Rapid Communications

Zero-point excitation of a circularly moving detector in an atomic condensate and phonon laser dynamical instabilities

Jamir Marino,[1,2,3,*] Gabriel Menezes,[4,5,*] and Iacopo Carusotto[6]

[1]*Department of Physics, Harvard University, Cambridge, Massachusetts 02138, USA*
[2]*Department of Quantum Matter Physics, University of Geneva, 1211, Geneve, Switzerland*
[3]*Institut für Physik, Johannes Gutenberg Universität Mainz, D-55099 Mainz, Germany*
[4]*Department of Physics, University of Massachusetts Amherst, Massachusetts 01003, USA*
[5]*Departamento de Fisica, Universidade Federal Rural do Rio de Janeiro, 23897-000, Seropedica, RJ, Brazil*
[6]*INO-CNR BEC Center and Department of Physics, University of Trento, I-38123 Povo, Italy*

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We study a circularly moving impurity in an atomic condensate for the realization of superradiance phenomena in tabletop experiments. The impurity is coupled to the density fluctuations of the condensate and, in a quantum field theory language, it serves as an analog of a detector for the quantum phonon field. For sufficiently large rotation speeds, the zero-point fluctuations of the phonon field induce a sizable excitation rate of the detector even when the condensate is initially at rest in its ground state. For spatially confined condensates and harmonic detectors, such a superradiant emission of sound waves provides a dynamical instability mechanism, leading to a phonon lasing concept. Following an analogy with the theory of rotating black holes, our results suggest a promising avenue to quantum simulate basic interaction processes involving fast-moving detectors in curved space-times.

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I. INTRODUCTION

Since Unruh's pioneering proposal in 1981 [1], the last decades of research activities have witnessed a surge of a field where concepts of general relativity and quantum field theories in curved backgrounds are investigated in the so-called analog models of gravity [2]. As a most celebrated example, acoustic analogs of black holes have been studied in trans-sonically flowing atomic Bose-Einstein condensates: The acoustic black-hole horizon corresponds to the interface between regions of, respectively, sub- and supersonic flows, and was anticipated to emit a thermal radiation of phonons via Hawking processes [3]. The first experimental observations of such phenomena [4] were instrumental in triggering the ongoing explosion of the field, with a revived interest in using analog models to investigate a variety of different effects of quantum field theories in curved space-times, from the dynamical Casimir effect [5] to acceleration radiation [6] to vacuum friction and Casimir forces [7,8].

The subject of the present Rapid Communication is the phenomenon of *rotational superradiance* [9,10], namely, the amplification of classical waves reflected by a fast rotating body. In the simplest formulation, superradiance processes are efficient whenever the linear velocity of the object (or parts of it) exceeds the phase velocity of the waves, so the wave frequency seen in the comoving frame turns negative. Such negative-energy modes then provide the energy that is required to amplify positive-energy waves via superradiant effects. In cylindrical geometries, a mode of frequency ω and azimuthal quantum number n can be superradiantly amplified when the angular velocity Ω of the rotating body satisfies $\Omega > \omega/n$.

Being a consequence of basic kinematical arguments, superradiance is a ubiquitous phenomenon in physics. Its first incarnation was the discovery of amplification of acoustic waves hinging upon a supersonically moving boundary [11] or the amplification of cylindrical electromagnetic waves interacting with a rotating material [12]. In an astrophysical rotating black hole, the superradiant amplification of waves is a consequence of the spacelike character of the generator of time translations inside the ergosphere [10,13] and is at the root of the several instability phenomena of Kerr black holes [14,15]. In the framework of analog models, theoretical studies have investigated superradiant phenomena in rotating classical and quantum systems [16], offering interpretations of basic hydrodynamic phenomena [17]. Experimental evidence of superradiant scattering of classical surface waves on water was reported in Ref. [18]. Much less studied are the quantum features when superradiant processes are triggered by zero-point fluctuations and the even more intriguing quantum friction effects that result from back-reaction of superradiance on the rotational motion [19,20].

