Letter

Impact of trans-Planckian excitations on black-hole radiation in dipolar condensates

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(Received 8 November 2022; accepted 17 May 2023; published 21 June 2023)

We consider a quasi-one-dimensional dipolar condensate in a steplike analog black hole setup. It is shown that the existence of roton excitations leaves significant imprints onto the Hawking radiation spectrum. The emitted radiation depends on the depth of the roton minimum, and is in general more intense. In addition, we find a novel spontaneous particle creation mechanism with no counterpart in nondipolar condensates. Our results establish that dipolar condensates offer a richer and more versatile environment for the simulation of particle production from the quantum vacuum in the presence of horizon interfaces than contact-interaction condensates.

DOI: 10.1103/PhysRevD.107.L121502

Black holes are a versatile laboratory to probe particle production from the quantum vacuum in the presence of horizon interfaces separating distinct regions of spacetime [1–3]. A suitably clean and controllable arena to produce analog black holes are Bose-Einstein condensates [4], in which the first unambiguous detection of spontaneous quantum Hawking radiation has been achieved [5,6]. It has been argued in the years since its inception that Hawking radiation is a rather universal phenomenon, as the thermality of the Hawking spectrum is generally robust against trans-Planckian deformations of the spectrum breaking Lorentz invariance [7–11], even though the emitted quanta at infinity, when traced back to the horizon, experience an infinite blueshift [12]. Black holes as well as their analogs are however conventionally set up in the field-theoretical context of contact interactions. In the following, we demonstrate that admitting nonlocal field theories offers a much richer arena to harness the impact of trans-Planckian excitations on black-hole radiation. Specifically, we show that due to the increased complexity of the scattering problem at the horizon, novel features emerge which clearly distinguish dipolar black holes from their contact counterparts. To this end, we take into account that interactions can be long range and in particular anisotropic, for which dipole-dipole interactions between atoms or molecules with magnetic or electric dipole moments are the archetype. Due to the roton minimum in their elementary excitation spectrum [13–16], black holes in dipolar condensates sensitively probe the robustness of Hawking radiation thermality to high frequency dispersion, thereby also thoroughly addressing one of the major original motivations of the whole analog black hole program [17].

We thus provide below the first simulation of black hole (BH) analogs from dipolar Bose-Einstein condensates, which has important ramifications in analog gravity physics and cold atoms theory, as well as more generally in quantum nonlocal field theory and in phenomenological theories of Lorentz-violating quantum gravity.

Our system is assumed to be an elongated radially harmonic trapped flowing quasi-one-dimensional (quasi-1D) dipolar condensate, systems routinely realized in experiment [18]. We assume that the system is stationary and sufficiently strongly elongated such that the details of how the flow is sustained can be neglected to a first approximation cf., e.g., Refs. [19–27].

We recall that three ingredients are necessary to trigger spontaneous particle production of quantum origin in stationary condensates: The existence of negative energy excitations, a mechanism of mode conversion [20,23,28], also cf. Ref. [29], and that the fundamental commutation relations for the field in question are fulfilled [30]. For contact BH analogs, negative energy superluminal excitations reach the event horizon from within the BH, which converts a part of these modes into outgoing spontaneous radiation, thus giving rise to the system's vacuum decay [20,23]. We show in what follows that spontaneous particle production also occurs in dipolar condensates when these three criteria are met.

