

Bose Polarons as Relativistic Unruh-DeWitt Detectors: Entanglement Harvesting from Bose-Einstein Condensates

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We show that a bound impurity in a Bose-Einstein condensate can be directly mapped to an Unruh-DeWitt detector interacting with a relativistic quantum field. We provide explicit experimental parameters for an implementation using ^{41}K impurities coupled to a ^{87}Rb condensate via finite-time Feshbach tuning. As an application, we study the extraction of vacuum entanglement from distant regions of the condensate and find viable parameters for the implementation of entanglement harvesting.

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Introduction—The advent of quantum field simulators capable of reproducing aspects of quantum field dynamics in curved spacetime has created unprecedented opportunities to test predictions of relativistic quantum field theory (QFT) [1–3]. A growing number of experimental platforms—ranging from ultracold atoms; trapped ions [4]; photonic [5], optomechanical [6,7], and superconducting architectures [8] to polariton fluids [9]—can engineer effective metrics, horizons, and nontrivial dynamics [10].

Consequently, phenomena long viewed as purely theoretical, such as horizon-induced correlations [11], particle production in time-dependent backgrounds [1,12], and entropy in QFT [13], are becoming experimentally accessible. This opens a path toward using laboratory systems to probe regimes that lie far beyond the reach of astrophysical or cosmological measurements.

The key feature of all experimental systems so far is the ability to detect spatial correlations of analog relativistic quantum fields. However, implementations of probes that can couple to the field locally in space and time are still missing. Local measurements are essential for probing fundamental features of QFTs that become evident only in finite spacetime regions, such as properties of local algebras [14,15], covariant measurement schemes [16,17], and, relevant for this work, entanglement between non-complementary regions [18–20].

Local probes in QFT are usually implemented as Unruh-DeWitt (UDW) detectors [21–23]—localized quantum systems that couple to a quantum field [24,25]. For instance,

UDW detectors can be implemented in practice with electro-optic sampling [26,27] and superconducting qubits [28], and proposals for interferometric schemes in Bose-Einstein condensates have been put forward [29,30].

When coupled for finite times, UDW detectors can implement quantum information protocols that fully utilize the local degrees of freedom of quantum fields [31–34]. In particular, they can quantify entanglement in the vacuum of relativistic QFTs through entanglement harvesting [35–38]. In the protocol, yet to be implemented experimentally, causally disconnected probes extract entanglement from a quantum field.

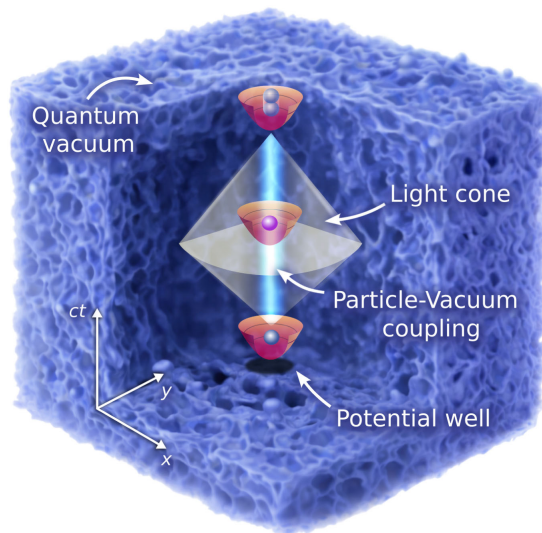


FIG. 1. A bound Bose polaron (UDW detector) in a harmonic trap couples for a finite time with a Bose-Einstein condensate through temporal Feshbach tuning. The detector probes local BEC degrees of freedom in a causal diamond determined by the speed of sound. After the interaction, the impurity accesses the BEC's quantum information content.

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Here, we show that a bound Bose polaron—a trapped impurity immersed in a Bose-Einstein condensate (BEC) [39–42]—realizes a microscopic implementation of an Unruh-DeWitt detector coupled for finite times to the effective relativistic phonon field of the condensate (see Fig. 1). Ultracold gases offer pristine control over motional degrees of freedom, i.e., spatial mode structure of the detector, and here we show how one can realize independent temporal control of the impurity-phonon coupling (see also Refs. [43,44]). This realizes a spatiotemporally localized detector suitable for probing quantum-information-theoretic properties such as entanglement. Our microscopic derivation expresses the detector gap, coupling strength, and switching function in terms of experimentally controllable parameters, providing concrete prescriptions for probing local QFT effects in current ultracold atom setups.