In this Rapid Communication, we investigate superradiant phenomena that can occur in ultracold atomic systems. In

[*]These authors contributed equally to this work.

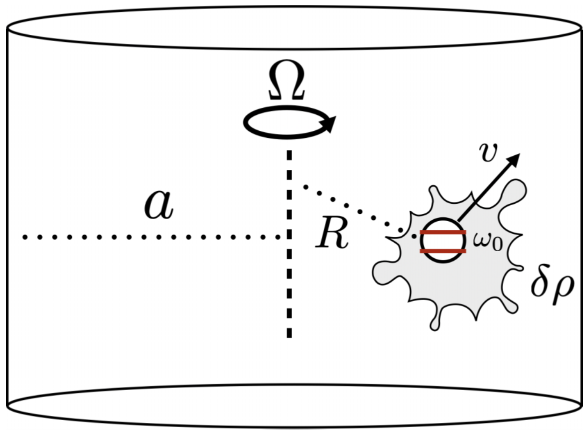

FIG. 1. A two-level detector with internal frequency, ω_0, rotates with uniform angular velocity Ω at a distance R from its rotational axis. The impurity couples to the density fluctuations, $\delta\rho(\mathbf{r})$, of a weakly interacting Bose gas, which is trapped inside a cylindrical cavity of radius a; the cavity is modeled via Dirichlet boundary conditions on the density fluctuations, $\delta\rho(\mathbf{r})|_{r=a} = 0$.

contrast to the conventional case of rotating fluids considered in superradiance, we study a configuration where the quantum fluid is at rest but a neutral impurity moves (classically) through the cold gas at sufficiently large speeds, as sketched in Fig. 1. As has been originally discussed in Refs. [6,8], the neutral impurity plays the role of a two-level detector in a canonical quantum field theory setup, and it allows us to explore superradiant phenomena beyond the usual amplification of incident waves [10,16]. Starting from the quantum vacuum state of the phonon field, we predict the spontaneous excitation of the internal degrees of freedom of the impurity in response to zero-point quantum fluctuations in the condensate, which in turn get amplified into real phonons. In a trapped geometry, the finite size of the fluid provides a perfectly reflecting cavity for the phonon modes; in this way, the self-stimulation of the detector leads to a dynamical instability for sound waves, which is rooted in superradiance. In contrast to the usual laser operation which requires an external pumping of the gain medium, the *phonon-lasing* mechanism envisioned here is driven by the mechanical motion of a detector that is initially prepared in its ground state.

II. A CIRCULARLY MOVING IMPURITY IN A UNIFORM CONDENSATE

Consider a two-level phonon detector with internal frequency ω_0 in circular motion with constant angular velocity Ω at distance R from its rotation axis (see Fig. 1 for an illustration). As originally proposed in Refs. [6,8], the detector is assumed to be coupled to density fluctuations of a weakly interacting three-dimensional Bose gas. Such a detector can be realized by means of an atomic quantum dot [21,22], namely, an impurity atom immersed in the condensate and collisionally coupled to the Bose gas via two channels. The first term is reminiscent of the interaction of a charged particle to an external scalar potential and can be canceled via proper tuning of the interaction constants (e.g., via Feshbach resonances). In this case, only the second interaction term survives, and we find the Hamiltonian

$$H = \frac{\omega_0}{2}\sigma_z + g_-\sigma_x\delta\rho(\mathbf{r}) + H_B, \tag{1}$$

where σ_x, σ_z are quantum operators (proportional to the Pauli matrices) associated with the two-level detector, while $\delta\rho(\mathbf{r})$ are the density fluctuations of the BEC which couple to the detector via the coupling constant g_-. It immediately recognizable how this Hamiltonian closely resembles the dipole coupling between a neutral polarizable object and the electromagnetic field. Extension of the theory to detectors with a harmonic oscillator internal structure is straightforward and will be discussed in the last section.

