Probing an ultracold-atom crystal with matter waves

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Atomic quantum gases in optical lattices serve as a versatile testbed for important concepts of modern condensed-matter physics. The availability of methods to characterize strongly correlated phases is crucial for the study of these systems. Diffraction techniques to reveal long-range spatial structure, which may complement in situ detection methods, have been largely unexplored. Here we experimentally demonstrate that Bragg diffraction of neutral atoms can be used for this purpose. Using a one-dimensional Bose gas as a source of matter waves, we are able to infer the spatial ordering and on-site localization of atoms confined to an optical lattice. We also study the suppression of inelastic scattering between incident matter waves and the lattice-trapped atoms, occurring for increased lattice depth. Furthermore, we use atomic de Broglie waves to detect forced antiferromagnetic ordering in an atomic spin mixture, demonstrating the suitability of our method for the non-destructive detection of spin-ordered phases in strongly correlated atomic gases.

The diffraction of electron and neutron matter waves from crystalline structures is a standard tool in crystallography, complementing X-ray techniques. The advent of quantum gases in optical lattices has introduced a new class of highly controllable systems that mimic the physics of solids at lattice constants that are three orders of magnitude larger, and it appears natural to consider a possible role for atomic matter-wave diffraction in the characterization of these systems.

Several approaches for detecting the spatial structure of strongly correlated phases of ultracold atoms in optical lattices have so far been taken. These include the analysis of noise correlations in momentum space, accessible after release of the atoms from the lattice, as well as dedicated in situ detection methods based on optical imaging and electron microscopy. In this context, diffraction experiments have the potential to reveal important information on the existence of long-range order, such as spin-ordered phases in atomic mixtures, in a non-destructive manner and with substantially lower experimental requirements. Here, optical and atomic matter-wave diffraction are equivalent in the sense that scattering of photons and atoms can be sensitive to both the external and internal state of atomic scatterers. However, when using matter waves, the (de Broglie) wavelength of a probe can be tuned freely by controlling its velocity, thus precluding limits on spatial resolution and also providing access to Bragg resonances without the need to vary the angle of incidence. As in the optical case, probes of high spectral brightness are readily available by using atoms from a Bose–Einstein condensate.

In this work we study the scattering of an atomic probe, consisting of one-dimensional Bose gases, from target atoms confined to an optical lattice. For weak confinement, we observe free-particle-like, one-dimensional (1D) collisions, corresponding to inelastic band-structure excitations of the target by the incident probe field. The inelastic scattering is suppressed for more deeply confined target atoms, giving way to elastic scattering and Bragg diffraction of the probe, from which the underlying crystalline order can be inferred. We use this matter-wave scattering technique to characterize the introduction of forced antiferromagnetic ordering into the system.

Matter-wave probing of an atomic crystal

Our experiments begin with a virtually pure Bose–Einstein condensate of 87Rb atoms in the (F, m_F) = (1, −1) ground state (F the hyperfine quantum number and m_F the Zeeman sublevel), prepared in an optical dipole trap of nearly isotropic harmonic trapping frequency 2π × 50 Hz in the transverse (x and y) and vertical (z) directions. Along the transverse axes, far-detuned optical lattices (λ_z = 1.064 nm) are smoothly ramped up in 200 ms, creating an array of 1D Bose gases, with a trapping frequency ω_z/2π ≈ 70 Hz along z. The final depths of these transverse lattices, z = 40 (in units of the transverse recoil energy (h/λ_z^2)/2m, with h Planck’s constant and m the atomic mass), are sufficiently deep so as to energetically restrict all dynamics to the lowest mode of the transverse potential, and to suppress tunnelling between the 1D tubes on our experimental timescales.

We next create a mixture of probe (|p⟩ ≡ |2, −2⟩) and target (|t⟩ ≡ |1, 0⟩ or |1, −1⟩, see below) species through microwave manipulation of the atoms’ internal hyperfine state. By adiabatically ramping up a magnetic field gradient, we then fully separate the probe and target atoms along the longitudinal (z) axis. We thereafter smoothly ramp up a state-selective lattice along z, formed by light of wavelength λ_z = 785 nm, between the D_1 and D_2 lines of 87Rb. This allows (refs 19–21, see Methods) a full cancellation of light-shifts for the probe atoms, while the target atoms experience an attractive lattice along z, with period d = λ_z/2 and variable depth s_z (in units of the longitudinal recoil energy E_R = (h/λ_z^2)/2m). This allows the target to be driven to a 1D Mott insulator state, while the probe remains a 1D superfluid. Finally, as detailed below, we examine scattering processes between the probe and target at a well-defined relative velocity, v_rel, related to the probe de Broglie wavelength Λ_dB = h/v_rel.