Local probes in QFT—The degrees of freedom of relativistic QFTs are spread over spacetime regions that are fully determined by information at a spatial surface, the so-called causal diamonds [45–47]. Due to this fact, long-time interactions effectively probe degrees of freedom localized in a spatial region much larger than the probe itself. Local effects in QFT can then be seen only when probing the field for finite times [34], as UDW detectors do.

In its simplest form, a UDW detector consists of a two-level system linearly coupled to a free scalar quantum field $\hat{\phi}(\mathbf{x})$ in a background spacetime $\mathcal{M}$. Here, we focus on $3+1$ -dimensional Minkowski spacetime with an inertial probe that couples to the conjugate momentum $\hat{\pi}(\mathbf{x}) = \partial_t \hat{\phi}(\mathbf{x})$ [48–50], where $\mathbf{x} = (t, \mathbf{x})$ denotes inertial coordinates. The coupling between the probe and the field $\hat{\pi}(\mathbf{x})$ is described by the interaction Hamiltonian density [51]

$$\hat{\mathcal{H}}_I(\mathbf{x}) = \lambda[\Lambda^*(\mathbf{x})e^{i\Omega t}\hat{\sigma}^+ + \Lambda(\mathbf{x})e^{-i\Omega t}\hat{\sigma}^-]\hat{\pi}(\mathbf{x}), \quad (1)$$

where λ is a coupling constant, $\hbar\Omega$ is the energy gap between the detector’s ground and excited states, and $\Lambda(\mathbf{x}) = \chi(t)f(\mathbf{x})$ is a spacetime smearing function that factors as a (dimensionless) real switching function $\chi(t)$ and a (potentially complex) smearing function $f(\mathbf{x})$ with units of density. $\Lambda(\mathbf{x})$ determines the region of spacetime where the probe couples to the field, and $e^{\pm i\Omega t}\hat{\sigma}^\pm$ are the time-evolved raising and lowering operators in the detector Hilbert space.

The UDW model in curved spacetimes and noninertial trajectories has been proven effective for studying effects such as Hawking radiation [52–54] and the Unruh effect [21,55–58]. Moreover, the model reproduces the essential features of the interactions of quantum fields coupled to any bosonic field, such as atoms or superconducting qubits coupled to electromagnetism [59–61] and quantum systems coupled to linearized quantum gravity [62–65]. However, implementing finite-time couplings is a major challenge in realizing UDW detectors experimentally [66].

A two-level system that interacts with the field according to (1) samples field degrees of freedom localized within the

causal diamond that contains the support of $\Lambda(\mathbf{x})$, allowing direct access to local observables in QFT [16]. Indeed, to leading order in λ , the detector’s final state depends only on the two-point function of $\hat{\pi}(\mathbf{x})$ smeared against $\Lambda(\mathbf{x})$. Explicitly, the vacuum two-point function of the operator $\hat{\pi}(\mathbf{x})$ is [67]

$$\langle \hat{\pi}(\mathbf{x})\hat{\pi}(\mathbf{x}') \rangle = \frac{1}{(2\pi)^3} \int d^3\mathbf{k} \frac{\hbar\omega_{\mathbf{k}}}{2} e^{ik \cdot (\mathbf{x}-\mathbf{x}')}, \quad (2)$$

where $\mathbf{k} \cdot \mathbf{x} = -\omega_{\mathbf{k}}t + \mathbf{k} \cdot \mathbf{x}$, $\omega_{\mathbf{k}} = c|\mathbf{k}|$, and c is the speed of light.

