Observation of New Superradiant Regimes in a Bose-Einstein Condensate

Ivana Dimitrova, 1 William Lunden, 1 Jesse Amato-Grill, 1, 2 Niklas Jepsen, 1 Yichao Yu, 1, § Michael Messer, 1, † Thomas Rigaldo, 1, ¶ Graciana Puentes, 1, ‡ David Weld, 1, †* and Wolfgang Ketterle 1

1Research Laboratory of Electronics, MIT-Harvard Center for Ultracold Atoms, Department of Physics, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA
2Department of Physics, Harvard University, Cambridge, Massachusetts 02138, USA
(Dated: September 8, 2017)

New phenomena of collective light scattering are observed when an elongated Bose-Einstein condensate is pumped by two non-interfering beams counterpropagating along its long axis. In the limit of small Rayleigh scattering rates, the presence of a second pump beam suppresses superradiance, whereas at large Rayleigh scattering rates it lowers the effective threshold power for collective light scattering. In the latter regime, the quench dynamics of the two-beam system are oscillatory, compared to monotonic in the single-beam case. In addition, the dependence on power, detuning, and atom number is explored. The observed features of the two-beam system qualitatively agree with the recent prediction of a supersolid crystalline phase of light and matter at large Rayleigh scattering rates.\n
Collective scattering of light can dramatically enhance single-atom scattering and can lead to qualitatively new phenomena. Since the pioneering work of Dicke [2], it has been observed in many systems [3–6]. Due to their long coherence time, Bose-Einstein condensates are well suited for studying superradiance [7]. Furthermore, the coherence between the atoms, which is responsible for the superradiant scattering, can be directly observed from the momentum distributions and allows a detailed study of superradiance without detection of the emitted light. Different regimes and geometries have been explored by many experimental [8–20] and theoretical studies (see [21–25] and references therein). In elongated clouds collective light scattering occurs preferentially along the long axis and the superradiantly scattered light modes are referred to as the “end-fire modes”. If the pump laser beam also propagates along this axis, all relevant modes of both the light and the atoms are aligned allowing for a simple one-dimensional description.\n
Motivated by recent theoretical work [1], we have studied the effect of illuminating a Bose-Einstein condensate by two counterpropagating beams at the same intensity. Interference between the two beams can be suppressed by orthogonal polarization or rendered irrelevant by a large frequency offset. Remarkably, we never find a regime in which the two beams act independently, i.e. in which each of them independently induces superradiant scattering into its corresponding direction of propagation. At small Rayleigh scattering rates, collective light scattering is suppressed compared to the single-beam case, which is readily explained by bosonic stimulation. By contrast, when the Rayleigh scattering rate into the end-fire modes is larger than the recoil frequency $\omega_R = \hbar k^2 / 2m$, where $m$ is the atomic mass and $k$ is the wavenumber of the pump light, the threshold for collective light scattering is lowered in the presence of a counterpropagating beam. The two regimes are characterized by different microscopic mechanisms.\n
For low Rayleigh scattering rates, each scattered photon creates a quasiparticle in the form of a recoiling atom, or a phonon, with momentum equal to the difference between the momenta of the incoming and the scattered photons. Each phonon mode is a moving density modulation which is amplified by subsequent collective light scattering. For large scattering rates, it becomes possible to scatter two or several photons quasi-simultaneously. The analysis of Ref. [1] shows that this can lead to a stationary situation, where an atomic density modulation collectively creates two backscattered beams. Each of them interferes with its corresponding pump beam so that two standing waves are formed. In a self-consistent way, a stationary density modulation is stabilized by the stationary optical lattice potential, which, in contrast to usual optical lattices, consists of two different standing waves which are created and therefore phase-locked by the atoms. Above a critical pump power, the system develops an instability towards a periodic modulation and undergoes a phase transition to this crystalline phase.
contrast to the single-beam case, in which the backscattered beam is frequency-shifted by $4\omega_R$ with respect to the incident beam due to the $2\hbar k$ of recoil momentum imparted to the atoms, in the crystal phase both backscattered beams have the same frequency as their corresponding pump beams because the atomic density distribution is stationary. This establishes a new supersolid form of matter by spontaneous crystallization of light and atoms. For a sudden turn on of the pump beams, there should be an oscillatory behavior around the new equilibrium phase, in contrast to the exponential growth and eventual gain-depletion of single-beam superradiance, as in [7].