For quasi-1D dipolar condensates, after integrating out the radial directions transverse to the long x axis, the ratio $\beta = \ell_\perp/\xi_{\rm u}$, where $\ell_\perp = 1/\sqrt{\omega_\perp}$ is the transverse harmonic oscillator length and $\xi_{\rm u}$ is the subsonic healing length outside the analog BH (setting $\hbar = m = 1$), measures how deep into the quasi-1D regime the system is [15,31]. There is a critical $\beta_c \simeq 0.776$ for which the system becomes unstable against a proliferation of rotonic excitations in the crossover to three spatial dimensions. In the limit $\beta \to 0$, indicating a quasi-1D condensate with effective contact interactions, our model system coincides with the one studied in [22], see for further details below. At finite β, for which the anisotropy of the dipolar interaction becomes manifest, we find important differences between contact and dipolar BHs that can be summarized as follows: (i) Increasing β from zero in general leads to larger Hawking radiation power than in the contact case. (ii) The rotonic/maxonic dispersion relation typical of dipolar gases leads to a strong nonthermality of the Hawking radiation spectrum. (iii) We find a clear correlation between the BH temperature and the appearance of a second spontaneous mode conversion mechanism, in addition to the usual Hawking mode mechanism, which has no analog in contact BHs. These findings demonstrate that dipolar condensates offer a much richer environment to simulate the Hawking phenomenon in comparison to contact condensates.

Within the mean-field approximation, the evolution of the condensate order parameter ϕ is described by the nonlocal Gross-Pitaevskii equation (GPE) [32]

$$i\partial_t\phi = \left[-\frac{1}{2}\partial_x^2 + U + g_{\rm dd}|\phi|^2\right]\phi - 3g_{\rm dd}\phi G * |\phi|^2, \quad (1)$$

where $U = U(x)$ is the trap potential, $g_{\rm dd} > 0$ is the quasi-1D dipolar interaction strength, and $G * f$ denotes the convolution $\int dx' G(x - x')f(x')$. We assume that the dipoles have a common direction relative to the long x axis, which fixes $g_{\rm dd}$ [31,33]. Note that we have separated off the contact part of the dipolar interaction (third term in square brackets), such that $\int dx G(x) = 0$, and we assume other contact interaction contributions coming from s-wave scattering to be negligible. This regime can be achieved by using Feshbach resonances [18]. For computational convenience, we discretize the quasi-1D dipolar kernel G in Eq. (1) [32,33]. We have

$$G(x) = \delta(x) - \frac{1}{2\ell_\perp}\frac{\sum_{j=1}^{\mathcal{N}} j^2\Delta q^3 e^{-j^2\Delta q^2/2 - j\Delta q|x|/\ell_\perp}}{\sum_{j=0}^{\mathcal{N}} j\Delta q^2 e^{-j^2\Delta q^2/2}}, \quad (2)$$

where the limit $\mathcal{N} \to \infty$, $\Delta q \to 0$ is understood, leading to the continuum expression $G(x) - \delta(x) = -(1/2\ell_\perp)f(|x|/\ell_\perp)$, $f(y) = -y + (1 + y^2)\exp(y^2/2)\times \mathrm{Erfc}(y/\sqrt{2})\sqrt{\pi/2}$, Erfc being the complementary error function [31]. This form of writing the quasi-1D dipolar kernel enables the construction of essentially analytical solutions: We solve for finite $\mathcal{N}$ and Δq and take the limits afterward. Also, by keeping $\mathcal{N}$ and Δq finite we can produce approximate solutions to the problem to the desired accuracy. For instance, as shown in [33], for $\mathcal{N} = 10$ and $\Delta q = 1/3.4$ the error is below 1%, maintained throughout our simulations.

To convey the essential physics, we construct our stationary dipolar BH analog as the solution $\phi = \sqrt{\rho}\exp(-i\mu t + ivx)$ $(v > 0)$ of Eq. (1) for a piecewise constant density: $\rho = \rho_{\rm u}$ for $x < 0$ and $\rho = \rho_{\rm d} < \rho_{\rm u}$ if $x > 0$. A subscript on local quantities "u" here and in what follows denotes the upstream region ($x < 0$), and "d" the downstream region ($x > 0$) (Fig. 1 top). This model was recently studied for contact condensates in [22]. We assume zero temperature throughout. Equation (1) then implies continuity $\rho_{\rm u}v_{\rm u} = \rho_{\rm d}v_{\rm d}$ and $U = \mu + \frac{\partial_x^2\sqrt{\rho}}{2\sqrt{\rho}} - \frac{1}{2}v^2 - g_{\rm dd}\rho + 3g_{\rm dd}G * \rho$, which fixes the external potential U (cf. Fig. 1 top).