Phonons as relativistic quantum field—Since the seminal paper by Unruh on experimental black hole evaporation [68], numerous theoretical and experimental results building on the analogy between propagation of fields in curved spacetimes and propagation of sound waves in fluids have been obtained [10]. Especially, atomic superfluids are an ideal platform to simulate quantum fields, as low-energy phonons behave as a quantum massless scalar field in a Lorentzian metric determined by the background condensate [69,70]. By splitting the atomic quantum field as a classical background and quantum fluctuations, $\hat{\Psi}(\mathbf{x}) = \Psi_0(\mathbf{x}) + \delta\hat{\Psi}(\mathbf{x})$ [71], the BEC density can be written (to leading order in the fluctuations) as

$$\hat{\rho}_b(\mathbf{x}) = \hat{\Psi}^\dagger(\mathbf{x})\hat{\Psi}(\mathbf{x}) \approx \rho_0(\mathbf{x}) + \delta\hat{\rho}(\mathbf{x}). \quad (3)$$

Here, $\rho_0 = \Psi_0^\dagger\Psi_0$ is the (classical) background density and $\delta\hat{\rho}(\mathbf{x})$ encodes its quantum fluctuations. In the low-energy regime, $\delta\hat{\rho}$ behaves analogously to the momentum of a real scalar field. This quantum field evolves in a spacetime with Lorentzian metric where the effective speed of light is given by the speed of sound in the BEC, given by $c_s^2 = \rho_0 g_{bb}/m_b$ with $g_{bb} = 4\pi\hbar^2 a_{bb}/m_b$. Here, a_{bb} is the s -wave scattering length characterizing the boson-boson interaction and m_b is the mass of the boson. In particular, when the density of the condensate is constant, $c_s(\mathbf{x}) = c_s$, and the effective spacetime is Minkowski. The density fluctuations $\delta\hat{\rho}(\mathbf{x})$ can then be written as [72,73]

$$\delta\hat{\rho}(\mathbf{x}) = \frac{1}{\sqrt{g_{bb}}V} \sum_{\mathbf{k} \neq 0} \sqrt{\frac{\hbar\omega_{\mathbf{k}}}{2}} \left(\hat{b}_{\mathbf{k}} e^{ik \cdot \mathbf{x}} + \hat{b}_{\mathbf{k}}^\dagger e^{-ik \cdot \mathbf{x}} \right), \quad (4)$$

with V the normalization volume, $\hat{b}_{\mathbf{k}}, \hat{b}_{\mathbf{k}}^\dagger$ the phonon creation and annihilation operators, $\mathbf{k} \cdot \mathbf{x} = -\omega_{\mathbf{k}}t + \mathbf{k} \cdot \mathbf{x}$, and the relativistic dispersion relation $\omega_{\mathbf{k}} \approx c_s|\mathbf{k}|$ holds in the low-energy regime. The two-point function of the density fluctuations is

$$\langle \delta\hat{\rho}(\mathbf{x})\delta\hat{\rho}(\mathbf{x}') \rangle = \frac{1}{g_{bb}V} \sum_{\mathbf{k} \neq 0} \frac{\hbar\omega_{\mathbf{k}}}{2} e^{ik \cdot (\mathbf{x}-\mathbf{x}')}, \quad (5)$$

analogous to the two-point function of the momentum operator (2) in a finite volume [74], with the additional factor of g_{bb}^{-1} .

Impurities as local probes—Immersing a distinguishable particle of a different atomic species in Bose-Einstein condensates has been studied in detail and is connected to polarons in solid state physics [40]. Here, the additional control to trap the impurity in a harmonic trap gives rise to a bound Bose polaron. This has been studied experimentally with a mixture of lithium (impurity) and sodium (BEC), revealing decoherence of motional degrees of freedom [75] and the analog of the Lamb shift due to the coupling to the phononic quantum field [76]. In the weakly interacting limit, which we exploit here, the Bose polaron is known as the Fröhlich polaron [77–79]. Effectively, the impurity interaction with the BEC is a density-density coupling controlled by the parameter $g_{ab} = 2\pi\hbar^2 a_{ab}/\mu_{ab}$. Here, $\mu_{ab} = m_a m_b/(m_a + m_b)$ is the reduced mass, m_a is the impurity atom's mass, and a_{ab} is the interspecies s -wave scattering length. Explicitly, the interaction Hamiltonian density (in the interaction picture) is

$$\hat{\mathcal{H}}_I(\mathbf{x}) = g_{ab}\hat{\rho}_a(\mathbf{x})\hat{\rho}_b(\mathbf{x}), \quad (6)$$

where $\hat{\rho}_a(\mathbf{x})$ is the impurity's time-evolved density. Denoting the localized eigenstate of the impurity atom with energy $E_n = \hbar\Omega_n$ by $|n\rangle$ with wave function $\psi_n(\mathbf{x}) = \langle \mathbf{x} | n \rangle$, we have

$$\hat{\rho}_a(\mathbf{x}) = \sum_{nm} \psi_n^*(\mathbf{x})\psi_m(\mathbf{x})e^{i(\Omega_n - \Omega_m)t}|n\rangle\langle m|. \quad (7)$$