In this work, we characterize the regimes of low and high Rayleigh scattering rates by determining the threshold power for collective light scattering and by monitoring the time-evolution of the atomic momentum distribution both for a single beam and for two beams. We observe qualitatively different behaviors, including oscillatory dynamics for two pump beams at high Rayleigh scattering rates. Our observations are consistent with the predictions of the crystal phase using the 1D theoretical model [1]. However, the experimental system is 3D and its lifetime is severely limited by Rayleigh scattering into free space. Therefore, the crystal phase can only form transiently and cannot reach equilibrium.

Experiments are performed with a new $^7$Li machine which produces Bose-Einstein condensates of typically $4 \times 10^5$ atoms in 10 s. Atoms from an effusive oven are laser-cooled with a spin-flip Zeeman slower and $5 \times 10^9$ atoms are captured in a 3D Magneto-Optical trap operated on the $D_2$ line. After a compression step, further sub-Doppler cooling is performed using Grey Molasses on the $D_1$ line, as in [27], which reduces the temperature to $90 \mu K$. Dark state optical pumping on the $D_1$ line prepares the atoms in the magnetically trappable $|F=2, m_F=2\rangle$ ground state. They are then transferred to a quadrupole magnetic trap with a repulsive “plug” optical beam used to inhibit losses from the center of the trap [28]. During RF evaporation the atomic density is kept at $1 \times 10^{13}$ cm$^{-3}$ by gradually opening up the magnetic trap to prevent strong losses due to three-body recombination. The negative scattering length of the $|F=2, m_F=2\rangle$ state prevents the formation of large stable condensates [29]. The evaporation is terminated just before degeneracy is reached and the atoms are transferred to a 1064 nm optical dipole trap. They are spin-flipped to the lowest hyperfine state $|F=1, m_F=1\rangle$ state, which has a Feshbach resonance at 737 Gauss [30]. The scattering length is tuned to about $125 a_B$, where $a_B$ is the Bohr radius, and the atoms are evaporated to degeneracy.

For the current experiment, a high magnetic field is chosen to realize a scattering length of $15 a_B$ to avoid strong scattering between atoms in different momentum states. We obtained a cloud of typical size of about $20 \mu m \times 120 \mu m$ by releasing it from a crossed optical dipole trap and letting it expand into a single beam dipole trap. A tunable Ti-Sapph laser generates 671 nm light detuned by 1 to 20 GHz from atomic resonance. The two pump beams have $e^{-2}$ radii of 140 $\mu m$ and propagate along the long axis of the condensate. Interference between them is suppressed on experimentally relevant timescales by offsetting them in frequency by 160 MHz using two acousto-optic modulators. Both beams have the same polarization and drive a $\pi$-transition. Rectangular pump pulses are applied, after which the trap is suddenly switched off and momentum distributions are recorded after ballistic expansion. Momentum distributions are then characterized by the number of atoms in the satellite peaks, which are separated from the main cloud by recoil momentum.

![FIG. 2. Observation of superradiant light scattering in different regimes. Shown are momentum distributions after 2ms time of flight for three different pump times and different pump powers at 18.62 GHz red of the $^7$Li $D_2$ line. For each time, the single-beam case with pump propagating from the left (top picture, blue data) is compared to the case of two balanced beams from opposite sides (bottom, red data). The strength of collective light scattering vs. laser power is characterized by the number of atoms with momentum $2\hbar k$ (right peak). Solid lines are a guide to the eye. The images are taken at the powers indicated by the dashed vertical lines.](image-url)

The onset of collective light scattering is studied by measuring the number of atoms in the $2\hbar k$ peak in time of flight as a function of laser power (Figure 2). All three plots on the right side show that there is an effective critical power for the onset of superradiance. For long pump times (Figure 2a-d), the thresholds are lower, corresponding to small Rayleigh scattering rates. For high...
pump powers (large Rayleigh scattering rates), the dynamics occur already on short time scales and the critical powers are lower for two beams than for a single beam (Figure 2a-f). By contrast, at low Rayleigh scattering rates collective light scattering is suppressed in the presence of two beams. This is best seen in Figure 2c-d: At the same powers, at which the system exhibits strong superradiance with a single pump beam, there are no recoiling atoms visible in the presence of two pump beams.