We scale lengths in units of $\xi_{\rm u}$, e.g., $x = x[\xi_{\rm u}]$, so that the energy and (inverse time unit) becomes $1/\xi_{\rm u}^2$. A dipolar BH is then specified by the set of parameters $\{\mathfrak{m}_{\rm u}, \mathfrak{m}_{\rm d}, \beta = \beta_{\rm u} = \ell_\perp\}$, where the local Mach number is $\mathfrak{m} = v/c$

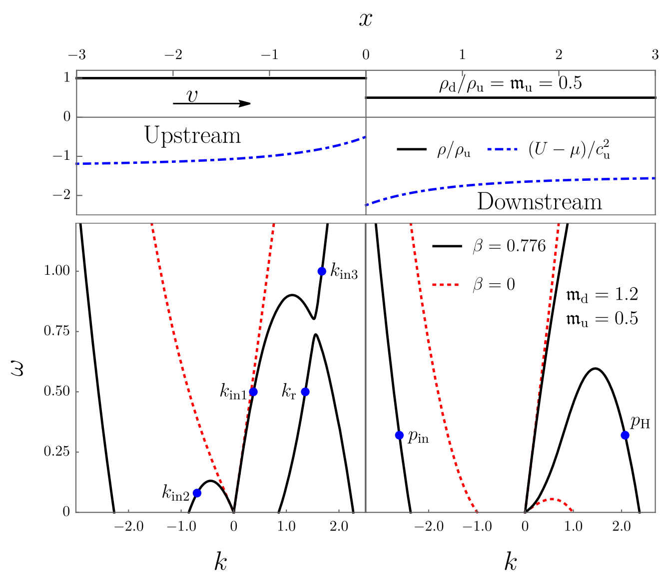

FIG. 1. Top: density jump model imposed by the external potential U. Bottom: Bogoliubov dispersion relation for two BH analogs with distinct β and fixed Mach numbers. Each line of constant ω intercepts the dispersion relation at the real wave vector solutions, corresponding to plane waves propagating rightwards (leftwards) if the slope at the interception point is positive (negative). Left: dispersion relation for the upstream region. Each plane wave propagating to the right gives rise to a distinct quasiparticle mode, indexed by $k_{\rm in1}$, $k_{\rm in2}$, $k_{\rm in3}$, and $k_{\rm r}$. Note that the (dispersive) "rotonic" branch $k_{\rm r}$ has negative energy. We also note that, importantly, whenever $k_{\rm r}$ exists, the upstream region also contains an outgoing negative energy channel. Right: downstream dispersion relation. Each leftwards propagating wave gives rise to a quasiparticle, indexed by $p_{\rm in}$ and the negative energy $p_{\rm H}$ (Hawking branch), which is also associated to an outgoing negative energy channel. Dipolar interactions, in particular, increase the cutoff frequency of the local maximum on the right.

and $c = \sqrt{g_{\rm dd}\rho}$ is the local sound speed. Then $\mathfrak{m}_{\rm u} < 1 < \mathfrak{m}_{\rm d}$ defines an analog BH. Our goal is to determine how small fluctuations over this BH background lead to spontaneous radiation. The Bogoliubov expansion of the wave function reads $\hat{\Psi} = (\sqrt{\rho} + \hat{\psi})\exp(-i\mu t + ivx)$, where the bosonic operator $\hat{\psi}$ models the small quantum fluctuations.

