Published in final edited form as:

Nat Phys. 2018 May ; 14(5): 442–446. doi:10.1038/s41567-018-0054-7.

Observation of Roton Mode Population in a Dipolar Quantum Gas

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Abstract

The concept of a roton, a special kind of elementary excitation, forming a minimum of energy at finite momentum, has been essential to understand the properties of superfluid ⁴He ¹. In quantum liquids, rotons arise from the strong interparticle interactions, whose microscopic description remains debated ². In the realm of highly-controllable quantum gases, a roton mode has been predicted to emerge due to magnetic dipole-dipole interactions despite of their weakly-interacting character ³. This prospect has raised considerable interest ^{4–12}; yet roton modes in dipolar quantum gases have remained elusive to observations. Here we report experimental and theoretical studies of the momentum distribution in Bose-Einstein condensates of highly-magnetic erbium atoms, revealing the existence of the long-sought roton mode. Following an interaction quench, the roton mode manifests itself with the appearance of symmetric peaks at well-defined finite momentum. The roton momentum follows the predicted geometrical scaling with the inverse of the confinement length along the magnetisation axis. From the growth of the roton population, we probe the roton softening of the excitation spectrum in time and extract the corresponding imaginary roton gap. Our results provide a further step in the quest towards supersolidity in dipolar quantum gases ¹³.

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Competing Interests: The authors declare that they have no competing financial interests.

Author contributions: FF, LC, DP, GF, MJM, JHB, SB conceived and supervised the experiment and took the experimental data. LC analysed them. RMWvB developed the BdG calculations. FW, RMWvB and LS performed the real time simulations. LS derived the analytical and SSM. LC, FF, RMWvB and LS wrote the paper with the contributions of all the authors.

Data availability: The data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.

Quantum properties of matter continuously challenge our intuition, especially when many-body effects emerge at a macroscopic scale. In this regard, the phenomenon of superfluidity is a paradigmatic case, which continues to reveal fascinating facets since its discovery in the late 1930s 2, 12. A major breakthrough in understanding superfluidity thrived on the concept of *quasiparticles*, introduced by Landau in 1941 1. Quasiparticles are elementary excitations of momentum k , whose energies ε define the dispersion (energy-momentum) relation $\varepsilon(k)$.

To explain the special thermodynamic properties of superfluid ^4He , Landau postulated the existence of two types of low-energy quasiparticles: phonons, referring to low- k acoustic waves, and rotons, gapped excitations at finite k initially interpreted as elementary vortices. The dispersion relation continuously evolves from linear at low k (phonons) to parabolic-like with a minimum (roton) at a finite $k = k_{\text{rot}}$. Neutron scattering experiments confirmed Landau's remarkable intuition 14. In liquid ^4He , k_{rot} scales as the inverse of the interatomic distance. This manifests a tendency of the system to establish a local order, which is driven by the strong correlations among the atoms 2.

In the realm of low-temperature quantum physics, ultra-cold quantum gases realise the other extreme limit for which the interparticle interactions - and correlations - are typically weak, meaning that classically their range of action is much smaller than the mean interparticle distance. Because of this diluteness, roton excitations are absent in ordinary quantum gases, i. e. in Bose-Einstein condensates (BECs) with contact (short-range) interactions 12.

However, about 15 years ago, seminal theoretical works predicted the existence of a roton minimum both in BECs with magnetic dipole-dipole interactions (DDIs) 3 and in BECs irradiated by off-resonant laser light 15. Following the lines of the latter proposal, a roton softening has been recently observed in BECs coupled to an optical cavity 16. Here, the excitation arises from the infinite-range photon-mediated interactions and the inverse of the laser wavelength sets the value of k_{rot} . Additionally, roton-like softening has been created in spin-orbit-coupled BECs 17, 18 and quantum gases in shaken optical lattices 19 by engineering the single-particle dispersion relation.

Our work focuses on dipolar BECs (dBECs). As in superfluid ^4He , the roton spectrum in such systems is a genuine consequence of the underlying interactions among the particles. However, in contrast to helium, the emergence of a minimum at finite momentum does not require strong inter-particle interactions. It instead exists in the weakly-interacting regime and originates from the peculiar anisotropic and long-range character of the DDI in real and momentum space 3–12. Despite the maturity achieved in the theoretical understanding, the observation of dipolar roton modes has remained so far an elusive goal. For a long time, the only dBEC available in experiments consisted of chromium atoms 20, for which the achievable dipolar character is hardly sufficient to support a roton mode. With the advent of the more magnetic lanthanide atoms 21, 22, a broader range of dipolar parameters became available, opening the way to access the regime of dominant DDI. In this regime, novel exciting many-body phenomena have been recently observed, as the formation of droplet states stabilised by quantum fluctuations 23–25, which may become self-bound 26. Lanthanide dBECs hence open new roads toward the long-sought observation of roton modes.

Prior to this work, dipolar rotons have been mostly connected to pancake-like geometries 3–6, 8–12. Here, we extend the study of roton physics to the case of a cigar-like geometry with trap elongation along only one direction (y) transverse to the magnetisation axis (z) (Fig. 1a). The anisotropic character of the DDI (Fig. 1b, inset) together with the tighter confinement along z is responsible for the rotonization of the excitation spectrum along y (Fig. 1b). To illustrate this phenomenon, we consider an infinite cigar-shaped dBEC and focus on its axial elementary excitations, of momentum k_y . These excitations correspond in real space to a density modulation along y of wavelength $2\pi/k_y$. For low k_y , the atoms sit mainly side-by-side and the repulsive nature of the DDI prevails, stiffening the phononic part of the dispersion relation (Fig. 1b1). In contrast, for $k_y \ell_z \gtrsim 1$, ℓ_z being the characteristic z confinement length, the excitation favours head-to-tail alignments and the DDI contribution to $\epsilon(k_y)$ eventually changes sign 3 (Fig. 1b2). The resulting softening of $\epsilon(k_y)$ is counterbalanced by the contributions of the repulsive contact interaction, and of the kinetic energy, which ultimately dominates at very large k_y . For strong enough DDI, this competition gives rise to a roton minimum in $\epsilon(k_y)$, occurring at a momentum $k_y = k_{\text{rot}}$ set by the geometrical scaling $k_{\text{rot}} \sim 1/\ell_z$ (see below and e.g. 3, 8, 9).