For Rayleigh scattering rates smaller than the recoil frequency (Figure 2a-d), a quasiparticle picture can be used to describe the onset of superradiance (for discussion see e.g. [7, 20, 21, 24]). Recoiling quasiparticles are created by Rayleigh scattering into the end-fire mode, which occurs at a rate of \( R_{ef} = R N_0 f \), where \( R \) is the total Rayleigh scattering rate per atom, \( N_0 \) is the number of atoms in the condensate, and \( f \) is the effective solid angle for scattering into the end-fire mode, approximately given by \( \lambda^2/D^2 \) where \( \lambda \) is the wavelength of the scattered light and \( D \) is the diameter of the condensate. This rate is enhanced via bosonic stimulation by a factor \( N_1 + 1 \) with \( N_1 \) atoms with recoil momentum \( 2 \hbar k \) already present. The recoiling atoms are lost from the system at a rate \( L \), either by collisions or because they move out of the condensate volume.

The resulting rate equation describes both the threshold and the initially exponential gain [7]:

\[
\frac{dN_1}{dt} = R_{ef}(N_1 + 1) - LN_1 = R_{ef} + (RN_0 f - L)N_1 \tag{1}
\]

For weak pump beams and negligible source depletion, one would expect, at least in the perturbative regime, that the addition of a second counterpropagating pump beam would trigger superradiant scattering into the opposite direction. However, due to bosonic stimulation, the rates of scattering into the end-fire modes are proportional to the number of atoms in the initial and the final states. For the case of two counterpropagating beams which can transfer equal but opposite momenta, the stimulated scattering rates, which are responsible for superradiance, cancel in the rate equation

\[
\frac{dN_1}{dt} = R_1 N_0 f(N_1 + 1) - R_2 N_1 f(N_0 + 1) - LN_1 \tag{2}
\]

if the single-particle scattering rates are equal \( (R_1 = R_2) \). The remaining terms simply reflect spontaneous scattering and loss \( L \) as described above. A similar equation can be written for the \(-2\hbar k\) atoms.

Complete suppression of superradiance for the two-beam case is observed for pump times on the order of tens of recoil times \( \omega_R^{-1} = 2.5 \mu s \) (e.g. 50 \( \mu s \) case in Figure 2c-d). For even smaller pump powers and therefore longer pump times, as in the 1 ms data, the suppression is incomplete. This is possibly due to effects of decoherence, or defocusing of the recoiling atoms by atom-atom interactions, which could begin to have an effect after long pump times. For detailed experimental studies of the behavior of the single-beam system at low Rayleigh scattering rates, see [7, 9, 11, 14-17].

**FIG. 3.** Time dynamics of collective light scattering for two beams in the large Rayleigh scattering rate regime. Each beam is at power \( P \) which is above the effective threshold power \( P_0 \) at blue detuning \( \delta \) from the \(^7\)Li \( D_2 \) line. (a) Comparison of the single-beam pump to the two-beam pump. In both cases, the single-beam power is the same and \( P/P_0 = 1.2 \), where \( P_0 = 2.8 \) mW at \( \delta = 9.9 \) GHz. (b) For different initial condensate numbers \( (P/P_0 = 1.8, \text{where } P_0 = 8.6 \) mW at \( \delta = 17 \) GHz); (c) For different pump powers \( (P_0 \text{ and } \delta \text{ are the same as in (a)} \); (d) For different detunings but at constant Rayleigh scattering rate, which was measured for a single beam in each case to be \( 1.7 \times 10^4 \) s\(^{-1} \), which corresponds to \( P/P_0 = 2.3 \). Here \( \delta_0 = 9.9 \) GHz. Solid lines are a guide to the eye.