1D collisions and inelastic scattering

We begin by studying collisions between the two species, as the target atoms become localized to a state-selective lattice of increasing...
depth. In the limit of zero lattice confinement along \( z \), the collisions occurring between the \(| p \rangle \) and \(| t \rangle \) atoms are free-particle-like, with all scattering restricted to the 1D tubes. Such binary collisions have previously been studied with a single species as a quantum analogue of a ‘Newton’s cradle’ (NC; ref. 17), and with a two-species mixture in the context of damped spin impurity transport\(^{23}\). Here, we realize such collisions by accelerating the probe atoms into an initially separated sample of target atoms, which is itself at rest.

As illustrated in Fig. 1a, we set the relative velocity of the collisions by controlling the magnetic field gradient along \( z \) that initially separates the two species. The probe atoms experience a longitudinal trapping potential shifted by a distance \( \Delta z \), whereas the target atoms (\(| F, m_F = 1, 0 \rangle \), insensitive to the applied gradient) remain at the trap centre. The magnetic field gradient is quickly extinguished, and after a quarter oscillation in the trap, the probe atoms have accelerated to a nominal velocity \( v_p^0 = \omega_T \Delta z \) on interacting with the stationary (\( v_t^0 = 0 \)) target atoms. At this point, we access the momentum distributions of the probe and target atoms by imaging the probe atoms at the trap centre. The magnetic field gradient is quickly extinguished, and after a quarter oscillation in the trap, the probe atoms are displaced by a distance \( 2\mu_0 z^2 \Delta z \). Here, two distinct velocity components at 0 and \( v^0 = 2v_R \) can be seen for either species. As the target atoms are initially at rest and the accelerated probe atoms are initially at \( v^0 = 2v_R \) before colliding, the scattering spectra show energy- and momentum-exchanging binary collisions between probe and target atoms (that is, reflections in the centre-of-mass frame of atomic pairs). Whereas reflection and transmission events are indistinguishable for colliding atoms of the same spin\(^{17}\), a spin-mixture gives experimental access to the reflection probability\(^{16}\). The atoms collide with a high kinetic energy that far exceeds the mean-field energy of either species, such that the collisions have free-particle character. Using a slightly uneven mixture of the two species (3:2 target to probe), we find that 11 ± 1% (14 ± 3%) of the target (probe) atoms are reflected, in fair agreement with the calculated reflection probability of \( R = 10\% \) (ref. 16) for our system parameters.

As mentioned above, much effort has been devoted to studying collisions of lattice-free 1D Bose gases, both for the spin-polarized\(^{16}\) and spin-mixed cases\(^{23}\). We now investigate what happens when the individual dispersion relations in a collision are qualitatively different, with one species subject to a longitudinal lattice. To discuss in simple terms our expectations for the case that the target resides in the context of damped spin impurity transport\(^{23}\), we now investigate what happens when the individual dispersion relations in a collision are qualitatively different, with one species subject to a longitudinal lattice. To discuss in simple terms our expectations for the case that the target resides

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**Figure 1 | Interspecies collisions of one-dimensional bosons.**

a. Probe atoms of species \(| p \rangle \) are displaced by a distance \( \Delta z \) and then accelerated within a trap potential of trap frequency \( \omega_T \) to a final velocity \( v_p^0 = \omega_T \Delta z \). The \(| p \rangle \) atoms then interact with the stationary target atoms \(| t \rangle \) at the trap centre. b. TOF absorption images of the target (top) and probe (bottom), with velocity components at \( v_p^0 = 0 \) and \( 2v_R \) from resonant momentum-exchanging NC collisions. c. Illustration of momentum and energy exchange (red and blue arrows) in the band structure of the optical lattice, for a lattice depth \( \Delta z = 8 \) (in units of the recoil energy \( E_R \)). d. Percentage of target atoms participating in NC-type collisions, plotted as a function of the energy mismatch \( \Delta E / E_R \).