Under the standard weak interaction limit for Bose gases, density fluctuations can be treated within Bogolyubov theory [23]. The Bose gas Hamiltonian is, therefore, given by the usual expression $H_B = \int d\mathbf{q}\, \omega_\mathbf{q}\, b^\dagger_\mathbf{q} b_\mathbf{q}$, where $b^\dagger_\mathbf{q}, b_\mathbf{q}$ are the creation and annihilation operators of Bogolyubov quanta of momentum $\mathbf{q}$, whose dispersion relation is $\omega_\mathbf{q}$. As is customary, we refer with m to the mass of the condensate's particles, with μ to the condensate's chemical potential, with c to its speed of sound, and with ξ to its healing length.

III. VACUUM EXCITATION RATE

Since the motion of the detector is noninertial, we can expect a nonvanishing transition probability $A_\uparrow$ for the detector to jump from the ground to the excited state even for a condensate initially in its ground state. This effect is due to the zero-point quantum fluctuations in the phonon quantum vacuum and is associated with the emission of a phonon. Related excitation mechanisms have been discussed for a detector in uniform supersonic motion along a rectilinear trajectory in standard electromagnetism, the so-called Ginzburg effect [24], as well as in a BEC analog model of the latter [8].

While the emission from circularly moving detectors with relativistic accelerations shares analogies with the Unruh effect [25], it is important to highlight a crucial difference. For circular motion, the detector is not expected to emit radiation thermally equilibrated at the Unruh temperature [26–30], contrary to its linearly accelerating counterpart: attempts to define an effective temperature become problematic in the spatial region $r > c_l/\Omega$ (where c_l is the speed of light), since inconsistencies related to causality prevent defining unambiguously a concept of a rotating vacuum and to build excitations on top of it [31].

The transition probability $A_\uparrow$ can be calculated making use of second-order perturbation theory (see Ref. [32]). As sketched in Fig. 1, the condensate is assumed to be radially confined in a cylinder of radius a and to extend indefinitely along z. The cylindrical confinement is modeled by imposing Dirichlet boundary conditions, $\delta\rho(r)|_{r=a} = 0$, on the density perturbation (see Refs. [31,32] for further details) and implies quantization of the Bogolyubov cylindrical waves, with spatial mode profiles proportional to $J_n(\xi_{n\nu}r/a)\, e^{in\theta}\, e^{ikz}$: for each value n of the angular momentum, the radial momenta $q_{n\nu} \equiv \xi_{n\nu}/a$ are determined by the νth zero $\xi_{n\nu}$ of the Bessel function $J_n(\cdot)$. Given the infinite size of the BEC along z, the linear momentum k can have arbitrary values.

A plot of the dimensionless ground-state excitation rate, $\Gamma_\uparrow$, as a function of the rescaled detector speed $\bar{v} = \Omega a/c$, is