In order to study the Hawking radiation in our system, we expand the quantum field $\hat{\psi}$ in a basis of quasiparticle modes whose vacuum state represents a zero flux of phonons sent towards the event horizon in the laboratory frame [20]. Hence for this vacuum choice any spontaneous radiation is linked to Hawking-like processes. Our representation for the interaction kernel (2) can be used to find such a basis as follows [32]. We first define the Nambu spinor in particle-hole space, $\hat{\Phi} = \binom{\hat{\psi}}{\hat{\psi}^\dagger}/\sqrt{\rho}$, and expand $\hat{\Phi} = \sum_n (\hat{a}_n\Phi_n + \hat{a}_n^\dagger\sigma_1\Phi_n^*)$, where σ_i, $i = 1, 2, 3$ are Pauli matrices and $\{\Phi_n\}$ is a complete set of positive norm quasiparticle solutions with respect to the Bogoliubov scalar product: $\int dx\rho\Phi_n^\dagger\sigma_3\Phi_{n'} = \delta_{n,n'}$. The functions Φ_n are solutions of the Bogoliubov–de Gennes equation

$$i\partial_t\sigma_3\Phi_n = -\frac{1}{2\rho}\partial_x(\rho\partial_x\Phi_n) - iv\sigma_3\partial_x\Phi_n + \rho g_{\rm dd}\sigma_4\Phi_n - 3g_{\rm dd}\sigma_4 G * \Phi_n, \quad (3)$$

where we defined $\sigma_4 = 1 + \sigma_1$. We find that $\sigma_1\Phi_n^*$ is a negative norm solution. The quantum number n will be identified below with the quasiparticle frequency ω.

Solutions to Eq. (3) can be determined assuming the time dependence $\exp(-i\omega t)$, with $\omega > 0$. Furthermore, because Eq. (2) is suppressed for $|x| \gg 1, \ell_\perp$, all solutions of Eq. (3) far from the horizon are written in terms of the local homogeneous condensate perturbations, i.e., we have superposition of plane waves in the form $\Phi(t, x) = \exp(-i\omega t + ikx)\Phi_k$ for constant Φ_k. This gives rise to the dispersion relation $\omega \coloneqq \omega^\pm(k)$, where

$$\omega^\pm = \mathfrak{m}_{\rm u}\frac{\rho_{\rm u}}{\rho}k \pm k\sqrt{(\rho/\rho_{\rm u})[1 - 3\tilde{G}(\beta k)] + k^2/4}, \quad (4)$$

and $\tilde{G}(\beta k) = \tilde{G}(\ell_\perp k) = \int {\rm d}x G(x)\exp(-i\ell_\perp kx/\ell_\perp)$.

Equation (4) enables us to identify the required quasiparticle basis whose vacuum state is by definition characterized by no quasiparticles propagating towards the event horizon. This state is, therefore, suitable for studying the spontaneous Hawking radiation in BH analogs. A positive (negative) sign of the dimensionless group velocity $v_g(k) = {\rm d}\omega/{\rm d}k$ (sign of the slopes in Fig. 1 bottom) determines whether the plane wave is propagating to the right (left), and each plane wave propagating towards the event horizon gives rise to a distinct quasiparticle mode found by solving the scattering problem within Bogoliubov theory [20,23,33,34]. We note that in this work we reserve the terms quasiparticle and channel to denote solutions of the Bogoliubov–de Gennes equation and the plane waves given by Eq. (4), respectively.

Let us recollect the salient features of the dispersion relation in the contact BH regime (dotted curves of Fig. 1, bottom). When $\beta = 0$ there is only one plane wave going towards the horizon from the upstream region (Fig. 1, bottom left) for each ω, whereas in the downstream region (Fig. 1, bottom right) for frequencies below the local maximum (Hawking cutoff frequency) there are always two dispersive channels of opposite energy sign propagating towards the horizon. When $\beta > 0$, the observed effect inside the black hole is the increase of the local maximum (Hawking cutoff) frequency. We denote by $p_{\rm in}$, $p_{\rm H}$ the wave vectors of the incoming channels, with the negative energy quasiparticles indexed by $p_{\rm H}$ from the "Hawking branch" leading to spontaneous radiation process. Furthermore, novel phenomena are expected to occur outside the black hole. As β continuously increases from zero, initially no qualitative distinction from the contact case occurs (dotted curves in Fig. 1, bottom left). However, when the roton minimum emerges (continuous curves in Fig. 1, bottom left), we find that four plane waves $k_{\rm in1}$, $k_{\rm in2}$, $k_{\rm in3}$, $k_{\rm r}$ approach the horizon, and the rotonic branch $k_{\rm r}$ has negative energy. The latter is present whenever $\beta > \beta_{\rm r} \simeq 0.63$ for $\mathfrak{m}_{\rm u} = 0.5$, and it is indicative of a strong departure from a contact-dominated BH analog. The increased number of mode conversion mechanisms in comparison to contact BHs shows that the scattering problem in our dipolar case is significantly more intricate than for the contact BH analogs.