Similar to the helium case, the roton energy gap, $\epsilon(k_{\text{rot}})$, depends on the density and on the strength of the interactions. In ultracold gases, both quantities can be controlled. In particular, the scattering length a_s , setting the strength of the contact interaction, can be tuned using Feshbach resonances 12. As a_s is reduced, $\epsilon(k_{\text{rot}})$ decreases, vanishes, and eventually becomes imaginary (Fig. 1b). In the latter case, the system undergoes a roton instability and the population at $k_y = 0$ is transferred in $\pm k_{\text{rot}}$ at an exponential rate 5, 6. The population of the roton mode is then readily visible in the momentum distribution of the gas (Fig. 1c-d). In the extensively-studied pancake geometries, the roton population in k -space spreads over a ring of radius $k = k_{\text{rot}}$ because of the radial symmetry of the confinement (Fig. 1d). Such a spread can be avoided using a cigar geometry. Here, the roton population focuses in two prominent peaks at $k_y = \pm k_{\text{rot}}$, enhancing the visibility of the effect (Fig. 1c).

We explore the above-described physics using strongly magnetic ^{166}Er atoms. The experiment starts with a stable dBEC in a cigar-shaped harmonic trap of frequencies $\nu_{x,y,z}$ elongated along the y axis. The trap aspect ratio, $\lambda = \nu_z/\nu_y$, can be tuned from about 4 to 30, corresponding to ν_z ranging from 150 Hz to 800 Hz, whereas ν_y and ν_z/ν_x are kept constant at about 35 Hz and 1.6, respectively (Methods). An external homogeneous magnetic field, B , fixes the dipole orientation (magnetisation) with respect to the trap axes and sets the values of a_s through a magnetic Feshbach resonance, centered close to $B = 0\text{G}$. In previous experiments, we precisely calibrated the B -to- a_s conversion for this resonance 25. The BEC is prepared at $a_s^i = 61 a_0$ ($B = 0.4\text{G}$) with transverse (z) magnetisation. The characteristic dipolar length, defined as $a_{\text{dd}} = \mu_0 \mu_m^2 m / 12\pi \hbar^2$, is $65.5 a_0$, with m the mass and μ_m the magnetic moment of the atoms, $\hbar = h/2\pi$ is the reduced Planck constant, μ_0 the vacuum permeability and a_0 the Bohr radius.

To excite the roton mode, we quench a_s to a desired lower value, a_s^f , and shortly hold the atoms in the trap for a time t_h . We measure that a_s converges to its set value with a characteristic time constant of 1 ms during t_h . We then release the atoms from the trap,

change a_s back close to its initial value and let the cloud expand for 30 ms. We probe the momentum distribution $\tilde{n}(k_x, k_y)$ by performing standard resonant absorption imaging on the expanded cloud (Methods). The measurement is then repeated at various values of $a_s < a_{dd}$ in a fixed trap geometry. The momentum distribution shows a striking behaviour (Fig. 2). For large enough a_s , $\tilde{n}(k_x, k_y)$ shows a single narrow peak with an inverted aspect ratio compared to the trapped gas, typical of a stable BEC 12 (Fig. 2a). We define the center of the distribution as the origin of k . In contrast, when further decreasing a_s below a critical value a_s^* , we observe a sudden appearance of two symmetric finite-momentum peaks, of similar shape and located at $k_y = \pm k_{\text{rot}}$ (Fig. 2b-c). By repeating the experiment several times, we observe that the peaks consistently appear at the same locations, and they are visible in the averaged distributions. To quantitatively investigate the peak structures, we fit a sum of three Gaussian distributions to the central cuts of the average $\tilde{n}(k_x, k_y)$ (Fig. 2a1-c1). From the fit we extract the central momentum, k_{rot} , and the amplitude, A_* , of the side peaks.

A major fingerprint of the roton mode in dBECs is its geometrical nature, leading to a universal scaling $k_{\text{rot}} \sim 1/\ell_z$ (see e.g. 3, 4, 8, 9). In addition, the dependency of k_{rot} on a_s close to the instability is expected to be mild as k_{rot} remains mainly set by its geometrical nature 3, 8. We investigate both dependencies in the experiment. In a first set of experiments, we repeat the quench measurements for various trap parameters and extract k_{rot} . We clearly observe the expected geometrical scaling. k_{rot} shows a marked increase with $1/\ell_z = 2\pi\sqrt{mv_z^2/\hbar}$, matching well with a linear progression with a slope of 1.61(4) (Fig. 2d). Note that no dependence of k_{rot} on v_y is observed. In a second set of experiments, we fix the trap geometry and explore the dependence of k_{rot} on a_s . We observe that, within our experimental uncertainty, k_{rot} stays constant when decreasing a_s (Fig. 2e). This behaviour contrasts the one expected for a phonon-driven modulation instability that exhibits a strong a_s -dependence 27.

To gain a deeper understanding of the roton excitations in our system and its dynamical population, we develop both an analytical model and full numerical simulations and compare the findings with our experimental data. Our analytical model starts by calculating the roton spectrum of the stationary BEC, generalizing the results of Ref. 3 to a non-radially symmetric configuration. Since the roton wavelength is much smaller than the extension of our 3D BEC along y , this mode can be evaluated using a local density approximation in y . Hence, for our model, we consider a dBEC homogeneous along y , of axial density n_0 , harmonically confined along x and z . To analytically evaluate the roton spectrum, we approximate the BEC wavefunction using the Thomas-Fermi (TF) approximation. For dominant DDI, $\epsilon_{dd} = a_{dd}/a_s \gg 1$, we find that $\epsilon(k_y)$ indeed rotonizes (Methods). In the vicinity of the roton minimum and for $\epsilon_{dd} \sim 1$, the computed dispersion acquires a gapped quadratic form similar to that of helium rotons:

$$\epsilon(k_y)^2 \simeq \Delta^2 + \frac{2\hbar^2 k_{\text{rot}}^2}{m} \frac{\hbar^2}{2m} (k_y - k_{\text{rot}})^2. \quad (1)$$

The roton momentum reads as $k_{\text{rot}} = \sqrt{2m}(E_0^2 - \Delta^2)^{1/4}/\hbar$, with

$\Delta = \sqrt{E_0^2 - E_I^2}$, $E_I = 2gn_0(\epsilon_{dd} - 1)/3$, $g = 4\pi\hbar^2 a_s/m$, and $E_0^2 = 2g\epsilon_{dd}n_0\frac{\hbar^2}{2m}(X^{-2} + Z^{-2})$. The TF radii X, Z satisfy $X^2, Z^2 \propto gn_0$ so that $E_0 \propto \hbar v_z$ with a scale set by ϵ_{dd} and v_z/v_x but independent of gn_0 15. Close to the instability ($\Delta \simeq 0$), the roton momentum thus follows a simple geometrical scaling, $k_{\text{rot}} = \kappa/\ell_z$ with the geometrical factor κ depending on v_z/v_x alone. For completeness, we have also performed full 3D numerical calculations of the static Bogoliubov spectrum of a finite trapped dBEC, confirming the existence and scaling of the roton mode (Supplementary Information).