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**Caption:**

Typical TOF absorption images are displayed in Fig. 1b, for the case of free (along \( z \)) target atoms and probe atoms incident at \( v_{\text{rel}} = v_p^0 = 2v_R \) (\( z \)-lattice recoil velocity \( v_R = \hbar / 2m_d \approx 5.8 \) mm s\(^{-1}\)). Here, two distinct velocity components at 0 and \( v^0 = 2v_R \) can be seen for either species. As the target atoms are initially at rest and the accelerated probe atoms are initially at \( v^0 \) before colliding, the scattering spectra show energy- and momentum-exchanging binary collisions between probe and target atoms (that is, reflections in the centre-of-mass frame of atomic pairs). Whereas reflection and transmission events are indistinguishable for colliding atoms of the same spin\(^{17}\), a spin-mixture gives experimental access to the reflection probability\(^{16}\). The atoms collide with a high kinetic energy that far exceeds the mean-field energy of either species, such that the collisions have free-particle character. Using a slightly uneven mixture of the two species (3:2 target to probe), we find that 11 ± 1% (14 ± 3%) of the target (probe) atoms are reflected, in fair agreement with the calculated reflection probability of \( R = 10\% \) (ref. 16) for our system parameters.
in a lattice of finite depth, we describe the momentum and energy exchange in the periodic zone scheme associated with the lattice. For zero lateral depth, the dispersion relations of target and probe atoms coincide, such that the momentum exchange is resonant. However, as illustrated in Fig. 1c, for non-zero lateral depth \( s_z \neq 0 \), each collision denoted an incident probe atom from the first band to the ground band, and vice versa for the target atom involved in the collision. With an energy mismatch \( \delta E \) that increases with \( s_z \), the momentum-exchange becomes off-resonant, amounting to an inelastic band-structure excitation of the target. This process can occur as long as the energetic uncertainty \( \Delta E \sim h/\tau \) associated with the finite interaction time \( \tau \) is larger than \( \delta E \).

The observed changes of the target’s TOF spectra with increasing lateral depth are analysed in Fig. 1d. In the absence of collisions, \( s_z \)-dependent, symmetric peaks due to optical diffraction are observed at \( v_\text{f} = \pm 2v_R \), as expected (for up to \( s_z \sim 20 \)). Collisions at the incident velocity \( v_\text{i} = 2v_R \) give rise to an asymmetry between the two peaks, depending on the depth of the lattice. The observed asymmetry decays with the mismatch \( \delta E \), in qualitative agreement with our expectation. We estimate the duration of probe-target interaction to be \( \tau \sim 10 \mu s / 2v_R \sim 0.9 \) ms, giving an associated energy uncertainty \( \Delta E \sim 0.3E_R \). This value is in good agreement with the observed 1/e exponential decay constant of 0.35E

Elastic Bragg diffraction of matter waves

Whereas inelastic scattering events die off with increasing \( z \)-lattice confinement, we instead can expect to observe elastic scattering. The distribution of elastically scattered probe atoms is determined by the static structure factor \( S(q) \) of the crystalline target (where \( \hbar \dot{q} \) is the probe momentum transfer). As \( S(q) \) is given by the square of the Fourier transform of the target’s density, probe scattering can thus reveal information about the underlying spatial distribution of the target atoms. As a caveat, we briefly mention that the study of the dynamic structure factor through low-energy inelastic scattering, which could provide insight into the target’s correlation properties, is largely prevented in our system owing to dominant reflection at low energies (described below), as well as our use of a superfluid probe gas, which itself supports low-energy collective excitations. Such a study may thus be better suited to energy-deposition measurements based on optical Bragg spectroscopy.25,26