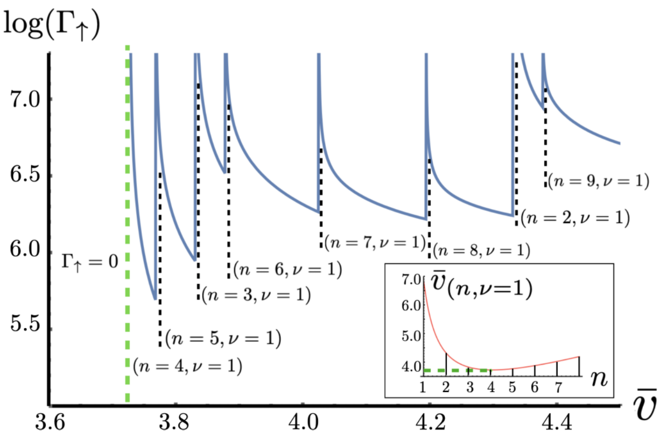

FIG. 2. The dimensionless excitation rate, $\Gamma_\uparrow = A_\uparrow\,\omega_0/P(\omega_0, R)$ [with $P(\omega_0, R) = g_-^2\rho_0/(2m\omega_0 R^5)$], evaluated for $\bar{y} = \xi/a = 0.4$, $\bar{R} = R/a = 0.6$, and $\omega_0/\Omega = 0.3$, as function of the rescaled rotation speed $\bar{v} = (\Omega a)/c$ of the detector. For these specific values of parameters, the rate drops to zero for $\bar{v} \lesssim 3.73$. The first threshold corresponds to setting $n = 4$ and $\nu = 1$ in Eq. (3) and is indicated by a dotted green line in the main panel. Each peak represents the resonant contribution associated to the excitation of one eigenmode of the cylindrical cavity. In the figure, we indicate some of the resonance channels that open upon increasing $\bar{v}$, labeled by the pair of quantum numbers (n, ν). The inset shows that the minimal velocity for the onset of spontaneous excitation of the ground state occurs at $n = 4$ for $\nu = 1$ (green dotted line), in agreement with the main panel. Higher values of ν yield resonances falling outside the range of $\bar{v}$ plotted in the main panel [for instance, $\bar{v}_{(n=1,\nu=2)} \simeq 17$ and $\bar{v}_{(n=2,\nu=2)} \simeq 9.64$].

reported in Fig. 2. For the Bogolyubov mode of azymuthal and radial quantum numbers (n, ν), such excitation rate is nonvanishing only for

$$n\Omega > q_{n\nu}c\sqrt{1 + \left(\frac{q_{n\nu}\xi}{2}\right)^2} + \omega_0, \qquad (2)$$

which provides a generalized superradiant condition on the angular velocity to excite a phonon with a given angular momentum n. This condition involves the frequency ω_0 of the detector and the cutoff frequency $\omega_{n\nu,k=0}$ of the branch of Bogolyubov modes propagating along the axis of the cylinder in the $n\nu$ radial-azymuthal mode; in our case, $\omega_{n\nu,k=0} = q_{n\nu}c[1 + (q_{n\nu}\xi/2)^2]^{1/2}$.

A further trend in the strength of the emission is due to the Bessel factor $J_n(\xi_{n\nu}R/a)$ appearing in the mode profile that suppresses the coupling of the detector to the high angular momentum phonon modes. This Bessel factor and, in particular, its strong suppression at short radii R (for $n > 0$), puts on rigorous grounds the usual qualitative reasoning based on the local dispersion relation of the waves in the rotating frame and on the necessity of a local supersonic motion [19].

In the geometry under consideration here, $\Gamma_\uparrow$ displays peaks whenever the condition Eq. (2) holds as an equality for a given $n, \nu > 0$ pair; this indicates the opening of a new emission channel occurring at

$$\bar{v}_{n,\nu} = \frac{\xi_{n\nu}\sqrt{4 + (\xi_{n\nu}y)^2}}{2|n - \omega_0/\Omega|}. \qquad (3)$$

The excitation rate can be then be rewritten as (see Ref. [32])

$$\Gamma_\uparrow \propto \sum_{n=1}^{\infty}\sum_{\nu=1}^{\infty} \gamma_{n\nu}(\bar{v})\theta(n - \omega_0/\Omega)\theta(\bar{v} - \bar{v}_{n,\nu}), \qquad (4)$$

where the square-root divergence $\gamma_{n\nu}(\bar{v}) \propto [\bar{v} - \bar{v}_{n,\nu}]^{1/2}$, visible in Fig. 2, follows from the effective one-dimensional density of states of each radial-azymuthal branch of Bogolyubov eigenmodes in the cylindrically shaped condensate. Interestingly, the dependence of $\bar{v}_{n,\nu}$ on n is nonmonotonous for a fixed value of ν as illustrated in the inset of Fig. 2: This feature explains the nonmonotonous labeling of the peaks visible in the main panel of Fig. 2.

Even though we have restricted our attention to the excitation of the detector, it is useful to recall that this process is always strictly associated with the emission of phonons propagating away from the detector along the BEC axis $\hat{z}$ (which may be detected following, for instance, Ref. [4]). In passing, we notice that stimulation of the process by an external incident field would lead to a superfluid analog of Zeldovich amplification of electromagnetic waves by a rotating dielectric [12].