The above stationary description accounts for the existence of the roton mode in the cigar geometry used in experiments, and predicts the scaling of k_{rot} and Δ with the system parameters. However the quench of a_s introduces dynamics, which is crucial to quantitatively reproduce the experimental observations. The reduction of a_s decreases the contact interaction and additionally induces a compression of the cloud. This yields a dynamical modification of the local roton dispersion relation, and the roton may be destabilised during the evolution. This dynamical destabilisation is well accounted for by a self-similar model (SSM) describing the evolution of the cloud shape after the quench 12. In particular, we consider a 3D harmonic confinement, and, starting from the stationary TF profile at a_s^i , we evaluate the evolution of the cloud along the change of a_s assuming that the TF shape is maintained but with time-dependent TF radii (Methods). We then estimate the local (along y) instantaneous roton spectrum $\epsilon(k_y, y, t_h)$ using a local density approximation, i.e. evaluating Eq. (1) with the experimentally calibrated $a_s(t_h)$, and the $n_0(y, t_h)$, $X(y, t_h)$, $Z(y, t_h)$ estimated from the 3D profile (Methods). We indeed find that the roton gap decreases and eventually turns imaginary ($\Delta(t_h) < 0$) (Fig. 3a). When this occurs the population at $k_y \approx k_{\text{rot}}$ exponentially grows with an instantaneous local rate $2\text{Im}[\epsilon(k_y, y, t_h)]/\hbar$, $\text{Im}[\cdot]$ indicating the imaginary part, giving rise to two symmetric side peaks in the axial momentum distribution $\tilde{n}_1(k_y, t_h) \propto \int \exp\left(2\int_0^{t_h} \text{Im}[\epsilon(k_y, y, t)]dt\right) dy$ (Fig. 3b) (Methods). The center of the peaks, $\pm k_{\text{rot}}(t_h)$, also evolves with t_h but quickly converges after a few ms to its final value, k_{rot} (Fig. 3c). The measured k_{rot} and the calculated values from our parameter-free theory are in remarkable agreement (Fig. 2d-e and 3c).

Our SSM quantitatively explains the experimental observations and provides us with a physical understanding. For completeness, we perform numerical simulations (NS) of the system dynamics. We calculate the time evolution of the generalized non-local Gross-Pitaevskii equation, which accounts also for quantum fluctuations, three-body loss processes, and finite temperature, not included in the SSM (Supplementary Information). The first two contributions limit the peak density of the atomic cloud and stabilise the dBEC against collapse 4–6, 24, 25, whereas the latter term thermally seeds the initial roton-mode population. The NS confirm both the SSM results and the experimental observations. Indeed the calculations show that, a few ms after the quench, the system develops roton peaks in momentum space (Fig. 3d) and short-wavelength density modulations at the center of the BEC (Fig. 3e), showing the predicted roton confinement 9. The extracted value of k_{rot} and its geometrical scaling are in very good quantitative agreement with both the experimental

data and the SSM calculations (Fig. 2d-e and 3c). Interestingly, the timescales for the emergence of roton peaks in the NS are a few ms longer than those observed in the experiment. The origin of this time shift remains an open question, whose answer could e.g. require a refinement of current models of quantum fluctuations. However, despite this delay, the growth rate of the roton population in the NS matches both the experimental observations and the SSM predictions in the early dynamics (Fig. 3b).

In the experiment, we study the time evolution of the roton mode in a fixed geometry (v_z , λ). In a first set of measurements, we fix a_s^f and follow the dynamics by recording the momentum distribution at various t_h . We observe that k_{Tot} does not change significantly while the roton population initially grows, in excellent agreement with the theories (Fig. 3b-c). The growth rate is a particularly relevant quantity as it is directly connected to the imaginary excitation energy in the Bogoliubov description (for the roton, its gap), as revealed by our theory. In a second set of experiments, we thus systematically study the growth rate of the roton population for various a_s^f (Fig. 4). Our data show that the roton mode begins to become populated from a critical value of the scattering length, a_s^* . For $a_s^f > 52a_0$, we do not observe the roton peaks at any time. For $a_s^f \leq 52a_0$, after a time delay, A^* undergoes an abrupt increase. At longer time, A^* then saturates and eventually slowly decreases while atoms are coincidentally lost (Fig. 4a). By further lowering a_s^f , the roton population exhibits a faster growth rate and shorter time delay.

From the growth rate of the roton population we now extract an overall roton gap $\bar{\Delta}(a_s)$ (Methods). In brief, the dynamical variation of the gap is determined by the leading time dependence of a_s and $A^*(a_s^f, t_h) \propto \exp\left(2 \int_0^{t_h} dt \text{Im}[\bar{\Delta}(a_s(t))]/\hbar\right)$. To investigate the scaling $\bar{\Delta}(a_s)$, we first consider the analytical expression of the elementary local roton gap, $\Delta(a_s)$ as a function of a_s (see Eq. (1)). Developing $\Delta(a_s)$ in the vicinity of $a_s = a_s^*$ yields $\Delta(a_s) \propto (a_s^* - a_s)^{1/2}$. Here we confirm this scaling by considering the generic power-law dependence $\text{Im}[\bar{\Delta}(a_s)]/\hbar = \Gamma \delta(a_s)$ with $\delta(a_s) = \left(\frac{a_s^* - a_s}{a_0}\right)^\beta$ ($a_s \leq a_s^*$). Consequently, A^* grows as $\exp\left(2 \Gamma \int_0^{t_h} \delta(a_s(t)) dt\right) = \exp(2 \Gamma T_h)$. By rescaling the time variable as $t_h \rightarrow T_h = \int_0^{t_h} \delta(a_s(t)) dt$, we determine the parameters a_s^* and β such that all the curves of Fig. 4a fall on top of each other (Fig. 4b). The best overlap of the experimental curves is found for $a_s^* = 53.0(4) a_0$ and $\beta = 0.55(8)$. We then perform an exponential fit to all the rescaled data for $T_h < 3$ ms and extract $\Gamma = 465(83) \text{ s}^{-1}$. The same analysis, applied to the calculated population growth from the SSM for different a_s^f , gives $a_s^* \simeq 52.8a_0$, $\beta \simeq 0.55$, and $\Gamma \simeq 472 \text{ s}^{-1}$, which are very close to the experimental values. This time-resolved study allows to readily extract the imaginary roton gap, $\text{Im}[\bar{\Delta}]$, from both the experiment and the SSM, as a function of a_s and t_h (Fig. 4c). Our observations show the softening of the roton mode at $a_s = a_s^*$ and the expected increase of $\text{Im}[\bar{\Delta}]$ for $a_s < a_s^*$, nicely grasping the dynamics during the growth of

the roton population (inset of Fig. 4c). The observed behaviour shows again a remarkable agreement with the theory predictions.