When the Rayleigh scattering rate into the end-fire modes \( R_{ef} \) becomes on the order of the recoil frequency \( \omega_R \), the quasiparticle picture can no longer be used. In this regime the system displays the opposite behavior compared to the low scattering rates regime. The presence of the second beam lowers the apparent threshold power for non-zero momentum atoms to appear in time of flight (see Figure 2d). In addition, the time dynamics of the system differ qualitatively for the single-beam and the two-beam cases (Figure 2). When the two beams are suddenly turned on, we observe temporal oscillations of the number of atoms with non-zero momentum. By contrast, with a single pump beam the number of recoil atoms grows continuously until the Bose-Einstein condensate gets depleted and the gain decreases, as shown in Figure 2a, in agreement with the predictions of eq. (1). The frequency of the oscillations is on the order of the recoil frequency \( \omega_R \). The amplitude decays with time due to the loss of atoms to Rayleigh scattering into modes other than the end-fire modes. We observe non-zero momentum atoms in time-of-flight with two beams for about 12 \( \mu s \) at an inverse Rayleigh scattering rate of
about 30 μs. The amplitude and period of the oscillations depend strongly on the initial number of condensate atoms (Figure 3b). This is characteristic of collective light scattering effects and for the two-beam system has been studied in [31]. The frequency and the amplitude of the oscillations increase with power in the pump beams as shown in Figure 4c. When the Rayleigh scattering rate is kept constant and the detuning is varied, oscillations with larger amplitude are observed at higher detunings, as evidenced in Figure 3c. This is consistent with the optical dipole potential rather than Rayleigh scattering governing the dynamics, as in the model used to describe the crystal phase in Ref [1]. A constant Rayleigh scattering rate \( R \) requires the power to scale with detuning \( \delta \) as \( P' \propto R' \delta^2 \), so that the AC Stark shift and, therefore, the potential depth \( \propto P' / \delta \propto \delta \).

In addition to the different time dynamics, the single-beam and the two-beam cases also differ in the shape of the observed momentum peaks (Figure 4). For the short times of flight used in this experiment, the observed density distributions within each momentum peak still reflect the in-trap density distribution. The shape of the \( \pm 2 \hbar k \) peaks follows closely the elongated shape of the condensate, while the shape of the \( 2 \hbar k \) peak for a single beam is shorter and more rounded. This can be explained by the inhomogeneous intensity distribution of the backscattered end-fire beam along the condensate. The power of the end-fire mode is largest where the pump beam enters the cloud, increasing the generation of recoiling atoms via Bragg scattering. In contrast, the combined potential of the two stationary optical lattices is predicted to be almost homogeneous along the condensate [1]. At even higher powers (see the rightmost column in Figure 4a), there is a backward peak of atoms with \( -2 \hbar k \) momentum in the one-beam case. It is the result of the second-order process of re-scattering of a backscattered photon. This is the Kapitza-Dirac regime, commonly associated with pump times smaller than the recoil time [32]. We also observe momentum peaks at \( 4 \hbar k \) due to higher-order superradiance as studied earlier [7]. In all cases, the two-beam geometry shows patterns distinctly different from those obtained by adding up the peaks for the two one-beam geometries, indicative of a different mechanism of collective light scattering.

We have also studied the detuning dependence of the effective threshold power \( P_0 \) for collective light scattering (Figure 4b). Due to the fine structure splitting in the excited state (10 GHz), the effective detuning \( \tilde{\delta} \) for a pump beam at frequency \( \omega \) is given by \( \tilde{\delta}^{-1} = 0.8/(\omega - \omega_{D2}) + 0.2/(\omega - \omega_{D1}) \), where \( \omega_{D1} \) and \( \omega_{D2} \) are the frequencies of the \( D_1 \) and \( D_2 \) lines and the coefficients reflect the relative strength of the dipole matrix elements. The observed thresholds are consistent with \( P_0 \propto \delta^2 \), i.e. the onset requires a critical Rayleigh scattering rate. This agrees with the threshold power predicted for the crystal phase in [1]. However, due to scatter in the data, the experimental results would also agree almost equally well with a linear fit \( P_0 \propto \delta \), i.e. the onset is driven by a critical AC Stark shift. Note that in-between the \( D_1 \) and \( D_2 \) lines optical pumping to other hyperfine states limits the gain into the end-fire modes. For the single-beam pump, this leads to Raman superradiance, as observed in [33, 34].