Figure 2 | Probe scattering from a crystalline target. a, Target atoms |(t), strongly confined to a moving state-selective lattice, impinge on a stationary probe |(p) at a velocity \( v_i = \hbar\omega(\Lambda/4\pi) \), where \( \delta \omega \) is the frequency detuning between the two lattice beams. b, TOF images show the |(t) and |(p) species after interaction. c, Probe TOF spectra after interaction with the target at \( s_z = 50 \), for varying relative velocity \( v_{rel} \). Each horizontal line is obtained from a TOF image as in b by integrating along \( y \) and normalizing to the total |(p) atom number. The solid white lines near \( v_{rel}/v_R \sim 2 \) illustrate the slope of the second-order Bragg resonance. The lines are extrapolations of a linear fit to the positions of maximum out-coupling: \( 1.8 \leq v_{rel}/v_R \leq 2.2 \). d, Dependence of reflection from the target on \( v_{rel} \). We count the number of atoms scattered to \( v_f \sim -v_{rel} \), that is within the dashed lines shown in c. Bragg resonances are observed for \( v_{rel}/v_R = 1, 2, 3 \), on top of a background contribution due to specular reflection, which decays with increasing \( v_{rel} \). e, Reflected probe population as a function of the time during which the crystal is moved (\( T_{\text{move}} \)), for fixed velocities \( v_{rel}/v_R = 1.5 \) (black data) and \( v_{rel}/v_R = 2 \) (red data). The data fall off at long times owing to evolution in the trapping potential. f, Percentage of Bragg-reflected probe atoms (with \( v_f = 2v_R \)) as a function of target confinement, \( s_z \). The dashed curve is proportional to the form factor \( (2 \times 2\pi / d)^2 \), calculated as a function of \( s_z \); the solid curve accounts for additional heating effects (see Methods). The black square is a reference measurement without \(|t|\) atoms. Error bars represent 1s standard deviation, with thin-line error bars in e representing estimated errors for individual points, set equal to the largest statistical error of the set.
The static structure factor for a deeply modulated target can be approximated as \( S(q) = |\sum f_j(q) e^{i q \cdot j}|^2 \), where the \( f_j(q) \) are the amplitudes of isolated scatterers at positions \( z_j = j d \), with \( j \) an integer. If one assumes them to be identical \( f_j(q) = f(q) \), as in a Mott phase with uniform filling, the structure factor is a product of two terms, \( |\sum \cos(q \cdot j)|^2 \) and \( |\sum \sin(q \cdot j)|^2 \), where \( \phi_j(z) \) describes the on-site density distribution of each scatterer. The first term determines the positions of Bragg resonances, and reveals the periodicity of the array of scatterers. The second term, the atomic form factor \( |f(q)|^2 \), or single-scatterer envelope, reveals the on-site density distribution.

To perform crystallographic measurements, we introduce a controlled method to vary the relative velocity (de Broglie wavelength) of matter waves on incident on a deeply confined crystal. Such a wavelength scan is necessary to identify Bragg resonances, which are expected when a multiple of the probe’s de Broglie wavelength coincides with the lattice spacing of the target (that is, \( 2d = n \lambda_{\text{DB}} \), or \( v_{\text{rel}} = n n_{\text{p}} \), for \( n \)th order diffraction), and more generally it allows a study of the dependence of scattering on the probe’s de Broglie wavelength. As illustrated in Fig. 2a, the target atoms (here in the \( |1, -1\rangle \) state) are loaded into a deep state-selective lattice initially at rest in the lab frame. This lattice is then made to move at a well-defined velocity \( v_1 = \delta \omega (\lambda_z/4\pi) \) by detuning the relative frequency of the two interfering laser beams that create the lattice by an amount \( \delta \omega \). The pinned target atoms follow the moving lattice minima for a total time \( T_{\text{move}} \), which includes a set delay time before the two species—initially separated—make contact. The large initial separation helps to avoid the formation of a quantum emulsion that could affect interspecies transport properties.

