{X}, \hat{Y}]]] \cdots, \tag{57}$$

and the commutation relation between $\hat{a}_{\mathbf{k}}$ and the field Hamiltonian,

$$[\hat{a}_{\mathbf{k}}, \hat{H}_\varphi] = \omega(\mathbf{k}) \hat{a}_{\mathbf{k}}, \tag{58}$$

we find that

$$\hat{a}_{\mathbf{k}} e^{-\beta \hat{H}_\varphi} = e^{-\beta \omega(\mathbf{k})} e^{-\beta \hat{H}_\varphi} \hat{a}_{\mathbf{k}}. \tag{59}$$

Using this and Eq. (15), it is straightforward to show that

$$\langle \hat{a}^\dagger_{\mathbf{k}} \hat{a}_{\mathbf{k}'} \rangle_\beta = \frac{e^{-\beta\omega(\mathbf{k}')}}{1 - e^{-\beta\omega(\mathbf{k}')}} \delta(\mathbf{k} - \mathbf{k}'), \tag{60}$$

which implies by induction that

$$\left\langle \prod_{i=1}^{m} \hat{a}^\dagger_{\mathbf{k}_i} \prod_{j=1}^{m} \hat{a}_{\mathbf{k}'_j} \right\rangle_\beta = \sum_{i=1}^{m} \langle \hat{a}^\dagger_{\mathbf{k}_i} \hat{a}_{\mathbf{k}'_m} \rangle_\beta \left\langle \prod_{i' \neq i} \hat{a}^\dagger_{\mathbf{k}'_{i'}} \prod_{j=1}^{m-1} \hat{a}_{\mathbf{k}'_j} \right\rangle_\beta. \tag{61}$$

It follows now that

$$\langle (\hat{a}^\dagger)^m (\hat{a})^{m'} \rangle_\beta = m \langle \hat{a}^\dagger \hat{a} \rangle_\beta \langle (\hat{a}^\dagger)^{m-1} (\hat{a})^{m-1} \rangle_\beta, \tag{62}$$

from which Eq. (56) can be obtained recursively. Finally,

$$\langle e^{i2\lambda \hat{a}^\dagger} e^{i2\lambda \hat{a}} \rangle_\beta = e^{-4\lambda^2 \langle \hat{a}^\dagger \hat{a} \rangle_\beta} = e^{2\lambda^2 [\hat{a}, \hat{a}^\dagger]} e^{-2\lambda^2 \langle \hat{\varphi}^2_{f_0} \rangle_\beta}, \tag{63}$$

which completes the proof.

Looking back at Eq. (53) and noting that $\mathcal{C}_{\ell_1}(\Phi) = |\mathrm{Im} z|$ we observe, perhaps unsurprisingly, that a thermal field is incapable of generating coherence through an instantaneous interaction. On the other hand, its decohering power is equal to

$$\mathcal{D}_{\ell_1}(\Phi) = 1 - e^{-\lambda^2 I(\beta)}, \tag{64}$$

where

$$I(\beta) = \frac{1}{(2\pi)^3} \int \frac{|F(\mathbf{k})|^2}{\omega(\mathbf{k})} \coth\left(\frac{\beta\omega(\mathbf{k})}{2}\right) d^3\mathbf{k} = \frac{1}{2\pi} \int_0^\infty \frac{k^2 e^{-\frac{\pi k^2 R^2}{8}}}{\sqrt{k^2 + m^2}} \coth\left(\frac{\beta\sqrt{k^2 + m^2}}{2}\right) dk. \tag{65}$$

Since

$$\frac{\partial I(\beta)}{\partial R} < 0, \quad \frac{\partial I(\beta)}{\partial m} < 0, \quad \frac{\partial I(\beta)}{\partial \beta} < 0, \tag{66}$$

the decohering power decreases for increasing values of the detector's radius and the mass of the field, whereas it increases with temperature. In Figs. 3(a) and 3(b), we present the ℓ_1-decohering power of a thermal field as a function of the detector's mean radius and the temperature of the field for a detector with the same Gaussian smearing function as in Eq. (41). It can be seen that compared to a massless one a massive field performs better at preserving the coherence in the detector. The loss of coherence of a qubit due to its interaction with the environment is a major obstacle in quantum computing. By making a two-level system interact with a massive instead of a massless field (such as an electromagnetic field in plasma [85], a waveguide [86], or a Proca metamaterial [71] for example) it might be possible to protect against this form of decoherence.

As in the previous case of a coherent field, in the limit of a detector much larger than the Compton wavelength of the field, effectively no interaction takes place between the two, so the detector loses no coherence and the decohering power of the field is equal to zero, whereas in the limit of a pointlike detector or a very strong coupling the quantum channel reduces to the bit flip channel with unit decohering power. The same conclusion also holds true when the temperature of the field takes very large values,