We show in the Supplemental Material [32] that the normalized quasiparticles for $\omega > 0$ can be written as

$$\Phi_\omega^{(\alpha)} = e^{-i\omega t}\begin{cases} \sum_p S_p^{(\alpha)} e^{ipx}\Phi_p, & x > 0 \\ \sum_k S_k^{(\alpha)} e^{ikx}\Phi_k, & x < 0 \end{cases}, \quad (5)$$

where we denote the downstream wave vectors by p. Here, α can be any of the incoming wave vectors displayed in Fig. 1, namely, $k_{\rm in1}, k_{\rm in2}, k_{\rm in3}, k_{\rm r}, p_{\rm in}, p_{\rm H}$. For each given α, the sums in Eq. (5) include the incoming channel with $S_{\alpha'}^{(\alpha)} = \delta_{\alpha,\alpha'}$, all outgoing propagating channels, and all evanescent waves. Both $S_k^{(\alpha)}$ and $S_p^{(\alpha)}$ are fixed by Eq. (3), and the sign of the norm

$$\Phi_k^\dagger\sigma_3\Phi_k = \frac{1}{2\pi\rho|v_g(k)|}{\rm sgn}(\omega - \mathfrak{m}_{\rm u}k\rho_{\rm u}/\rho) \quad (6)$$

determines whether an incoming channel propagating towards the horizon has positive or negative energy for real k [32]. For each given frequency, we let the set $\Gamma_\omega^{(+)}$ (respectively, $\Gamma_\omega^{(-)}$) contain the positive (respectively negative) energy incoming channels. Accordingly, the full field operator expansion reads

$$\hat{\Phi} = \int_0^\infty \mathrm{d}\omega \Bigg[\sum_{\alpha\in\Gamma_\omega^{(+)}} (\hat{a}_\omega^{(\alpha)}\Phi_\omega^{(\alpha)} + \hat{a}_\omega^{(\alpha)\dagger}\sigma_1\Phi_\omega^{(\alpha)*}) + \sum_{\alpha\in\Gamma_\omega^{(-)}} (\hat{a}_\omega^{(\alpha)\dagger}\Phi_\omega^{(\alpha)} + \hat{a}_\omega^{(\alpha)}\sigma_1\Phi_\omega^{(\alpha)*}) \Bigg]. \quad (7)$$

Furthermore, $[\hat{a}_\omega^{(\alpha)}, \hat{a}_{\omega'}^{(\alpha')\dagger}] = \delta_{\alpha,\alpha'}\delta(\omega - \omega')$, and the vacuum state $|0\rangle$ is defined by $\hat{a}_\omega^{(\alpha)}|0\rangle = 0$.

The operator $\hat{\psi}$ is the upper component of $\sqrt{\rho}\hat{\Phi}$, and once the quantum field expansion is obtained, we find that the (normal ordered) system Hamiltonian assumes the diagonal form [32]

$$\hat{H} = \int_0^\infty \mathrm{d}\omega\omega \Bigg[\sum_{\alpha\in\Gamma_\omega^{(+)}} \hat{a}_\omega^{(\alpha)\dagger}\hat{a}_\omega^{(\alpha)} - \sum_{\alpha\in\Gamma_\omega^{(-)}} \hat{a}_\omega^{(\alpha)\dagger}\hat{a}_\omega^{(\alpha)} \Bigg], \quad (8)$$