To conclude, our work demonstrates the power of weakly-interacting dipolar quantum gases to access the regime of large-momentum, yet low-energy, excitations dressed by interactions. This newly accessible regime, which is largely unexplored in ultracold gases, raises fundamental questions and opens novel directions. Future key developments are to study the impact of such low-lying excitations on the superfluid behavior of a dBEC 8–10, 12, the interplay between phononic and rotonic modes in the stability diagram of the quantum gas 4, 7, and the possible role of roton excitations as triggering mechanism of the recently observed instability leading to the formation of metastable droplet arrays in elongated traps 24, 28. Of particular interest is the prospect of creating a supersolid and striped ground-states in dBECs 13, 28. Indeed, the short-wavelength density modulation tied to the roton softening together with the quantum stabilization mechanism can favor the formation of such an exotic phase of matter, in which crystalline order coexists with phase coherence. Contrary to recent experiments 29, 30, where the density modulation is imposed by external fields yielding a supersolid-like arrangement of infinite stiffness, a dipolar supersolid would be compressible. Experiments on dBECs provide hence the exciting opportunity to unveil similarities and differences among complementary approaches to supersolidity.

Methods

Trapping geometries

The BEC is confined in a harmonic trapping potential $V(\mathbf{r}) = 2m\pi^2(v_x^2x^2 + v_y^2y^2 + v_z^2z^2)$, characterized by the frequencies (v_x, v_y, v_z) . The ODT results from the crossing of two red-detuned laser beams of 1064 nm wavelength at their respective foci. One beam, called vODTb, propagates nearly collinear to the z -axis and the other, denoted hODTb propagates along the y -axis. By adjusting independently the parameters of the vODTb and hODTb, we can widely and dynamically control the geometry of the trap (see Supplementary Information). In particular, v_y is essentially set by the vODTb power while v_z (and v_x) independently by that of the hODTb. This yields an easy tuning of the trap aspect ratio, $\lambda = v_z/v_y$, relevant for our cigar-like geometry.

After reaching condensation (see Supplementary Information), we modify the beam parameters to shape the trap into an axially elongated configuration, favourable for observing the roton physics ($v_y \ll v_x, v_z$). The trapping geometries probed in the experiments, whose (λ, v_x, v_y, v_z) are reported in Supplementary Table I, are achieved by changing the hODTb power with the vODTb power set to 7W so that v_y and v_z/v_x are kept roughly constant. Only the green triangle in Fig. 2d is obtained in a distinct configuration, with a vODTb power of 2W, leading to $(v_x, v_y, v_z) = (156, 17, 198)$ Hz and $\lambda = 11.6$. The (λ, v_x, v_y, v_z) are experimentally calibrated via exciting and probing the center-of-mass oscillation of thermal samples. We note that the final atom number N , BEC fraction f , and temperature T after the shaping procedure depend on the final configurations, as detailed in Supplementary Table I.

Quench of the scattering length a_s

To control a_s we use a magnetic Feshbach resonance between ^{166}Er atoms in their absolute ground state, which is centered around $B = 0\text{ G}$ 31. The B -to- a_s conversion has been previously precisely measured via lattice spectroscopy, as reported in Ref. 25. Errors on a_s , taking into account statistical uncertainties of the conversion and effects of magnetic field fluctuations (e.g. from stray fields), are of 3-to-5 a_0 for the relevant a_s range 27-67 a_0 in this work. After the BEC preparation and in order to investigate the roton physics via an interaction quench, we suddenly change the magnetic field set value, B_{set} , twice. First we perform the quench itself and abruptly change B_{set} from 0.4 G ($a_s^i = 61 a_0$) to the desired lower value (corresponding to a_s^f) at the beginning of the hold in trap ($t_h = 0$ ms). Second we prepare the ToF expansion and imaging conditions (see Method Imaging procedure) and abruptly change B_{set} from the quenched value back to 0.3 G ($a_s = 57 a_0$) at the beginning of the ToF expansion ($t_f = 0$ ms). Due to delays in the experimental setup, e.g. coming from eddy currents in our main chamber, the actual B value felt by the atoms responds to a change of B_{set} via $B(t) = B_{\text{set}}(t) + \tau dB/dt$ 32. By performing pulsed-radio-frequency spectroscopy measurements (pulse duration 100 μs) on a BEC after changing B_{set} (from 0.4 to 0.2 G), we verify this law and extract a time constant $\tau = 0.98(5)$ ms. Hence a_s is also evolving during t_h and t_f on a similar timescale. This effect is fully accounted for in the experiments and simulations, and we report the roton properties as a function of the effective value of a_s at t_h . We use t_h ranging from 2 to 7 ms. The lower bound on t_h comes from the time needed for a_s to effectively reach the regime of interest. We then consider the initial evolution for which t_h/v_y ; $t_h/\tau_{\text{coll}} \ll 1$, $1/\tau_{\text{coll}}$ being the characteristic collision rate. We estimate that τ_{coll} ranges typically from 40 to 90 ms in the initial BECs of Supplementary Table I at $a_s^i = 61 a_0$. Experimentally we observe that the roton, if it ever develops, has developed within the considered range of t_h .

Imaging procedure

In our experiments, we employ ToF expansion measurements, accessing the momentum distribution of the gas 12, to probe the roton mode population. We let the gas expand freely for $t_f = 30$ ms, which translates the spatial imaging resolution ($\sim 3.7 \mu\text{m}$) into a momentum resolution of $\sim 0.32 \mu\text{m}^{-1}$ while our typical roton momentum verifies $k_{\text{rot}} \gtrsim 2 \mu\text{m}^{-1}$. After ToF expansion, we record 2D absorption pictures of the cloud by means of standard resonant absorption imaging on the atomic transition at 401 nm. The imaging beam propagates nearly vertically, with a remaining angle of $\sim 15^\circ$ compared to the z -axis within the xz -plane. Thus the ToF images essentially probe the spatial density distribution $n_{\text{tof}}(x, y, t_f)$ in the xy plane. When releasing the cloud ($t_f = 0$ ms), we change B back to $B = 0.3$ G (Method Quench of the scattering length a_s). This change enables constant and optimal imaging conditions with a fixed probing procedure, i. e., a maximal absorption cross-section. In addition, the associated increase of a_s to 57 a_0 allows to minimize the time during which the evolution happens in the small- a_s regime where the roton physics develops, such that we effectively probe only the short-time evolution in this respect. In this work, we use the simple mapping:

$$\tilde{n}(k_x, k_y) = \left(\frac{\hbar t_f}{m}\right)^2 n_{\text{tof}}\left(\frac{\hbar k_x t_f}{m}, \frac{\hbar k_y t_f}{m}, t_f\right), \quad (2)$$

which neglects the initial size of the cloud in trap and the effect of interparticle interactions during the free expansion. Using real-time simulations (see Supplementary Information), we simulate the experimental sequence and are able to compute both the real momentum distribution from the in-trap wavefunction and the spatial ToF distribution 30 ms after switching off the trap. Using the mapping of Eq. (2) and our experimental parameters, the two calculated distributions are very similar, and, in particular, the two extracted momenta associated with the roton signal agree within 5%. This confirms that the interparticle interactions play little role during the expansion and justifies the use of Eq. (2).

Fit procedure for the ToF images

For each data point of Figs. 2-4, we record between 12 and 25 ToF images. By fitting a two-dimensional Gaussian distribution to the individual images, we extract their origin $(k_x, k_y) = (0, 0)$ and recenter each image. From the recentered images, we compute the averaged $\tilde{n}(k_x, k_y)$, from which we characterise the linear roton developing along k_y . To do so, we extract a one-dimensional profile $\tilde{n}_1(k_y)$ by averaging $\tilde{n}(k_x, k_y)$ over k_x within

$$|k_x| \leq k_m = 3.5 \mu\text{m}^{-1} : \tilde{n}_1(k_y) = \int_{-k_m}^{k_m} \tilde{n}(k_x, k_y) dk_x / \int_{-k_m}^{k_m} dk_x.$$

To quantitatively analyse the observed roton peaks, we fit a sum of three Gaussian distributions to $\tilde{n}_1(k_y)$. One Gaussian accounts for the central peak and its centre is imposed to $k_0 \sim 0$. Its amplitude is denoted A_0 . The two other Gaussians are symmetrically located at $K_0 \pm k_y^*$, and we constrain their sizes and amplitudes to be identical, respectively equal to σ^* and A^* . We focus on the roton side peaks by constraining $k_y^* > 0.5 \mu\text{m}^{-1}$ and $\sigma^* < 3 \mu\text{m}^{-1}$ (peak at finite momentum of moderate extension compared to the overall distribution).

In order to analyse the onset and evolution of the roton population (see Fig.4), we perform a second run of the fitting procedure, in which we constrain the value of k_y^* more strictly. The interval of allowed values is defined for each trapping geometry based on the results of the first run of the fitting procedure. We use the results of the (a_s^f, t_h) -configurations where the peaks are clearly visible and we set the allowed k_y^* -range to that covered by the 95% confidence intervals of the first-fitted k_y^* in these configurations. This constraint enables that, for $a_s > a_s^*$, the fitting procedure estimates the residual background population on the relevant momentum range for the roton physics (see e.g. Fig.4a).

Time-rescaling analysis and roton gap estimate

In Fig. 4a-b, we systematically analyse the time evolution of the roton population for various a_s^f and link it to the roton spectrum in a quench picture. The roton population is embodied by the amplitude A^* of the three-Gaussian fit (Method Fit procedure for the ToF images), which

measures the density $\tilde{n}_1(k_{\text{rot}})$. A_* is observed to initially increase if $a_s^f < a_s^*$, its growth rate increases for decreasing a_s^f and a_s^* depends on the trap geometry.

If dynamically unstable, the k_{rot} component is expected to grow exponentially with an instantaneous rate $\text{Im}[\varepsilon(k_{\text{rot}}, t_h)]/\hbar$. We hence expect an initial growth of A_* of the form:

$$A_*(a_s^f, t_h) \propto \exp\left(2 \int_0^{t_h} dt \text{Im}\left[\bar{\Delta}(a_s^f, t)/\hbar\right]\right), \quad (3)$$

$\bar{\Delta}(a_s^f, t_h)$ being the instantaneous value of the overall roton gap, corresponding to an average over the cloud, after quenching a_s to a_s^f . Our results show that the most relevant effect of the quench on the roton spectrum is given by the reduction of a_s itself. Hence the time dependence of $\bar{\Delta}$ is determined by $a_s(t)$ and, by monitoring $A_*(a_s^f, t_h)$, one can readily extract the scattering-length dependent gap $\bar{\Delta}(a_s)$, $\bar{\Delta}(a_s^f, t) = \bar{\Delta}(a_s(t))$. To investigate the scaling $\bar{\Delta}(a_s)$, we first consider the simpler case of the static and local roton gap $\Delta(a_s)$. From Eq. (1) of the main text characterising this elementary gap, assuming n_0 , X and Z fixed, and using the fact that $\Delta(a_s = a_s^*) = 0$, one easily obtain that for a_s in the vicinity of

a_s^* , $\Delta(a_s) \propto (a_s^* - a_s)^{1/2}$. Here we verify this scaling for the global and dynamical quantity $\bar{\Delta}(a_s)$, considering the generic power-law dependence $\text{Im}[\bar{\Delta}(a_s)]/\hbar = \Gamma \delta(a_s)$ with $\delta(a_s) = \left(\frac{a_s - a_s^*}{a_0}\right)^\beta$ ($a_s(t) \leq a_s^*$), in which the parameters Γ , a_s^* and β mainly depend on the trap geometry.

Our full set of data $A_*(a_s^f, t_h)$ for the geometry $(\nu_z, \lambda) = (149 \text{ Hz}, 4.3)$ enables us to assess this scaling. Indeed Eq. (3) then reads:

$$A_*(a_s^f, t_h) \propto \exp\left(2 \Gamma \int_0^{t_h} \delta(a_s(t)) dt\right). \quad (4)$$

This defines a time rescaling $t_h \rightarrow T_h = \int_0^{t_h} \delta(a_s(t)) dt$ along which all our experimental data $A_*(a_s^f, t_h)$ should collapse in a unique curve, marked by an initial exponential growth of rate 2Γ . The relevant values of a_s^* , β and Γ are the one that result in the best overlap of the data for the initial growth of A_* (minimal spread in T_h).

In order to determine a_s^* and β , we then plot our full dataset as a function of T_h for various trial values of these parameters and evaluate the relevance of the trial couple (a_s^*, β) .