In Fig. 2b, we show TOF absorption images of both species following interaction, taken for the case of deep confinement (s_2 = 50) and with a velocity \( v_{\text{rel}} = 2v_{\text{B}} \). The target’s velocity distribution is very broad, as expected for a deep 1D Mott state, and is centred around the velocity of the moving lattice (\( -2v_{\text{B}} \)). The velocity distribution of the probe, initially centred around \( v_0 = 0 \), displays a peak of atoms transferred to \( -4v_{\text{B}} \). A series of such scattering spectra (integrated along the transverse \( y \) axis) is shown in Fig. 2c, for a range of relative velocities \( v_{\text{rel}} \). For each velocity, the time \( T_{\text{move}} \) is chosen so that the target enters the probe and then moves for 1.4 ms (much less than the trap period \( T_z = 2\pi/\omega_z \sim 14 \) ms to avoid the dispersion of out-coupled atoms to different velocities in the trapping potential). From the scattered-probe spectra, we observe that as well as the line of ‘transmitted’ atoms at \( v_p = 0 \), some probe atoms are out-coupled (that is, reflected). The reflection of probe atoms arises from two apparent elastic mechanisms—specular reflection and resonant Bragg diffraction. We note that both of these mechanisms are to be distinguished from free-space four-wave mixing in two-component mixtures, which does not persist when one species loses matter-wave coherence.

The specular contribution is due to the reflection of probe atoms from the potential ‘step’ of the target crystal (energy mismatch outside and inside the crystal). It is dominant at low velocities and gives way to transmission at larger velocities, where the band structure of the target is free-particle-like. Specular reflection transfers probe atoms to \( v_p = -2v_{\text{B}} \) and shows no resonant structure, whereas Bragg diffraction occurs at values \( v_{\text{rel}}/v_{\text{B}} \approx \pi \) and transfers probe atoms to \( v_p = -v_{\text{rel}} - n v_{\text{B}} \). Such a ‘locking’ to a Bragg resonance is observable through a change in the slope of the out-coupled branch of probe atoms. Results of the velocity scan, which displays these features, are shown in Fig. 2c.

In Fig. 2d we count the number of probe atoms reflected to \( v_p \approx 2v_{\text{B}} \), and plot this number as a function of \( v_{\text{rel}}/v_{\text{B}} \). Three resonant peaks are observed for \( v_{\text{rel}}/v_{\text{B}} = 1, 2, 3 \), corresponding to first-, second, and third-order Bragg processes. As well as these resonances a significant non-resonant contribution due to specular reflection is observed, which as expected decays with increasing velocity \( v_{\text{rel}} \). We fit the second-order resonance peak with a Gaussian on top of a linear slope (red line in Fig. 2d), and extract a \( 1/\sqrt{\pi} \)-width of \( 2\sigma_r = 0.3v_{\text{B}} \), which represents the in-trap velocity width of the probe species. This width is a factor of two smaller than the directly observed TOF velocity width of the probe atoms, probably owing to the effects of interaction-induced expansion during TOF (ref. 30).

Using our ability to vary the total time \( T_{\text{move}} \) during which the target moves at constant velocity, we can also study the temporal buildup of the out-coupled probe population. In contrast to optical scattering, this may in the future be used to provide ‘line-of-sight’ information in matter-wave scattering experiments. In Fig. 2e, we plot the temporal evolution of the reflected probe population at \( v_p = -2v_{\text{rel}} \), for the cases \( v_{\text{rel}} = 1.5v_{\text{B}} \) and \( 2v_{\text{B}} \). The difference between incoherent (specular reflection) and coherent (Bragg) processes should lead to differences in growth behaviour. Indeed, there are some indications of a more nonlinear initial growth for \( v_{\text{rel}} = 2v_{\text{B}} \).
than for \( v_{m \text{cl}} = 1.5 v_{\text{Bragg}} \). However, the relatively short coherence length of the probe (\( L \sim h/\hbar (\sigma r) = 0.8 \mu m \)) precludes a fully coherent temporal evolution, such that the observed growth will be mostly determined by the time-dependent overlap of the inhomogeneous profiles of the probe and target.