similar to the contact BH [20]. The Hamilton operator above demonstrates that exciting a quasiparticle mode with index in $\Gamma_\omega^{(-)}$ diminishes the system energy. Note that only when $\beta = 0$ and thus $G = 0$ (contact-only case) energy is locally conserved in the system, whereas for any finite β there is no local energy conservation [32]. Nevertheless, we find generally that energy is globally conserved: $\partial_t H = -S_\infty + S_{-\infty} = 0$, where H is the system total energy. Using the unitarity of the scattering process, we show in the Supplemental Material [32] that the power radiated at $x \to -\infty$ can be expressed by the scattering coefficients of the negative energy quasiparticles alone:

$$S_{-\infty} = \frac{1}{2\pi}\int_0^\infty \mathrm{d}\omega\omega\mathcal{F}_\omega,$$
$$\mathcal{F}_\omega = \sum_{\alpha\in\Gamma_\omega^{(-)}} \sum_{k\,\mathrm{real}} |S_k^{(\alpha)}|^2 \mathrm{sgn}[v_g(k)(\omega - \mathfrak{m}_\mathrm{u} k)]. \quad (9)$$

The sum in k here is performed over all upstream propagating waves. The quantity $\mathcal{F}_\omega$ is the power spectrum, containing the negative energy quasiparticle modes from the expansion in Eq. (7). Note that for the dipolar gas, the set $\Gamma_\omega^{(-)}$ contains both the rotonic branch k_r (when $\beta > \beta_\mathrm{r}$) and the Hawking branch p_H, which is the only member of that set for contact interactions (also see Fig. 1).

We recall that for thermal radiation at temperature T the power spectrum reads $n_\omega = (e^{\omega/T} - 1)^{-1}$, and in particular $\omega n_\omega \to T$ when $\omega \to 0$. It is instructive to assign a temperature for $\mathcal{F}_\omega$ as $T = \lim_{\omega\to 0}\omega\mathcal{F}_\omega$, such that we can assess the thermality of the spectrum through graybody factors σ_ω defined by $\sigma_\omega = \mathcal{F}_\omega(e^{\omega/T} - 1)$. $\sigma_\omega = 1$ then corresponds to a pure thermal spectrum. We depict T in Fig. 2 top left panel as function of β for $\mathfrak{m}_\mathrm{u} = 0.5$ and several values of $\mathfrak{m}_\mathrm{d}$. The figure shows that the observed T can be higher or smaller in comparison to contact-only BH ($\beta = 0$) due to the dipolar interactions. In particular, a strong increase in T is observed near the point $\beta = \beta_\mathrm{r}$ where the rotonic branch emerges, representing a novel temperature signature of the latter. Figure 2 top right depicts the graybody factors for $\mathfrak{m}_\mathrm{u} = 0.5$, $\mathfrak{m}_\mathrm{d} = 1.2$, and several values of β. BHs within the range $0 < \beta \leq 0.4$, where there are no rotons and maxons present, possess similar graybody factors (and spectrum), with the difference that the cutoff frequency is higher, corresponding to a higher radiated power. However, we see from the figure that, when rotons are present, i.e., for $\beta \geq 0.52$ when $\mathfrak{m}_\mathrm{u} = 0.5$, a strong departure from thermality is observed. For instance, when $\beta = 0.55$, $\sigma_\omega \sim 30$ for some frequencies (dotted line in Fig. 2 top right). In order to understand the origin of this departure, we plot in Fig. 2 (bottom) the spectrum $\mathcal{F}_\omega$ and the dispersion relation, from which we see a clear correlation between the formation of the roton minimum (right panel) and an increase in the power spectrum (left panel) when $\beta = 0.55$, $\mathfrak{m}_\mathrm{u} = 0.5$.