Precisely, we assess the dispersion in T_h of the full dataset for few fixed $A_* = A_*^i$. We use a

panel of 10 values A_*^i within $[80,200]\mu\text{m}^2$, corresponding to the range of the initial growth of A_* for the geometry of Fig. 4a-b. We interpolate the experimental data $A_*(T_h)$ for each a_s^f using piecewise cubic polynomial interpolation, extract the corresponding set of T_h^i at which $A_*(T_h^i) = A_*^i$. For each i , we evaluate the spread of the $T_h^i(a_s^f)$ by two complementary quantities: (a) the square-root of their variance and (b) the discrepancy between the maximal and minimal value of the set. We finally estimate the accuracy of a unified dependence $A_*(T_h)$ for a given (a_s^*, β) via the geometrical average of (a) and (b) for all the ten A_*^i values. We fit the relevant a_s^* and β by minimizing this averaged spread. For the experimental data of Fig. 4a-b, this results in $a_s^* = 53.0(4)a_0$ and $\beta = 0.55(8)$.

Using these values of a_s^* and β , we observe that the initial growth of A_* extends for $T_h \sim 3$ ms (before saturating and decreasing for longer T_h). We estimate Γ by performing an exponential fit on the full set of data $A_*(T_h)$ with $T_h \sim 3$ ms. This gives $\Gamma = 465(83) \text{ s}^{-1}$. Note that Fig. 4a-b only show 5 values of a_s^f while the reported analysis considers all available experimental data, i.e. 11 values between $31 a_0$ and $52 a_0$.

From the formula with the estimated a_s^* , β and Γ , we can then compute the global roton gap $\bar{\epsilon}$. The extracted value is only meaningful for $\bar{\epsilon}^2 < 0$. For the experiments, we also restrict its relevance to $T_h \sim 3$ ms so that, e.g., in the inset of Fig. 4c, $\text{Im}[\bar{\epsilon}]$ is shown up to $t_h = 4$ ms after which the roton population is observed to deviate from the exponential growth picture (see Figs. 3b and 4a). Note that the $\text{Im}[\bar{\epsilon}]$ estimates at the different t_h are not independent.

In Fig. 2, we report on the roton momentum as a function of the system characteristics (a_s , ℓ). Here we estimate k_{rot} from a t_h value that is individually optimised for each a_s and ℓ investigated (largest visibility). We point out that the selected t_h correspond, in the A_* dynamics, to the late stage of the exponential growth, close to the maximum (i.e. $T_h \sim 3$ ms). Here, the atom loss remains at a few percents level.

In Fig. 2d, we show the value of k_{rot} at the onset of the population of the roton mode, i.e. at $a_s = a_s^*$. We estimate a_s^* for each trap geometry by analysing the evolution of the roton population with a_s . We employ a simplified approach with respect to the one in Fig. 4 (see Supplementary Information), which we estimate to lead to a maximum underestimate for a_s^* of about $1.5 a_0$, lying within our experimental uncertainty on a_s ; see Supplementary Table II.

Analytical dispersion relation for an infinite axially elongated geometry

Equation (1) in the main text results from a similar procedure as that used in Ref. 3 for rotons in infinite pancake traps. We consider a dBEC homogeneous along y but harmonically confined with frequencies ν_x and ν_z along x and z . For sufficiently strong interactions the BEC is in the TF regime on the xz plane, in which the BEC wavefunction acquires the form $\psi_0(\rho) = \sqrt{n(\rho)}$, with $n(\rho) = n_0 (1 - (x/X)^2 - (z/Z)^2)$, where X and Z are the TF radii, and $\rho = (x, z)$. The calculation of n_0 , X and Z is detailed at the end of this section.

Due to the axial homogeneity, the elementary excitations of the Bogoliubov-de Gennes spectrum have a defined axial momentum k_y , and take the form $\delta\psi(r, t) = u(\rho)e^{ik_y y - iet/\hbar} - v(\rho)e^{-ik_y y + iet/\hbar}$, where u, v denote the amplitudes of the spatial modes oscillating in time with characteristic frequency $\epsilon/\hbar$ (see Supplementary Information for more details). We consider the GPE without LHY correction and three-body losses, and insert the perturbed solution $\psi(r, t) = (\psi_0(\rho) + \eta\delta\psi(r, t))e^{-i\mu t/\hbar}$ where μ is the chemical potential associated to ψ_0 and $\eta \ll 1$. After linearization we obtain the BdG equations for $f_{\pm}(\rho) = u(\rho) \pm v(\rho)$:

$$\epsilon f_{-}(\rho) = H_{\text{kin}} f_{+}(\rho), \quad (5)$$

$$\epsilon f_{+}(\rho) = H_{\text{kin}} f_{-}(\rho) + H_{\text{int}}[f_{-}(\rho)], \quad (6)$$

where

$$H_{\text{kin}} f_{\pm}(\rho) = \frac{\hbar^2}{2m} \left(-\nabla^2 + k_y^2 + \frac{\nabla^2 \psi_0}{\psi_0} \right) f_{\pm}(\rho), \quad (7)$$

$$H_{\text{int}}[f_{-}(\rho)] = 2 \int d^3 r' U(r - r') e^{-ik_y(y - y')} \psi_0(\rho) \psi_0(\rho') f_{-}(\rho'). \quad (8)$$

where $U(\mathbf{r}) = g \left(\delta(\mathbf{r}) + \frac{3\epsilon_{\text{dd}}}{4\pi} \frac{1 - 3\cos^2\theta}{|\mathbf{r}|^3} \right)$ is the binary interaction potential for two particles

separated by $\mathbf{r}$ including both contact and dipolar interactions and $g = \frac{4\pi\hbar^2 a_s}{m}$ (see Supplementary Information). θ is the angle between $\mathbf{r}$ and the magnetisation axis (z).