$$\beta \ll \lambda. \tag{67}$$

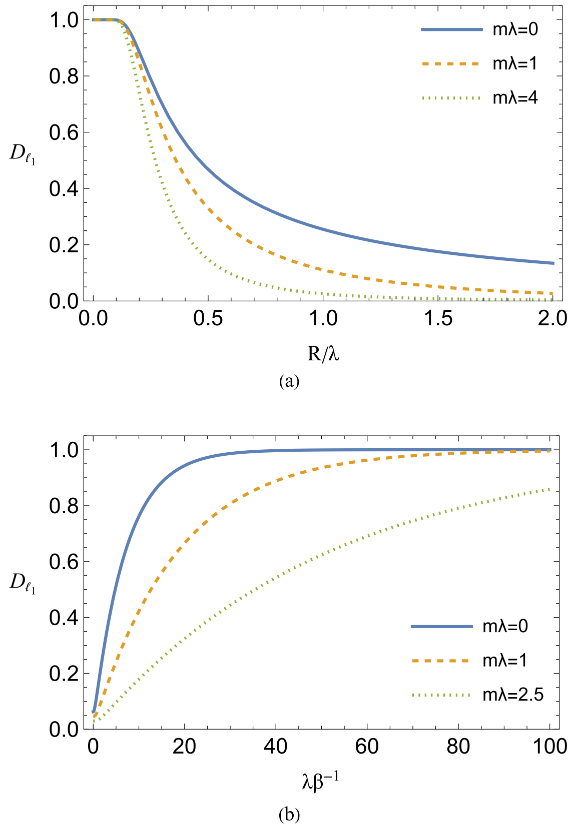

FIG. 3. ℓ_1-decohering power of a massive scalar field with respect to (a) the radius of the detector for a field with inverse temperature $\lambda\beta^{-1} = 2$ and (b) with respect to temperature for a detector with radius equal to $R/\lambda = 1$.

VI. DISCUSSION

Employing an instantaneous interaction between a two-level UDW detector and a massive scalar field, we investigated the ability of the field to generate or destroy coherence in the detector. This nonperturbative approach permits an exact examination of the effects that different parameters (such as the strength of the coupling constant, the size of the detector, the energy of the field, or its temperature, for example) have on the cohering and decohering power of the induced quantum evolution channel.

In the case of coherence generation by a coherent field state it was demonstrated that the success of the process depends on the size of the detector. More specifically, apart from the pointlike limit $R \ll \lambda$ where the mean radius of the detector is much smaller than the coupling strength and the macroscopic limit $R \gg \frac{2\pi}{m}$ of a detector much larger than the Compton wavelength of the field, there exist nontrivial values of the detector's radius for which it is impossible to generate any amount of coherence between its energy levels. This phenomenon demonstrates how the size of the system, which we wish to bring into a superposition of states, needs to be taken into consideration. For a suitable choice of detector radius, field energy, and coupling strength it is also possible to infer the mass of the field by measuring the amount of coherence present in the detector. Massive fields have been previously employed for distinguishing the kinematic state of a detector [68] for determining the distance of closest approach between two accelerating detectors [70] or for probing the mass of axion dark matter [64]. In all of these cases the quantity of interest under study each time is very small since it it is of the same order as the coupling constant or less. Even though instantaneous interactions are based on δ- switching functions and may appear as an idealization at a first glance, they can nonetheless be obtained from a Gaussian switching with an interaction duration which is much shorter than some characteristic time interval [87]. As we have demonstrated, in this case changes in coherence are no longer of the same order as the coupling constant, so they are easier to observe. It is expected that the above results directly apply in the generation of other quantum resources from the field as in entanglement harvesting [88–91], for example.

By calculating the decohering power of a thermal field, we also investigated the degradation caused by the field on the amount of coherence that is initially stored in a maximally coherent state of the detector. It was shown that for fixed values of the detector's radius, the field's temperature, and the coupling constant between the two, a massive field is better at preserving coherence. These results are a direct extension of similar perturbative observations [63] and suggest that massive fields could be employed in protecting against decoherence even in the case of a detector extended in space for an arbitrary coupling constant.

It is known that the UDW Hamiltonian contains all of the essential features of the interaction of matter with an electromagnetic field [79,80] where the analog of a massive field in this case is a Proca field [92]. Since massive electromagnetic theory can be realized as Maxwell theory in Proca metamaterials [71], studying the effects of mass on the generation and protection of coherence in a two-level system could potentially lead to the construction of novel technologies, such as new types of quantum memories, communication channels, and sensors. For this reason, a more complete investigation of cohering and decohering effects, of detectors interacting with massive fields, by making use of other nonperturbative methods [93,94] permitting a full dynamical analysis are certainly worth pursuing.

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APPENDIX: GENERALIZED COHERING POWER

Instead of the cohering power, one could enquire whether it is possible to harness the coherence already present in a state in order to obtain a greater amount of coherence from the action of a quantum operation, than what would otherwise be possible by only using incoherent states. In this case, one needs to specify the *generalized cohering power* of the channel defined by

$$\hat{\mathcal{C}}(\Phi) = \max_{\rho}[C(\Phi(\rho)) - C(\rho)]. \tag{A1}$$

It is obvious that $\mathcal{C}(\Phi) \leqslant \hat{\mathcal{C}}(\Phi)$. Depending on the choice of C as a coherence measure, there exist channels such that the cohering power is strictly smaller than its generalized definition. For the ℓ_1-norm of coherence in Eq. (9), it was shown that for channels acting on qubits [53],

$$\mathcal{C}_{\ell_1}(\Phi) = \hat{\mathcal{C}}_{\ell_1}(\Phi), \tag{A2}$$

so this is, in fact, the maximum possible amount that can be obtained by the action of Φ.

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