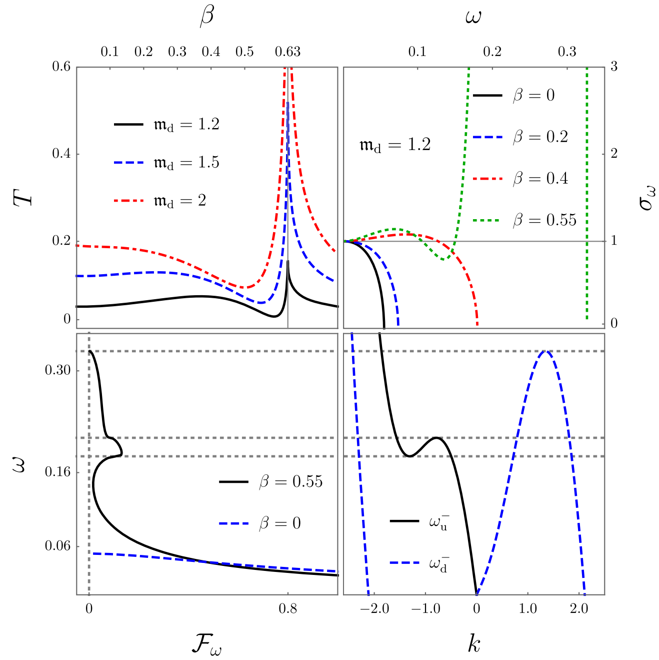

FIG. 2. Top left: radiation temperature of dipolar BH analogs as function of β for $\mathfrak{m}_\mathrm{u} = 0.5$ and several values of $\mathfrak{m}_\mathrm{d}$. The dipolar interactions can increase or decrease the temperature, and a strong increase is observed near the roton branch formation threshold $\beta = \beta_c = 0.63$. Top right: graybody factors as function of frequency for $\mathfrak{m}_\mathrm{u} = 0.5$, $\mathfrak{m}_\mathrm{d} = 1.2$, and several values of β. As β increases from zero, the graybody factors are similar to the ones found for local BH analogs. However, after the roton minimum formation occurs at $\beta \sim 0.52$, a strong departure from thermality is observed (dotted curve). In the bottom part, we also display the correlation between the spectral distribution for $\beta = 0.55$ (left) and the dispersion relation (right), showing the existing correlation between the increase in the spectral function and the appearance of the roton minimum.

As β approaches the deepest possible roton minimum at $\beta_c = 0.776$, the system enters the vicinity of a dynamical instability and the departure from thermality is so strong that the rotonic branch can, in principle, even suppress the radiation power or revert its direction, which represents a dramatic impact of trans-Planckian physics. Finally, a curious feature of our flowing condensate model is that the negative energy rotonic quasiparticles might be present even without a horizon ($\mathfrak{m}_{\rm d} < 1$), and can be scattered at the interface at $x = 0$ owing to nonthermal outgoing radiation.

In conclusion, we considered here for the first time the possibility of simulating BHs in dipolar Bose-Einstein condensates. We have shown that the presence of dipolar interactions leads to a marked departure from contact condensates, including a novel mechanism of mode conversion of the rotonic branch of quasiparticles at the interface, leading to radiating scenarios even when there is no horizon. The presence of a roton minimum leads to strong nonthermality of the spectrum as in the analog Unruh effect [35]. Furthermore, dipolar condensates feature enhanced radiation power, making them especially promising candidates to probe Hawking radiation.

Generally, we expect in dipolar gases significant alterations of the predictions of a Lorentz-invariant theory, also and in particular in cosmological scenarios that employ a similar dispersion relation see, e.g., Refs. [36–39]. Finally, our analysis focuses on the existence of quasiparticle radiation mechanisms which characterize a dipolar BH analog model. Future studies will investigate the effects of “zero modes” on the very dipolar condensate existence [25,40,41], and the radiation and pair-entanglement verification procedure via density-density correlations [27,42–45].

This work has been supported by the National Research Foundation of Korea under Grants No. 2017R1A2A2A05001422 and No. 2020R1A2C2008103.

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