Employing $f_{+}(\rho) = W(\rho) \psi_0(\rho)$, and for $k_y \gg 1/X, 1/Z$, we obtain the following equation for the function $W(\rho)$:

$$\begin{aligned}
0 = & 2gn_0(1 - \tilde{x}^2 - \tilde{z}^2) \left[\frac{1}{X^2} \frac{\partial^2 W}{\partial \tilde{x}^2} + \frac{1}{Z^2} \frac{\partial^2 W}{\partial \tilde{z}^2} \right] \\
& - gn_0(1 + 2\varepsilon_{dd}) \left[\frac{1}{X^2} \tilde{x} \frac{\partial W}{\partial \tilde{x}} + \frac{1}{Z^2} \tilde{z} \frac{\partial W}{\partial \tilde{z}} \right] \\
& + \left(\frac{2m}{\hbar^2} (\varepsilon^2 - E(k_y)^2) - 2g\varepsilon_{dd}n_0 \left(\frac{1}{X^2} + \frac{1}{Z^2} \right) \right) \\
& - \frac{4m}{\hbar^2} gn_0(1 - \varepsilon_{dd}) E(k_y) (1 - \tilde{x}^2 - \tilde{z}^2),
\end{aligned} \tag{9}$$

where $\tilde{x} = x/X$, $\tilde{z} = z/Z$, $E(k_y) = \hbar^2 k_y^2 / 2m$. For $\varepsilon_{dd} = 1$ the last term of Eq. (9) vanishes. In that case, the lowest-energy solution is given by $W = 1$, whose eigen-energy builds, as a function of k_y , the dispersion $\varepsilon_0(k_y)$ with

$$\varepsilon_0(k_y)^2 = E(k_y)^2 + E_0^2, \tag{10}$$

with $E_0^2 = 2g\varepsilon_{dd}n_0 \frac{\hbar^2}{2m} \left(\frac{1}{X^2} + \frac{1}{Z^2} \right)$. In the vicinity of $\varepsilon_{dd} = 1$, the effect of the last term in Eq. (9) may be evaluated perturbatively, resulting in the dispersion

$$\varepsilon(k_y)^2 \simeq \varepsilon_0(k_y)^2 - 2E_I E(k_y), \tag{11}$$

with $E_I = \frac{2}{3}gn_0(\varepsilon_{dd} - 1)$.

This expression for the dispersion presents a roton minimum for $\varepsilon_{dd} > 1$ at $k_{\text{rot}} = \frac{1}{\hbar} \sqrt{2mE_I}$.

Expanding Eq. (11) in the vicinity of the roton minimum,

$$\varepsilon(k_y)^2 \simeq \varepsilon(k_{\text{rot}})^2 + \frac{1}{2} \left[\frac{d^2 \varepsilon^2(k_y)}{dk_y^2} \right]_{k_y = k_{\text{rot}}}, \text{ we obtain Eq. (1) of the main text, with}$$

$$\Delta = \varepsilon(k_{\text{rot}}) = \sqrt{E_0^2 - E_I^2}. \text{ At the instability, } \Delta = 0, \text{ and } k_{\text{rot}} = \frac{1}{\hbar} \sqrt{2mE_0}.$$

Employing a similar procedure as in Ref. 15 we obtain that the BEC aspect ratio $\chi = Z/X$ fulfills:

$$\chi^2 \left[\frac{(1 - \varepsilon_{dd})(1 + \chi)^2 + 3\varepsilon_{dd}}{(1 + 2\varepsilon_{dd})(1 + \chi)^2 - 3\varepsilon_{dd}\chi^2} \right] = \lambda_{\perp}^2, \tag{12}$$

with $\lambda_{\perp} = v_x/v_z$ and

$$Z^2 = \frac{gn_0}{2\pi^2 m v_z^2} \left[(1 + 2\varepsilon_{dd}) - \frac{3\varepsilon_{dd}\chi^2}{(1 + \chi)^2} \right] \quad (13)$$

These two equations fully determine the TF solution for given ε_{dd} , gn_0 , and $\lambda_\perp$. By inserting the expressions of X^2 and Z^2 in E_0 , we find for $\varepsilon_{dd} \simeq 1$:

$$E_0^2 = \frac{\hbar^2 v_z^2}{6} (1 + \chi)^2 \left(\lambda_\perp^2 + \frac{1}{1 + 2\chi} \right) \quad (14)$$

whereas χ simplifies into $\chi = \lambda_\perp (1 + \sqrt{1 + 1/\lambda_\perp})$. As a result, at the instability $k_{\text{rot}} \ell_z$ depends only on the transverse confinement aspect ratio $\lambda_\perp$, giving the geometrical factor κ :

$$\kappa = k_{\text{rot}} \ell_z = \left(\frac{2}{3} \right)^{1/4} \sqrt{1 + \chi} \left(\lambda_\perp^2 + \frac{1}{1 + 2\chi} \right)^{1/4}. \quad (15)$$

Self similar model and instantaneous roton spectrum

The dynamics of the condensate during and after the quench is crucial to understand the resulting momentum peaks and their direct relation to the growth of roton excitations. Due to its large characteristic momentum, the roton spectrum may be evaluated at any time using local density approximation. On the other side, the evolution of the local density is determined by the global dynamics of the condensate induced by the quench. Hence, interestingly, the analysis of the effect of the quench on the roton spectrum may be performed by combining a self-similar theory of the global dynamics 33, describing the evolution at low k_y , and the model of the roton spectrum developed in the Method Analytical dispersion relation for an infinite axially elongated geometry, accounting for the high- k_y region of interest. We note that this analysis approximates the real time evolution provided by the standard NLGPE 12, 34, 35. This treatment is valid for moderate-enough quenches so that each of the descriptions applies (see below for an in-depth discussion) and as long as the high- k modes population only minimally impact the BEC dynamics, that is for short-enough time scales.

Here, we assume that the condensate preserves its TF shape during the evolution:

$$n(\mathbf{r}, t) = n_0(t) \left[1 - \left(\frac{x}{X(t)} \right)^2 - \left(\frac{y}{Y(t)} \right)^2 - \left(\frac{z}{Z(t)} \right)^2 \right], \quad (16)$$

where $X(t)$, $Y(t)$ and $Z(t)$ are the re-scaled TF radii. Their evolution after the quench can be deduce from solving the hydrodynamics equations (see Supplementary information). Prior to the quench of a_s , the stationary TF solution is obtained from the self-consistent equations deduced from the hydrodynamics equations by cancelling all time derivative. Solving these

equations and using normalisation provides $X(0)$, $Y(0)$, $Z(0)$, and $n(0)$ for known atom number N and trap frequencies $\nu_{x,y,z}$. We use this stationary solution as the initial condition at the start of the quench ($t = 0$), and solve the system of differential equations resulting from the hydrodynamics evolution to obtain the re-scaled radii. An example for the relevant parameters of Fig. 3 is shown in Supplementary Fig. 3.