Finally, we recall the dependence of the static structure factor \( S(q) \) on the atomic form factor \( |f(q)|^2 \), in relation to the three Bragg resonances observed in Fig. 2d. In the harmonic approximation, the form factor is proportional to \( \exp(-q^2 \alpha_r^2/2) \), where \( \alpha_r \) is the harmonic-oscillator length characterizing the extent of the atomic distribution on each site. Although the third-order peak (\( q = 3 \)) is smaller than the second-order peak (\( q = 2 \)), the expected momentum dependence is partially obscured by the fixed probe-target interaction time, which results in lower-velocity probe atoms interacting with fewer sites of the crystalline target. However, we can directly probe the contribution from the atomic form factor, which is formally identical to the Debye-Waller factor (\( \exp(-q^2 \alpha_r^2/2) \)) describing the reduction of elastic scattering due to position fluctuations (\( \alpha_r^2 \)) of scatterers in an ionic crystal\(^1\). In analogy to the decrease of the Debye-Waller factor with increasing temperature due to thermal fluctuations, \( |f(q)|^2 \) should decrease for smaller \( d_{ij} \) due to increased zero-point motion (larger \( \sigma_r \)). Using the aforementioned probe acceleration method (with \( |1,0 \rangle \) target atoms), we can study the dependence of the second-order Bragg diffraction amplitude on \( s_j \), as shown in Fig. 2f. We observe an increase of the diffraction efficiency as the target atoms become more localized, in good qualitative agreement with the calculated dependence of \( |f(q)|^2 \) on \( s_j \) (using Bloch functions for \( \psi_0 \)). Even better agreement can be obtained by taking into account \( s_j \)-dependent heating due to spontaneous photon scattering (see Methods).

**Detecting forced antiferromagnetic ordering**

To further demonstrate that matter-wave diffraction may be used to probe non-trivial structure, we show here that it gives access to a clear signature of forced antiferromagnetic ordering\(^18\) in a mixed-spin crystal. As before, we start with a mixture of atoms in one-dimensional tubes, with a probe that experiences no longitudinal lattice along \( z \) (\( \langle p \rangle \), here \( |1,1 \rangle \) atoms, with \( \lambda_{z} \approx 788 \text{ nm} \)). However, instead of using only one extra species, our target now consists of two different states. The first one \( \{|1, -1 \rangle \} \) experiences an attractive lattice potential as before, whereas for the second species \( \{|2, -2 \rangle \} \) the potential is repulsive. Thus, the \( \{|r \rangle \} \) atoms will be drawn to the intensity maxima, whereas the \( \{|l \rangle \} \) atoms will be forced to the minima, as illustrated in Fig. 3a. While keeping the total population of these two species fixed, we vary their relative population, which can be quantified by the net ‘polarization’ of the crystal \( P = \Delta N/N, \) with \( \Delta N = N_L - N_R \) and \( N = N_L + N_R \).

When the target is ‘spin-polarized’, consisting of either all \( \{|r \rangle \} \) or all \( \{|l \rangle \} \) atoms, the situation is as before, with a lattice constant of \( \lambda_{z} = 2/400 \text{ nm} \). When both species are present, however, a new spacing of \( \lambda_{z} = 2/400 \text{ nm} \) arises between atoms in the two states. This may be viewed as a crystal with the original periodicity and a two-atom basis, or a new structure with half the lattice period, given that the probe atoms interact approximately the same with either the \( \{|r \rangle \} \) or \( \{|l \rangle \} \) atoms. This change results in a different diffraction spectrum for the matter wave probes. With the addition of a second species, as scattering centres separated by \( d \) give way to those of a smaller spacing \( d' \), the original first-order diffraction peak at \( v_{\text{rel}}/v_{\text{Bragg}} = 1 \) will be diminished, and entirely disappear if an equal mixture of the two species uniformly fills the lattice, with the new first-order peak occurring at \( v_{\text{rel}}/v_{\text{Bragg}} = 2 \).

To probe the mixed-spin crystal, we move it at a constant velocity of \( v_{\text{rel}}/v_{\text{Bragg}} = 1 \) with respect to the probe atoms, for a fixed interaction time of 1 ms. In Fig. 3a, we show probe velocity distributions for the three cases \( P = -1, 0, \text{ and } 1 \). The number of atoms reflected to \( 2v_{\text{Bragg}} \) is found to be much lower for the spin mixture \( P = 0 \) than for either of the nearly spin-polarized cases. We note an appreciable transfer even for \( P = 0 \), most probably owing to specular reflection as in Fig. 2d, although site vacancies in either the attractive or repulsive lattice may also cause Bragg diffraction consistent with the original lattice spacing. To characterize how the crystal structure changes as the population imbalance is continuously tuned, we count the number of probe atoms transferred to a velocity region around \( v_{\text{rel}} \approx 2v_{\text{Bragg}} \). As shown in Fig. 3b, the transferred population shows a distinct minimum for a balanced mixture, near \( P = 0 \). This example readily demonstrates how matter-wave scattering can be used to detect changes to the crystal structure of an ultracold lattice gas, and through species-selective scattering may eventually be used to detect quantum-magnetic spin-ordered states.