Assuming that the only relevant energy levels in Eq. (7) are its ground (g) and excited (e) states, the two-level impurity field operator becomes

$$\begin{aligned} \hat{\rho}_a(\mathbf{x}) = & e^{i\Omega t}(\psi_e^*\psi_g)(\mathbf{x})\hat{\sigma}^+ + e^{-i\Omega t}(\psi_g^*\psi_e)(\mathbf{x})\hat{\sigma}^- \\ & + |\psi_e(\mathbf{x})|^2\hat{\sigma}^+\hat{\sigma}^- + |\psi_g(\mathbf{x})|^2\hat{\sigma}^-\hat{\sigma}^+, \end{aligned} \quad (8)$$

where $\Omega = \Omega_e - \Omega_g$ and $\hat{\sigma}^\pm$ are the raising and lowering operators between $|g\rangle$ and $|e\rangle$. Since the two diagonal terms in (8) do not contribute to the impurity's excitation probability [22], we drop them in the following. Analogously, in (6) we can replace $\hat{\rho}_b(\mathbf{x})$ with $\delta\hat{\rho}(\mathbf{x})$, as ρ_0 is a classical term proportional to the identity. Remarkably, the polaron model also allows for a time-varying coupling $g_{ab}(t)$ to be implemented, leading to the interaction Hamiltonian

$$\hat{\mathcal{H}}_I(\mathbf{x}) = \bar{g}_{ab}\chi(t)(f^*(\mathbf{x})e^{i\Omega t}\hat{\sigma}^+ + f(\mathbf{x})e^{-i\Omega t}\hat{\sigma}^-)\delta\hat{\rho}(\mathbf{x}), \quad (9)$$

where $f(\mathbf{x}) = (\psi_g^*\psi_e)(\mathbf{x})$, and we write $g_{ab}(t) = \bar{g}_{ab}\chi(t)$, factoring out the time dependence in a switching function $-1 \leq \chi(t) \leq 1$. Comparing (9) with (1) shows that the polaron is effectively a UDW detector for the phonon density field. Introducing a timescale T that controls the switching of the interaction through $\chi(t) = \beta(t/T)$, the momentum-coupled UDW detector can be mapped to the polaron model by

$$\hat{\pi}(\mathbf{x}) \mapsto \sqrt{\frac{c_s}{c}}\sqrt{g_{bb}}\delta\hat{\rho}(\mathbf{x}), \quad \lambda \mapsto \sqrt{\frac{c}{c_s}}\frac{\bar{g}_{ab}}{\sqrt{g_{bb}}} \quad (10)$$

and the rescalings $\Omega \mapsto (c/c_s)\Omega$ and $T \mapsto (c_s/c)T$.

Arguably, allowing finite-time couplings is the most important feature of an UDW detector. The control of $g_{ab}(t)$ is achieved by magnetically tuned Feshbach resonances [80]. Most distinct pairs of atomic species allow a value of the magnetic field B_0 such that scattering length $a_{ab}(B_0) = 0$, implying $g_{ab} = 0$, thus decoupling the impurity and the BEC. Applying a magnetic pulse to the system by varying the external magnetic field $B(t)$ away from B_0 implements a finite time coupling, in perfect analogy with the detector model in (1).

We can explicitly see how the impurity can measure localized field observables by noticing that its vacuum excitation probability (see Supplemental Material [81]) is

$$\mathcal{P} = \frac{\bar{g}_{ab}^2}{\hbar^2} \langle \delta\hat{\rho}(\Lambda^-)\delta\hat{\rho}(\Lambda^+) \rangle, \quad (11)$$

where $\delta\hat{\rho}(\Lambda^\pm) \equiv \int d^4\mathbf{x}\Lambda^\pm(\mathbf{x})\delta\hat{\rho}(\mathbf{x})$ are smeared field observables localized by $\Lambda^\pm(\mathbf{x}) = \chi(t)f(\mathbf{x})e^{\pm i\Omega t}$. The shape of the wave function $(\psi_g^*\psi_e)(\mathbf{x})$ determines the spatial localization of the observables probed and $\chi(t)$ their extension in time, encoding expected values of local QFT observables in directly accessible physical quantities.

