Ghost imaging is a counter-intuitive phenomenon—first realized in quantum optics\(^1,2\)—that enables the image of a two-dimensional object (mask) to be reconstructed using the spatio-temporal properties of a beam of particles with which it never interacts. Typically, two beams of correlated photons are used: one passes through the mask to a single-pixel (bucket) detector while the spatial profile of the other is measured by a high-resolution (multi-pixel) detector. The second beam never interacts with the mask. Neither detector can reconstruct the mask independently, but temporal cross-correlation between the two beams can be used to recover a ‘ghost’ image. Here we report the realization of ghost imaging using massive particles instead of photons. In our experiment, the two beams are formed by correlated pairs of ultracold, metastable helium atoms\(^3\), which originate from s-wave scattering of two colliding Bose–Einstein condensates\(^4,5\). We use higher-order Kapitza–Dirac scattering\(^6–8\) to generate a large number of correlated atom pairs, enabling the creation of a clear ghost image with submillimetre resolution. Future extensions of our technique could lead to the realization of ghost interference\(^9\), and enable tests of Einstein–Podolsky–Rosen entanglement\(^10\) and Bell’s inequalities\(^10\) with atoms.

Since the first experimental realization of ghost imaging with light more than twenty years ago\(^1,11\), following the original proposals some years previously\(^12,13\), there has been much interest in the technique\(^1,2\). Ghost imaging has found applications in areas including environmental monitoring\(^1\), and cryptography\(^2,6\). Extension of these techniques has enabled ghost imaging in the X-ray domain\(^12,13\), three-dimensional ghost imaging\(^14\) and even temporal ghost imaging, which could potentially lead to improved telecommunications\(^15\). Ghost imaging has produced sub-shot-noise images of weakly absorbing objects\(^16\) and, for low-light-level imaging when the number of registered photons per image pixel is less than one, has been shown to outperform conventional imaging in terms of contrast\(^17\).

There was considerable debate as to whether ghost imaging is a semi-classical or a quantum-optical phenomenon. It has been shown that ghost imaging can be realized either using a source of thermal\(^13,24\) or pseudothermal\(^25\) photons, or using correlated photon pairs (bi-photons) such as are created in spontaneous parametric downconversion (SPDC)\(^24\). Thermal or pseudothermal ghost imaging is not a quantum effect and can be explained semi-classically, whereas SPDC ghost imaging requires quantum theory to describe some aspects of its quantitative performance and can demonstrate violation of Bell’s inequalities\(^19\). However, the key enabler of ghost imaging (whether using thermal or quantum sources) is the existence of particle correlations in the source used for the ghost imaging.

Experimental ghost imaging has so far been achieved exclusively with photons. Here we report on the experimental realization of ghost imaging using massive particles. As well as demonstrating complementarity between light and matter for this phenomenon, realizing ghost imaging with atoms is a potential precursor to experiments that test fundamental concepts in quantum mechanics with massive particles, such as ghost interference, Einstein–Podolsky–Rosen entanglement and Bell’s inequalities\(^9,10,12\). Further, there could potentially be future applications such as real-time, \textit{in situ} control of atom lithography while imaging the deposition remotely via the spatial (multi-pixel) detector.

There are two main challenges to ghost imaging with massive particles: first, the need for a source of such particles with the required correlation properties; and second, the need for a sufficiently high flux of particles to provide adequate measurement statistics. Atom cooling techniques that enable the creation and manipulation of Bose–Einstein condensates (BECs) allow such a source to be created. A schematic of ghost imaging using pairs of correlated, colliding atoms is illustrated in Fig. 1.