In the future, the techniques presented in this paper may be extended to the characterization of various novel states of ultracold matter, such as charge- and spin-density waves, magnetically ordered ground states of quantum gas mixtures, and even self-assembled structures such as Tonks–Girardeau gases of fermionized bosons, Abrikosov vortex lattices\(^33\),\(^34\),\(^35\), and dipolar crystals\(^36\).

**Methods**

**Mixture preparation and detection.** Starting with an optically-trapped Bose–Einstein condensate of \( ^{87}\text{Rb} \) atoms in \( |1,0 \rangle \) \( \rightarrow |1, -1 \rangle \) hyperfine ground state, the system is split into an array of one-dimensional tubes in the horizontal \( xy \) plane. Next, at a magnetic bias field of 1.7 G along \( z \), we create hyperfine state mixtures by means of combinations of microwave pulses and Landau–Zener sweeps. Different combinations of hyperfine states are used to perform particular experiments on interspecies scattering, with slight variations in the mixture characteristics. For detection, all species are absorptively imaged on the \( F = 2 \rightarrow F' = 3 \) cycling transition, concurrent with optical pumping from \( F = 1 \rightarrow F = 2 \).

The initial two-species mixture of probe \( |p = 2, -2 \rangle \) and target \( |t = 1, 0 \rangle \) atoms, used to study the \( s_j \)-dependent scattering (Fig. 1; Fig. 2f), contains a total of (1.6 ± 0.2) \times 10^3 \text{ atoms} with 60\% of atoms in the target state. For this mixture, a final magnetic bias field of 7.4 G is employed to suppress \( |1,0 \rangle \leftrightarrow |1, 1 \rangle \) spin-changing collisions. A second binary mixture of probe \( |p = 2, -2 \rangle \) and target \( |t = 1, 0 \rangle \) atoms is used in conjunction with the moving optical lattice (Fig. 2a–e). This mixture contains a total of (2.8 ± 0.5) \times 10^3 \text{ atoms} with 33\% target atoms. A three-species mixture of one probe species \( |p = 1, 1 \rangle \) and two target species \( |2, -2 \rangle \) and \( |1, -1 \rangle \) is used to study forced antiferromagnetic ordering (Fig. 3). This mixture contains a total of (1.5 ± 0.2) \times 10^4 \text{ atoms} with 50\% of the atoms in the target. The target, consisting of two different species, has a fully tunable spin composition (see Fig. 3b). The intra- and interspecies scattering lengths for all the states used are approximately equal to 5.3 \text{ nm}, which is the background scattering length for \( ^{87}\text{Rb} \) atoms.

**One-dimensional tubes.** The array of one-dimensional tubes is created by ramping up two orthogonal lasers (along \( x \) and along \( y \)) within 200 ms to depths of \( s_z = 40 \), using partial retroreflection of the laser beams of the optical dipole trap\(^12\), resulting in a final trapping frequency \( \omega_z/2\pi \approx 70 \text{ Hz} \) along \( z \), as determined by dipole oscillations. In the harmonic approximation, the transverse oscillation frequency in each tube is \( \omega_z/2\pi = 26 \text{ kHz} \), resulting in a spacing to the first allowed transverse excited mode of 2\( \omega_z/2\pi = 14 \text{ kHz} \), which is greater than the probe atoms’ kinetic energy \( T = (v_z^2/v_{\text{Bragg}})^2/2E_k \) and all other relevant energy scales (thermal as well as all intra- and interspecies interaction energies). For all the cases studied, the probe, while 1D, is not deep within the Tonks–Girardeau regime\(^35\), and is characterized by a Lieb–Liniger parameter of \( \gamma \lesssim 1 \) (refs 36,37).