Experimental platform—The experimental system of ultracold mixtures offers a broad range of control parameters, which we discuss in the context of UDW detectors and relativistic quantum fields. As a suitable combination of species, we propose a potassium-rubidium mixture [82], where species-selective trapping of ^{41}K in a light-shift potential realizes the trapped impurity, while ^{87}Rb acts as the substrate for the phononic relativistic quantum field [70]. The large fine-structure splitting of rubidium allows for the realization of species-selective dipole potentials utilizing the existence of a tune-out wavelength [83], i.e., the light-shift potential is seen only by potassium atoms. The trap frequencies, realized with optical tweezers [84,85] or large-period optical lattices [86], can be tuned from a few hundred Hz to several tens of kHz.

The s -wave scattering length for rubidium, $a_{bb} = 100.4(1)a_0$, where a_0 is the Bohr radius, combined with the low three-body loss rate, allows for the preparation of large and high-density Bose-Einstein condensates [87]. For the reported peak density $\rho_0 = 5 \times 10^{14} \text{ atoms cm}^{-3}$, the speed of sound is $c_s \sim 4.2 \text{ mm/s}$, and the healing length is $\xi \sim 120 \text{ nm}$. Length scales larger than ξ are well captured by the relativistic (phononic) approximation, a prerequisite for the analogy to relativistic QFT. Additionally, temperatures on the order of nK ensure that the zero temperature approximation is valid for short temporal couplings.

The interaction between rubidium and potassium can be controlled via Feshbach resonances employing controlled

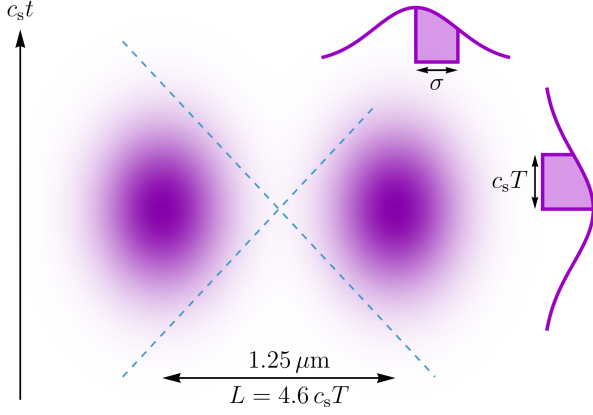


FIG. 2. Spacetime density plot of the interaction regions $\Lambda_I(\mathbf{x}) = \chi(t)f_I(\mathbf{x})$ for ^{41}K with $\sigma = 0.83c_s T$ ($\Omega = 26.3$ krad/s).

magnetic fields [88,89]. For our purposes, the zero crossing of the interaction, i.e., decoupling the detector from the field, is important and has been studied in detail [90,91]. Additionally, the interaction can be changed by small variations in the magnetic field (~ 1 G) to values as high as $a_{ab} = \pm 1000a_0$ [92,93]. The scattering properties of rubidium are not changed when varying the magnetic field in this range, ensuring that the substrate for the phononic field remains unchanged while coupling the detector.

Entanglement harvesting from a BEC—When systems interact with a quantum field they can become entangled in two ways. The first (and most commonly considered) method is when the probes exchange information, and entanglement is mediated by the field. The second method for a field to entangle two probes is when the probes extract entanglement from the field itself, effectively performing an entanglement swap operation. These methods contribute differently in different regimes: in the long time limit, entanglement extraction is negligible compared to mediated entanglement. In contrast, in the short time limit, when the probes effectively cannot communicate, they can become entangled only by extracting previously existing field entanglement (see Fig. 2). This can be used to infer the existence of entanglement between the regions that the probes couple to, corresponding to the protocol of entanglement harvesting [35,36,38].