We have so far emphasized the qualitative differences between the one- and two-beam cases. Those cases can be connected by using two beams with imbalanced intensities. In the large Rayleigh scattering limit, we observe that the number of atoms in both satellite peaks initially goes through an oscillation, but eventually the stronger beam wins over: With pump time, one momentum peak grows in number and the other one decreases. Ref. [31] has investigated the phase diagram of the imbalanced system showing that there is a similar instability to self-organization for all values of the beam asymmetry. Further studies are needed to explore this regime.

In conclusion, we have studied collective light scattering of an elongated Bose-Einstein condensate when pumped with one or two non-interfering beams. In the regime of small Rayleigh scattering rates, the behavior of the two-beam system is qualitatively different and consistent with the incip-
dent formation of the predicted crystal phase. The study of this phase is currently limited by the large Rayleigh scattering rates into free space relative to the scattering into the end-fire modes. The superradiant gain could be increased by increasing the atomic density or decreasing the cloud diameter $D$ while keeping the Fresnel number $F \propto D^2/L_0\lambda$ on the order of one, where $L_0$ is the length of the condensate (see 23 for discussion). This would allow the study of a supersolid formed by collective light scattering, different from the supersolid recently realized by stimulating light scattering with two optical cavities 35. In particular, it would be interesting to confirm the predictions that the backscattered light is not recoil-shifted and that the phase of the density modulation is spontaneously chosen.

We would like to thank Helmut Ritsch, Stefan Ostermann, and Francesco Piazza for fruitful discussions and Georgios Siviloglou for help and advice during the initial stages of building the $^7$Li machine. We acknowledge support from the NSF through the Center for Ultracold Atoms and from award 1506369, from ARO-MURI Nonequilibrium Many-body Dynamics (grant W911NF-14-1-0003), from AFOSR-MURI Quantum Phases of Matter (grant FA9550-14-1-0035), from ONR DURIP (N00014-16-1-3141) and from an ARO seedling grant. J.A.G. acknowledges support by the National Science Foundation Graduate Research Fellowship under Grant No. DGE 1144152.

---

[1] S. Ostermann, F. Piazza, and H. Ritsch, Phys. Rev. X 6, 021026 (2016)
[2] R. H. Dicke, Phys. Rev. 93, 99 (1954)
[3] M. Gross and S. Haroche, Physics Reports 93, 301 (1982)
[4] Q. H. F. Vrehen and H. M. Gibbs, “Superfluorescence experiments,” in Dissipative Systems in Quantum Optics: Resonance Fluorescence, Optical Bistability, Superfluorescence, edited by R. Bonifacio (Springer Berlin Heidelberg, Berlin, Heidelberg, 1982) pp. 111–147.
[5] J. C. MacGillivray and M. S. Feld, Phys. Rev. A 14, 1169 (1976)
[6] K. Baumann, C. Guerlin, F. Brennecke, and T. Esslinger, Nature 464, 1301 (2010)
and light,” in review.

[32] D. Schneble, Y. Torii, M. Boyd, E. W. Streed, D. E. Pritchard, and W. Ketterle, Science 300, 475 (2003)

[33] D. Schneble, G. K. Campbell, E. W. Streed, M. Boyd, D. E. Pritchard, and W. Ketterle, Phys. Rev. A 69, 041601 (2004)

[34] Y. Yoshikawa, T. Sugiura, Y. Torii, and T. Kuga, Phys. Rev. A 69, 041603 (2004)

[35] J. Léonard, A. Morales, P. Zupancic, T. Esslinger, and T. Donner, Nature 543, 87 (2017)