As with optics, two types of atom sources can be envisaged. Like thermal light sources, a thermal ensemble of ultracold atoms possesses second-\((1<\rho^{(2)}<2)\) and higher-order correlations, where \(\rho^{(2)}\) is the second-order correlation function between two particles, as demonstrated via the Hanbury Brown and Twiss effect\(^26–28\); however, the flux of correlated atom pairs is usually very low. Alternatively, the degree of correlation can be enhanced by making a source comprised largely of correlated atom pairs. In our experiment, we use s-wave collisions between ultracold atoms\(^4,5\), which can yield a high degree of second-order (two-atom) correlation\((\rho^{(2)}=2)\). This set-up is analogous to correlated bi-photon sources in that pairs of correlated particles are created at the same point in space and time. Like SPDC, these sources can exhibit quantum correlations\(^29\) and demonstrate Bell’s inequalities\(^30\). In fact, the high correlations and small spatial correlation width of our source mean it would be well suited to such measurements, although in our ghost imaging experiment it is the strong correlations between atom pairs rather than any quantum entanglement that enables ghost imaging.

**Figure 1 | Schematic of atomic ghost imaging.** Correlated pairs of atoms created in a collision form two beams. One beam passes through the object to be imaged (‘O’) and the arrival times of the individual atoms are detected by a bucket detector (‘B’). The second beam never interacts with the object, but is detected with full temporal and spatial resolution by a multi-pixel detector (‘M’). A correlator (‘C’) then reconstructs the image of the object.
The experiments start with a BEC of helium atoms in the metastable \((2S_1)\) state: \(^4\)He\(^*\). This state enables single-atom detection with high efficiency because of the large internal energy of the atoms. We magnetically trap a BEC of about one million \(^4\)He\(^*\) atoms in the \(m_f = +1\) sublevel with no discernible thermal fraction (see Methods for details). To produce an \(s\)-wave halo, we collide atoms in the BEC in two steps: (i) we out-couple nearly all atoms from the trap via a two-photon Raman transition to the magnetically insensitive sublevel \(m_f = 0\) and (ii) we diffract the cloud into multiple momentum modes. This last step introduces a relative momentum difference between atoms in the different diffraction orders, which collide and thus generate a series of \(s\)-wave scattering halos via binary collisions (Methods). We use the same laser beams for both Raman and Kapitza–Dirac pulses, with the latter Fourier-broadened to ensure we are in the Kapitza–Dirac regime, which populates multiple diffraction orders. Both Raman and Kapitza–Dirac pulses propagate along the \((e_x, e_y, e_z)\) directions as shown in Fig. 2a, where \(e_x, e_y, e_z\) are the unit vectors along the \(x, y, z\) directions, respectively. These propagation directions result in momentum transfer along the vertical \(z\) axis, with the momentum difference between any two adjacent diffracted orders \(\hbar \Delta k = \pm \sqrt{2} \hbar k_0 e_x\), where \(\hbar k_0\) is a single photon recoil, \(k_0 = 2\pi/\lambda\) and \(\lambda = 1,083.1979\) nm is the photon wavelength. In \(k\)-space associated with the centre-of-mass reference frame, each \(s\)-wave scattering halo comprises atoms on a sphere of radius \(k_r = \Delta k/2 = k_0/\sqrt{2}\), which reflects energy and momentum conservation, as shown in Fig. 2a.

After the collision, the expanding halo falls approximately 850 mm under gravity onto a detector that consists of two multi-channel plates and a delay line, which records the arrival positions and time \((x, y, t)\) of individual atoms. We then transform this three-dimensional information into the momenta \((k_x, k_y, k_z)\) of individual atoms, which is plotted in Fig. 2b to show the 11 halos produced in the collision process. The dark clouds represent BECs in different diffraction orders \(\ell\) with an \(s\)-wave halo situated between each of the corresponding orders \((\ell, \ell + 1)\). The centre of each halo is on the momentum-transfer axis \(z\).

Also clearly visible are the portions of larger-diameter halos from non-adjacent orders \((\ell, \ell + 2), (\ell, \ell + 3)\) and so on, as well as halos formed by single-photon spontaneous scattering from the diffraction laser. The halo populations follow that of different orders in the Kapitza–Dirac effect and, as a consequence, the average number of atoms per halo decreases from about 370 for \((\ell, \ell + 1) = (0, +1)\) to about 50 for \((+5, +6)\).