When two bound impurities (A and B) initialized in their ground state interact with a BEC for finite times, their final state is

$$\hat{\rho}_{AB} = \begin{pmatrix} 1 - \mathcal{P}_{AA} - \mathcal{P}_{BB} & 0 & 0 & \mathcal{M}^* \\ 0 & \mathcal{P}_{BB} & \mathcal{P}_{AB} & 0 \\ 0 & \mathcal{P}_{AB}^* & \mathcal{P}_{AA} & 0 \\ \mathcal{M} & 0 & 0 & 0 \end{pmatrix} + \mathcal{O}(g_{AB}^4), \quad (12)$$

where

$$\begin{aligned} \mathcal{P}_{IJ} &= \frac{\bar{g}_{ab}^2}{\hbar^2} \int d^4\mathbf{x} d^4\mathbf{x}' \Lambda_I(\mathbf{x}) \Lambda_J(\mathbf{x}') e^{-i\Omega(t-t')} \langle \delta\hat{\rho}(\mathbf{x}) \delta\hat{\rho}(\mathbf{x}') \rangle, \\ \mathcal{M} &= -\frac{\bar{g}_{ab}^2}{\hbar^2} \int d^4\mathbf{x} d^4\mathbf{x}' \Lambda_A(\mathbf{x}) \Lambda_B(\mathbf{x}') e^{i\Omega(t+t')} \\ &\quad \times \langle T[\delta\hat{\rho}(\mathbf{x}) \delta\hat{\rho}(\mathbf{x}')] \rangle. \end{aligned} \quad (13)$$

$I, J \in \{A, B\}$, $\Lambda_I(\mathbf{x}) = \chi(t)f_I(\mathbf{x})$ are the spacetime smearing functions of the interactions, where $\chi(t)$ is determined by the time profile of the magnetic field applied to the system (Feshbach tuning) and $f_I(\mathbf{x})$ are determined by the trapping potentials. The entanglement in the final state of the probes can be quantified through the negativity [94], a faithful entanglement quantifier for two-qubit systems. Assuming identical shapes for the detectors, $\mathcal{P}_{AA} = \mathcal{P}_{BB} = \mathcal{P}$, the negativity reads

$$\mathcal{N} = \max(|\mathcal{M}| - \mathcal{P}, 0). \quad (14)$$

Overall, the negativity is a competition between a nonlocal term $\mathcal{M}$ and a local noise term $\mathcal{P}$, corresponding to the detectors' excitation probability.

We now connect theoretical concepts of entanglement harvesting to experimentally accessible parameters of our platform. Estimates use analytical predictions within Gaussian approximations [95], with switching function $\chi(t) = e^{-t^2/2T^2}$ and spatial smearing $f(\mathbf{x})$ approximated by a normal distribution of width $\sigma = \sqrt{\hbar/m_a\Omega}$, where Ω is the impurity trap frequency (see Supplemental Material [81]).

Since the negativity scales as a_{ab}^2 , maximizing the interspecies scattering length enhances the signal. The signal is also favored by a large sound speed in the rubidium background, requiring high density ρ_0 . However, three-body losses scale as $a_{ab}^4\rho_0^2$ and deplete the impurities. We therefore fix $\rho_0 = 5 \times 10^{14} \text{ atoms cm}^{-3}$ ($c_s = 4.2 \text{ mm/s}$) and $a_{ab} = -1000a_0$. For the optimal coupling times, we find that the impurity's three-particle collision loss is below 20%.

In the upper panel of Fig. 3, we show the negativity (14) for the chosen atomic species versus trap frequency Ω and interaction time T , with detector separation fixed at $L = 4.6 \times c_s T \sim 1.25 \mu\text{m}$. The signaling estimator [96] vanishes on parameters corresponding to solid lines, confirming that the entanglement is harvested from the BEC.

The maximal negativity with negligible signaling $\mathcal{N} \sim 2 \times 10^{-4}$ at $\Omega = 26.3$ krad/s requires estimating $\mathcal{P}$ and $\mathcal{M}$ with comparable precision; both are about 1 order of magnitude larger than $\mathcal{N}$ (see the lower panel in Fig. 3). While $\mathcal{P}$ follows directly from local detection of excited-state populations, estimating $\mathcal{M}$ requires additional unitary transformations, achievable by controlling the trap potentials before population detection (see Supplemental Material [81]) [76]. For the potentials detailed in the Supplemental Material [81], the corresponding infidelities of $\mathcal{P}$ and $\mathcal{M}$ measurements can be realized below 10^{-5} more than 1 order of magnitude below the expected signal.