As a final remark, it is worth highlighting that the emission processes studied in this Rapid Communication have a spontaneous nature. They are thus very different from synchrotron radiation emitted by circularly moving charges in classical electrodynamics [33], in exactly the same way as Ginzburg emission from superluminally moving polarizable objects [24] is conceptually different from the Cherenkov radiation emitted by moving charges or static dipoles [34]. Analogs of such Cherenkov and synchrotron processes might occur if the tuning of the Feshbach resonance mentioned in the paragraph before Eq. (1) was not perfect; nevertheless, the statistical properties of the phonon radiation emitted would be drastically different in this case and, more importantly, such processes could not lead to the dynamical instabilities which we discuss in the following, and which represent one of the salient features of our study.

IV. DYNAMICAL INSTABILITIES

In the astrophysical context, rotational superradiance can give rise to different kinds of dynamical instabilities depending on the specific geometry, from black-hole bombs to ergoregion instabilities [10]. Analog effects are also at play in rotating superfluids [17]. In this final section, we explore dynamical instability mechanisms induced by the circularly moving detector considered in this Rapid Communication.

To favor self-stimulation of the superradiant process, it is convenient to focus on a fully confined, pancake-shaped condensate with discrete Bogolyubov modes. To avoid the instability being disturbed by saturation of the two-level detector, we must extend our model by assuming that the internal structure of the detector is well approximated by a harmonic oscillator coupled to the density fluctuations of the condensate. To this purpose, one can consider a large number N of two-level atoms whose average distance is smaller than the magnitude of the inverse wave-vector undergoing the dynamical instability: analogously to the Dicke model [35], all the atoms can then be grouped into a large collective spin of size proportional to N, with the consequence that saturation effects will start to become relevant only when an equally large number of phonons is emitted. Technically, this corresponds to approximate a large spin of size N with harmonic oscillator creation and annihilation operators using a Holstein-Primakoff transformation [35] and to introduce a cutoff on the number of excitations in the detector proportional to N.

Restricting for simplicity our attention to the lowest excitation mode along the $\hat{z}$ direction, the total Hamiltonian in the frame comoving with the rotating detector is then given by

$$H = \omega_0 d^\dagger d + \sum_{n,\nu} \bar{\omega}_{n\nu} \bar{b}^\dagger_{n\nu} \bar{b}_{n\nu} + \frac{\mathcal{G}}{\mathcal{N}} x\, \delta\rho(\mathbf{r}_D), \tag{5}$$

where $d, d^\dagger$ are the harmonic oscillator destruction and creation operators for the impurity and we have defined $x = (d + d^\dagger)/\sqrt{2\omega_0}$. Furthermore, we have set $\bar{\omega}_{n\nu} = \omega_{n\nu} - n\Omega$ equal to the Bogolyubov mode frequency in the rotating frame and we indicate with r_D the radial position of the detector. In practical calculations, the sums will be restricted to $-n_0 \leqslant n \leqslant n_0$ and $0 \leqslant \nu \leqslant \nu_0$. For convenience, a factor of $\mathcal{N} \equiv (2n_0 + 1)\nu_0$ has been included to count the number of cylindrical Bogolyubov modes to ensure a proper scaling of the coupling in the multimode limit $\mathcal{N} \gg 1$.