Ghost imaging is demonstrated by placing a thin metal mask 10 mm above the detector, which covers a portion of the detector’s surface such that only a fraction of the \(s\)-wave halo (containing at most one atom

However, such a high degree of correlation usually requires a very low flux of atoms, which in turn would necessitate an extremely long data acquisition time. To overcome this problem, we developed a technique to enhance the number of correlated atom pairs, while maintaining a high degree of correlation and therefore signal-to-noise ratio. We use higher-order Kapitza–Dirac diffraction\(^6\)–\(^8\) to produce multiple distinct sources of \(s\)-wave-scattered atoms in each experimental run. Using multiple sources enables more than a tenfold increase in the data acquisition rate, thus making the experiment feasible—the data we present here represent around three weeks of full-time data acquisition.

**Figure 2 | Schematic of the experiment and resulting ghost image.**

**a.** The experiment starts by using Kapitza–Dirac diffraction laser beams (yellow arrows) to split the trapped BEC cloud (dark red sphere) into different momentum states (for simplicity, only the first diffraction order is shown as the clouds evolve in time). Binary atomic collisions populate an \(s\)-wave scattering halo (pale red sphere) with correlated pairs of opposite momenta \((\pm k)\), which then expand as they fall under gravity (grey). In the momentum space associated with the centre-of-mass reference frame, the BECs are situated at opposite poles of the \(s\)-wave scattering sphere. Some of the halo atoms pass through a mask placed 10 mm above the bucket port of the single-atom detector, while their diametrically opposing counterparts (separated by the dashed grey line) are registered by the multi-pixel port. **b.** Experimental data from 2,000 individual experimental runs showing the 11 halos produced in the collision process, where the individual atom counts are reconstructed in three-dimensional momentum space. Kapitza–Dirac diffraction produces 12 diffraction orders \(-6, \ldots, +5\) along \(e_z\). Collisions between each pair of adjacent orders result in 11 independent scattering halos. **c.** Individual ghost images from each of the halos from 68,835 experimental runs are combined to form the final image (bottom; scale bar, 5 mm). Because of the difference in absolute velocities for different diffraction orders, the halos that land first cover only a fraction of the image.

**Figure 3 | Cross-correlation function.** The main plot shows \(g^{(2)}(0, 0, \Delta z)\) as a function of only the vertical coordinate \(\Delta z\). The solid line is a Gaussian fit, which has an r.m.s. width of \(\sigma_z = 0.37\) mm, corresponding to the correlation length. Error bars show the statistical error over the 54,473 experimental runs. The inset shows \(g^{(2)}(\Delta x, \Delta y, 0)\) for the experimental data.
from each correlated pair) passes through the mask. The rest of the atoms are detected directly, without any interaction with the object, as shown in Fig. 2a. The detector is artificially split into two regions: the bucket region, masked with the object we want to image, and the multi-pixel region, where the atoms are detected with full spatial and temporal resolution. On the bucket portion of the detector, only the time of arrival \( t \) is used for ghost image reconstruction, whereas for the multi-pixel detector we retain a full \((x, y, t)\) set. The ghost image is then reconstructed in \( k\)-space using coincidence-counting between atoms in the multi-pixel port and the bucket port. Combining the independent images from halos in the different diffraction orders results in the ghost image shown at the bottom of Fig. 2c. The ghost image clearly resolves the object, which is a 15.5-mm-wide mask of the letters ‘ANU’ with line widths of approximately 0.75 mm.