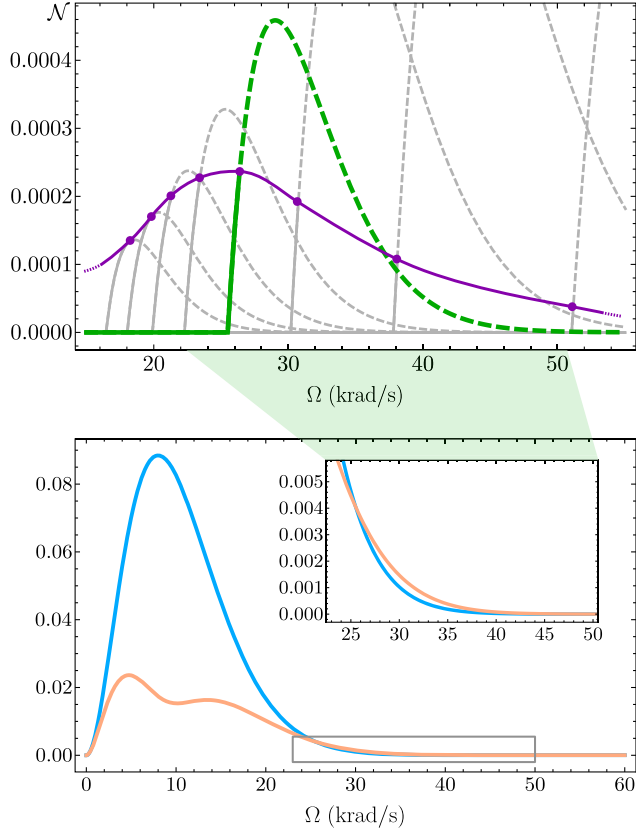


FIG. 3. Entanglement harvesting from a Rb BEC with speed of sound $c_s = 4.2$ mm/s using ^{41}K probes with $a_{ab} = -1000a_0$. The upper panel shows the negativity $\mathcal{N} = \max(|\mathcal{M}| - \mathcal{P}, 0)$ as a function of the impurity-trap frequency Ω for a distance of $L = 4.6 \times c_s T$ and interaction time $T \in \{105 \mu\text{s}, 95 \mu\text{s}, \dots, 35 \mu\text{s}\}$, with $T = 65 \mu\text{s}$ highlighted in green. Solid lines denote zero signaling estimator. The purple curve interpolates the corresponding negativity values, showing the maximal entanglement that can be harvested. $\mathcal{P}$ (blue) and $|\mathcal{M}|$ (orange) are displayed in the lower panel for $T = 65 \mu\text{s}$.

Thus, the ultimate precision is statistics limited: for $\mathcal{P} \approx 10^{-3}$, resolving the effect requires $\sim 10^5$ experimental realizations, straightforwardly obtained in modern experiments [97].

Our platform also allows the implementation of recently proposed routes to enhance harvested entanglement, i.e., increase negativity, including optimized switching functions [98] and utilizing multiple probes [99], each shown to increase the negativity by 10^2 . In particular, the optimization of [98] applies directly, since arbitrary $\chi(t)$ can be implemented by controlling the magnetic-pulse shape.

Conclusions—We show that a bound impurity coupled to a BEC via controlled interspecies interactions directly realizes an Unruh-DeWitt detector interacting with a relativistic quantum field. This brings entanglement harvesting, a paradigmatic relativistic quantum-information protocol, within experimental reach even without specific optimizations. Due to the pristine control of spatial and

temporal degrees of freedom, this BEC particle detector can probe local physics in scenarios not yet experimentally explored. This paves the way for implementations of other relevant protocols such as quantum collect calling [33] and quantum energy teleportation [31,32,100] (also see Ref. [101] for simulating protocols in Lorentz-violating QFTs).

Our proposal has three main advantages when compared to previous entanglement harvesting experimental proposals in BECs [30]. First, it makes spacetime locality explicit, allowing controllable qubits to extract and store local information in QFT. Second, it allows for the implementation of oscillatory switching functions, which can enhance entanglement harvesting by several orders of magnitude [37,98]. And third, it can be implemented in an effective $3+1$ -dimensional spacetime. Beyond controlling the probes, we can also envision pushing the relativistic field to theoretically challenging regimes, e.g., by implementing self-interactions. These experiments can give an alternative input for tackling these problems, opening the doors to phenomenological approaches to relativistic quantum information.

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Data availability—The data that support the findings of this article are openly available [102].

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