To identify dynamical instabilities, we consider the corresponding equations of motion,

$$\begin{aligned} \dot{d} &= -i\left[\left(\frac{\omega_0}{\Omega}\right)d + \sum_{n=-n_0}^{n_0}\sum_{\nu=1}^{\nu_0} g_{n\nu}(\bar{b}_{n\nu} + \bar{b}^\dagger_{n\nu})\right], \\ \dot{\bar{b}}_{n\nu} &= -i[\tilde{\omega}_{n\nu}\bar{b}_{n\nu} + g_{n\nu}(d + d^\dagger)], \end{aligned} \tag{6}$$

where we have introduced the dimensionless coupling between modes and impurities,

$$g_{n\nu} \equiv \frac{\bar{g}}{\mathcal{N}} \frac{\xi_{n\nu}}{\bar{v}^2} \frac{J_n(\xi_{n\nu}\bar{R})}{J_{n+1}(\xi_{n\nu})} \left(\frac{\Omega}{2\omega_0(\widetilde{\omega}_{n\nu} + n)}\right)^{1/2}, \tag{7}$$

with $\bar{g} = \mathcal{G}\sqrt{\rho_0/m}/c^2$, and $\widetilde{\omega}_{n\nu} = \bar{\omega}_{n\nu}/\Omega$. Furthermore, in the Heisenberg Eqs. (6), time has been rescaled by Ω^{-1}.

The different kinds of dynamical instabilities that this Hamiltonian can display are physically understood considering the simplified two-mode bosonic model Hamiltonian:

$$H = \omega_0 a^\dagger a + \bar{\omega} b^\dagger b + g(a^\dagger b + a^\dagger b^\dagger + \text{H.c.}). \tag{8}$$

The eigenvalues of the associated set of linear Heisenberg equations of motion,

$$\lambda = \pm \frac{i}{\sqrt{2}} \sqrt{\omega_0^2 + \bar{\omega}^2 - \sqrt{\left(\omega_0^2 - \bar{\omega}^2\right)^2 + 16 g^2 \omega_0 \bar{\omega}}}, \tag{9}$$

allow for direct inspection for the conditions of stability. If ω_0 and $\bar{\omega}$ are both positive, the onset of the first type of dynamical instability occurs for $g > \sqrt{\omega_0\bar{\omega}}/2$, therefore a tiny value of g can induce unstable dynamics near $\bar{\omega} \simeq 0$. On the other hand, exactly on the parametric resonance $\bar{\omega} + \omega_0 = 0$, imaginary values of λ can be found for any value of g. In the two-dimensional many-mode problem defined by Eq. (5), this latter condition is satisfied when Eq. (2) holds as an equality—a circumstance made possible by the rotational Doppler shift experienced by the Bogolyubov modes.

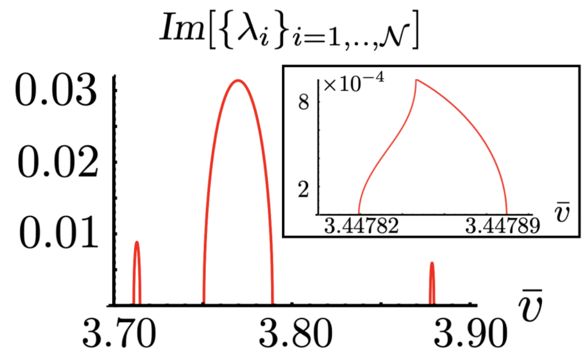

FIG. 3. Positive imaginary parts of the eigenvalues of Eq. (6) as a function of $\bar{v}$, evaluated for $\bar{y} = 0.4$, $\bar{R} = 0.6$, $\omega_0/\Omega = 0.3$, $\bar{g}/\mathcal{N} \simeq 1.53$; a numerical cutoff at $n_0 = 9$ and $\nu_0 = 1$ has been used ($\mathcal{N} = 19$). The nonvanishing imaginary parts visible in the $\bar{v}$ interval plotted in the figure correspond to the parametric resonances $\bar{\omega}_{n,\nu} + \omega_0 = 0$ for, respectively, $n = 4, 5, 6$. The parametric resonance for $n = 3$ triggers a small imaginary part of order 10^{-3} (not visible in the figure) around $\bar{v} \simeq 3.82$. The dynamical instabilities corresponding to the other peaks of Fig. 2 fall at larger $\bar{v}$ values. Other types of instabilities arise when the conditions $\bar{\omega}_{n,\nu} \simeq 0$ are satisfied; the one occurring at $n = 3$ and $\nu = 1$ is illustrated in the inset.