For a given position y we then evaluate the local TF profile: $n(\mathbf{r}, t) = n_0(y, t) [1 - (x/X(y, t))^2 - (z/Z(y, t))^2]$, where $n_0(y, t) = n_0(t) [1 - (y/Y(t))^2]$, $X(y, t) = X(t) [1 - (y/Y(t))^2]^{1/2}$, and $Z(y, t) = Z(t) [1 - (y/Y(t))^2]^{1/2}$, with $n_0(t) = n_0(0)X(0)Y(0)Z(0)/X(t)Y(t)Z(t)$. We may then employ a local density approximation and evaluate the local (in y) instantaneous roton spectrum using the results of the Method Analytical dispersion relation for an infinite axially elongated geometry:

$$\begin{aligned} \frac{\varepsilon^2(k_y, y, t)}{(h\nu_z)^2 t_z^4} &\simeq \frac{k_y^4}{4} - \frac{8\pi}{3} n_0(y, t) (a_{dd} - a_s(t)) k_y^2 \\ &+ 4\pi n_0(y, t) a_{dd} \left[\frac{1}{X(y, t)^2} + \frac{1}{Z(y, t)^2} \right]. \end{aligned} \quad (17)$$

The roton population grows exponentially when the spectrum becomes imaginary, leading to the appearance of peaks at large momenta. From the local instantaneous roton spectrum we may estimate the associated peak in the momentum distribution as

$$\tilde{n}_1(k_y, t) \sim \int dy \left[e^{2 \int_0^t dt' \text{Im}[\varepsilon(k_y, y, t')]} - 1 \right]. \quad (18)$$

where we assume equal seeding for all unstable modes. We then evaluate the total population of the peak $N_*(t) = \int dk_y \tilde{n}_1(k_y, t)$, as well as the roton momentum

$$\langle k \rangle = \int dk_y k_y \frac{\tilde{n}_1(k_y, t)}{N_*(t)}.$$

The previously discussed SSM does not describe properly the evolution when the shrinking of the condensate, and hence the increase of its peak density, is too large. As the gas gets dense, quantum fluctuations, which are not considered in the SSM, introduce in our experiments an effective repulsion that crucially prevents large densities or condensate collapse (see Supplementary Information). This limits the use of the SSM to the vicinity of a_S^* . Moreover, for our tightest trap the SSM results are also unreliable, since the theory predicts condensate collapse before the momentum peak develops. That said, as shown in the main text, the SSM provides not only a clear intuitive understanding of our measurements, but to a very large extent also a very good quantitative agreement with both our experimental results and our NS based on a generalized GPE, which is detailed in Supplementary Information.

Supplementary Material

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Acknowledgements

We are particularly grateful to B. Blakie for many inspiring exchanges. We thank D. O'Dell, M. Baranov, E. Demler, A. Sykes, T. Pfau, I. Ferrier-Barbut and H. P. Büchler for fruitful discussions, and G. Natale for his support in the final stage of the experiment. This work is dedicated in memory of D. Jin and her inspiring example. The Innsbruck group is supported through a ERC Consolidator Grant (RARE, no. 681432) and a FET Proactive project (RySQ, no. 640378) of the EU H2020. LC is supported within a Marie Curie Project (DipPhase, no. 706809) of the EU H2020. FW and LS thank the DFG (SFB 1227 DQ-mat). All the authors thank the DFG/WWF (FOR 2247). Part of the computational results presented have been achieved using the HPC infrastructure LEO of the University of Innsbruck.

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Figure 1. Roton mode in a dBEC.

a, axially elongated geometry with dipoles oriented transversely. **b**, real (solid lines) and imaginary (norm of the dotted line) parts of the dispersion relation of a dBEC in the geometry **a**, showing the emergence of a roton minimum for decreasing a_s (dashed arrow). The DDI changes from repulsive (blue) to attractive (red) depending on the dipole alignment (inset). **b1**, **b2**, the dipole alignment (color code as in inset) associated to small- (**b1**) and large- k_y (**b2**) density modulations. **c**, **d**, distributions on the $k_x k_y$ -plane associated with the roton population in cigar **a** or pancake **d** geometries with an identical roton population and colorscale.

Figure 2. Observed roton peaks and characteristic scalings.

a-c, $\tilde{n}(k_x, k_y)$ obtained by averaging 15-to-25 absorption images for $(v_z, \lambda) = (456 \text{ Hz}, 14.4)$, $t_h = 3 \text{ ms}$ and: **a**, $a_s = 54a_0$, **b**, $a_s = 44a_0$ and **c**, $a_s = 37a_0$. **a1-c1**, corresponding cuts at $k_x \approx 0$ (dots) and their fits to three-Gauss distributions (lines), from which we extract k_{rot} and A^* . **d**, Measured k_{rot} depending on $1/\ell_z$ at $a_s \lesssim a_s^*$ for $v_y \approx 35 \text{ Hz}$ (circles) and $v_y = 17(1) \text{ Hz}$ (triangle). **e**, k_{rot} depending on a_s for $(v_z, \lambda) = (149 \text{ Hz}, 4.3)$ (circles). In (d-e), error bars show the 95% confidence interval of the three-Gauss fits (see a1-c1). The squares

(diamonds) show predictions from the SSM (NS) in their range of validity (Methods). Lines are guides-to-the-eye.

Figure 3. Dynamics of the roton mode.

Results depending on t_h after quenching to $a_s^f = 50a_0$ in $(v_z, \lambda) = (149 \text{ Hz}, 4.3)$. **a-c**, Δ^2 at $y = 0$, amplitude and momentum of the roton peak, from the SSM (solid lines), the NS (dashed lines) and the experiments (circles) with their 95% confidence interval from the three-Gauss fit (error bars). Shaded areas identify the dynamical destabilization. For comparison, the amplitudes are renormalized to their $t_h = 3 \text{ ms}$ -values and the NS results are translated by -3.7 ms . k_{rot} is reliably extracted from the experiments (NS) for $t_h \geq 3 \text{ ms}$ (2.5 ms). Inset, same plot in log scale. **d, e**, Integrated 1D density profile, in momentum $\tilde{n}_1(k_y)$ and real space $n_1(y)$, from one run of the NS.

Figure 4. Population growth and roton gap.

Results for $(v_z, \lambda) = (149 \text{ Hz}, 4.3)$. **a, b**, A^* depending on t_h and on the rescaled time T_h , after quenching to $a_s^f = 39a_0$ (hexagons), $a_s^f = 44a_0$ (triangles), $a_s^f = 47a_0$ (diamonds), $a_s^f = 50a_0$ (circles) and $a_s^f = 52a_0$ (squares). Shaded areas show the 95% confidence interval from the three-Gauss fit. The black line shows an exponential fit to the full dataset with $T_h < 3$ ms. **c**, Extracted $\text{Im}[\bar{\Delta}]$ depending on a_s and on t_h (inset) from the experiments (dashed line with pentagons) with the propagated errors (shaded area) from the time-rescaling analysis

and exponential fit (see b), and from the SSM (solid line). The dotted line and the corresponding shaded area show the experimental a_s^* and its fit's confidence interval. The inset shows the same configuration as Fig. 3 ($a_s^f = 50a_0$).