**State-selective lattice potential.** The state-selective lattices along \( z \) are formed by interfering two laser beams (1/e^2 radius \( \approx 230 \mu m \), the same polarization) with tunable wavelength between the \( ^{85}\text{Rb} D_2 \) and \( D_1 \) lines to effect a light-shift cancellation for the probe atoms. In the case of a stationary target of \( |1, 0 \rangle \) atoms, the lattice is loaded to a variable depth \( s_z \) with an s-shaped curve in 75 ms and held for a further 5 ms before acceleration of the probe atoms. This lattice is formed by fully retroreflected laser light of wavelength \( \lambda = 785 \text{ nm} \) with \( \sigma \)-polarization. For the case of a moving target of \( |1, -1 \rangle \) atoms (wavelength and polarization as in the stationary case), the lattice is first smoothly ramped up in 45 ms to a depth of \( s_z = 10 \), exceeding the critical depth of the 1D Mott insulator transition, and then ramped in 5 ms to a depth of \( s_z = 50 \) to freeze the atoms to the sites of the lattice. The lattice is then moved, by introducing a relative frequency detuning of \( \delta\omega \) between the forward and retroreflected laser beams, by means of two acousto-optic modulators, which are driven by phase-locked function generators.
For the case of a spin-mixed crystal with forced antiferromagnetic ordering, the two species are quickly loaded into the state-dependent lattice in 2 ms following creation of the target mixture to avoid spatial separation in the small magnetic field gradient used to separate the probe. The lattice has modulation depths $s_\perp = 5$ and $s_\parallel = 12$, for the $|\ell\rangle = |2, -2\rangle$ (repulsive lattice) and $|\ell\rangle = |1, -1\rangle$ (attractive lattice) atoms, respectively. This lattice is formed by light of wavelength $\lambda = 788 \text{ nm}$, with a slightly elliptical polarization. After loading, it is moved at a fixed velocity as above, with a restriction to relatively low values to ensure that both target species faithfully follow the optical potential at these modest depths, as observed in their velocity distributions.

All the lattice depths are calibrated using Kapitza–Dirac atom diffraction\cite{1}, with a systematic uncertainty of about 3%. For probe atoms, the lattice potential is sufficiently ‘zeroed’, even for the greatest available optical potentials, such that no diffraction is observed for a pulsed-on lattice or for application of a linear potential gradient while the lattice is present (i.e., Rayleigh–Bragg oscillations). To further ‘zero’ the optical lattice, the probe atoms are loaded into 1D tubes, and we either minimize excitations of probe atoms undergoing dipole oscillations, or in the case of a moving lattice we move with a velocity of $v_p = v_R$, with no observable transfer of probe atoms to non-zero velocities.

Second-order Bragg peak population versus $s_\parallel$. The percentage of second-order Bragg-reflected probe atoms $N_{\text{diff}}$ plotted in Fig. 2f is determined from a fit of the TOF spectrum with Gaussian peaks for the transmitted probe around $\delta p = v_f^2 = 0$; a broad incoherent background centred at $v_p = 0$; Bragg-reflected atoms at $v_p = -2v_R$. $N_{\text{diff}}$ is determined as the amount of Bragg-reflected atoms normalized with respect to the total probe population.

**Heating effects.** The observed percentage of Bragg-diffacted atoms is in qualitative agreement with that expected from the form factor $|f(q)|^2$, up to an overall scaling factor (dashed line in Fig. 2f). However, the observed signal saturates at large $s_\parallel$ whereas the expected curve does not. We attribute this mainly to a heating of the probe, leading to reduced coherence and an increased spectral width, as also directly observed through a linear increase in the probe’s TOF velocity-within increased hold time of a deep lattice ($s_\perp = 25.5$) before probe acceleration to $v_f^2 = 2v_R$. The Bragg diffraction signal exhibits a roughly exponential decay with hold time ($1/e$-time $s_{\text{hold}} = 150–200$ ms). Assuming that contributions from heating due to Rayleigh scattering scale as the time integral over the lattice depth, we expect a correction factor of $e^{-\beta t_{\text{hold}}}$ to modify the form factor. A fit to the data with $|f(q)|^2 \times e^{\beta t_{\text{hold}}}$ yields $\beta = -0.013$ (the value used for the dashed curve in Fig. 2f), consistent with the role of such heating.

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