This physics is illustrated in Fig. 3 where the imaginary parts of the eigenvalues of the linear system Eqs. (6) for the fully multimode problem are evaluated for the same parameters employed in the plot of Fig. 2. The resonant condition underlying each instability window is specified in the caption. Different strengths are found for instabilities of the two types. For those of the second kind (akin to parametric instabilities), a crucial contribution is due to the spatial profile of the mode via the $J_n(\xi_{n\nu}\bar{R})$ factor. Physically, the onset of instabilities will be observable as an exponential growth of the amplitude of some Bogolyubov mode at a rate set by the imaginary part of the eigenvalue, in concomitance with an analogous exponential growth of the internal oscillation amplitude of the detector. Even though this instability mechanism would be quickly saturated for a single two-level impurity after the emission of the first phonon, it can lead to sizable excitations if several impurities are used to mimic a harmonic oscillator, as discussed at the beginning of this section.

As compared to the notoriously elusive nature of the superradiance effect in the electromagnetic context, we anticipate that the phonon instabilities discussed here can be employed as a way to reinforce the signature of superradiance by inspecting the quick exponential growth of the detector excitation. Further insight into the underlying process can be provided by the spectral selective and velocity-dependent nature of the amplification mechanism (recall that the unstable modes are those satisfying the resonance condition $\bar{\omega}_{n,\nu} \simeq 0$). Experimentally, information on the emitted phonons can in fact be retrieved from the density profile of the BEC and its noise properties. To give a concrete estimate on the timescales required to observe the lasing instability, one can consider the largest real positive eigenvalue $\lambda^* \simeq 0.03$ of Fig. 3, and evaluate the associated timescales in units of Ω^{-1} [cf. Eqs. (6)], which yields $t^* \simeq 1/(\lambda^*\Omega) \simeq 0.3$s, if $\Omega \simeq 100$ Hz as in typical experiments for rotating BECs [36]. Since the lifetime of cold atoms is of the order of many seconds, we expect that the lasing instability discussed in this Rapid Communication should be within reach of state-of-the-art experiments in the field.

In addition to their intrinsic interest as a novel manifestation of superradiance, dynamical instabilities triggered by moving detectors are also of great interest as a concept of phonon lasing, where the amplification mechanism is provided by the mechanical motion of a detector and not by some external pumping. Even though self-supported oscillations are a common feature in classical acoustics as well as in laser operations, nontrivial mechanisms for mechanical oscillation accompanied by the onset of quantum fluctuations are currently of high interest in a broad range of platforms, e.g., driven-dissipative coupled microcavities [37], ion-traps [38], nanomagnets [39], and optically driven quantum dots [40].

V. PERSPECTIVES

To summarize, we have shown in this Rapid Communication that a circularly moving impurity immersed in an atomic condensate at rest constitutes a promising avenue to investigate quantum features of rotational superradiance in a novel context. If the interaction of the impurity is tuned in a way to serve as a phonon detector, signatures of superradiance include the excitation of the impurity by zero-point quantum fluctuations of the phonon field in the condensate, and the onset of dynamical instabilities for the Bogolyubov modes which can serve as a new avenue for phonon lasing.

Beyond the specific configurations investigated in this Rapid Communication, our results suggest that moving impurities in condensates can be employed as a platform to investigate basic interaction processes between fields in intricate curved space-time geometries, with emitters that move at speeds comparable to the wave velocity. Different cosmological scenarios and new aspects of trans-Planckian physics with cold atoms can be addressed by tuning the microscopic properties of the BEC, e.g., introducing dipolar interactions as suggested in Ref. [41]. Another intriguing future direction consists of analyzing the impact of superradiant effects on higher order quantum vacuum processes such as the Casimir-Polder forces between a pair of circularly rotating impurities. Finally, an exciting challenge is to extend our proposal to photonic quantum simulators, in particular to identify a viable implementation of the moving detector concept in quantum fluids of light [42].

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