Following ref. 4, we characterize the correlations between atoms with almost equal, but opposite, momenta \( k_1 = -k_2 + \Delta k \) (where \( \Delta k = (\Delta k_x, \Delta k_y, \Delta k_z) \) and the centre of each sphere has \( k = 0 \)) by constructing the two-particle cross-correlation function \( \widetilde{g}^{(2)}(\Delta k) \). Figure 3 shows the second-order spatial correlation function between the two atoms \( g^{(2)} \) for the halo from diffraction orders \( \ell = +3 \) and \( \ell = +4 \). To facilitate analysis of the ghost image resolution, we convert \( \widetilde{g}^{(2)}(\Delta k) \) into spatial coordinates, denoted \( g^{(2)}(\Delta r) \) (where \( \Delta r = (\Delta x, \Delta y, \Delta z) \); Methods). The key factor limiting the image resolution is the finite width of \( g^{(2)} \) — the correlation length.

A smaller average number of counts per halo is beneficial, because it leads to a higher peak \( g^{(2)} \) (ref. 4) and, consequently, a lower probability of registering false coincidence counts, which contribute to the ghost image background. Multi-order Kapitza–Dirac diffraction is crucial for this type of experiment, because it allows a large number of relatively dilute \( s\)-wave halos to be populated and detected in a single experimental run, thereby substantially decreasing the required experimental run-time.

In the \( x-y \) plane of the detector, the absolute image resolution is limited by the real-space widths of \( g^{(2)}(\Delta x, \Delta y, 0) \), whereas the spread along \( z \) contributes to background counts during image reconstruction. The \( g^{(2)}(\Delta x, \Delta y, \Delta z) \) data in Fig. 3 are fitted with a three-dimensional Gaussian function \( f(\Delta r) = 1 + A \exp(-\Delta x^2/(2\sigma_{\Delta x}^2) - \Delta y^2/(2\sigma_{\Delta y}^2) - \Delta z^2/(2\sigma_{\Delta z}^2)) \), yielding root-mean-square (r.m.s.) widths of \([\sigma_{\Delta x,\Delta y,\Delta z}] = [0.43, 0.39, 0.37] \pm 0.01 \text{ mm} \) (errors given here and elsewhere represent the standard deviation), which are set by a combination of the spatial correlation length and the detector resolution.

Our ghost imaging resolution can also be estimated by analysing the image of a known shape, which we do in Fig. 4a for the bars of the ‘U’, which are known to be 0.77 mm wide. Fitting the convolution of the object with a one-dimensional Gaussian point spread function (PSF) gives the r.m.s. width of the PSF as 0.40 ± 0.03 mm along the \( x \) direction (Methods). This value is in good agreement with the value of \( \sigma_{\Delta x} = 0.43 \pm 0.01 \text{ mm} \) obtained from the analysis of the \( g^{(2)} \) widths.

In Fig. 4b we plot the visibility \( V = (I - B)/(I + B) \) of the ghost image\(^{22} \), where \( I \) is the average intensity of the ghost image within the mask region (yellow lines) and \( B \) is the average intensity outside this region. We show an increase in the visibility of the ghost image as we cumulatively add contributions from different diffraction orders. The optimal, final image utilizes nine halos and has a visibility of about 35%, in agreement with a simple analytical model for our experimental parameters (Methods).

In conclusion, by using higher-order Kapitza–Dirac diffraction, we are able to extend the technique of pair production via \( s\)-wave scattering and demonstrate more than a tenfold increase in the number of correlated pairs available for each single experimental run. Because we created a source of strongly correlated atom pairs, by analogy with SPDC we are able to generate very high second-order \( g^{(2)} \) correlation values that greatly exceed the maximum thermal value of two, with maxima in our case reaching \( g^{(2)}(0, 0, 0) \approx 250 \). Using this source of correlated twin beams we performed ghost imaging with atoms, achieving good visibility and showing that the sub-millimetre resolution of the image is limited by the two-particle correlation function of the atomic momenta. Extending the techniques demonstrated in this experiment could enable fundamental tests of quantum mechanics using massive particles.\(^{9,10} \)

**Online Content** Methods, along with any additional Extended Data display items and Source Data, are available in the online version of the paper; references unique to these sections appear only in the online paper.

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Author Information Reprints and permissions information is available at www.nature.com/reprints. The authors declare no competing financial interests. Readers are welcome to comment on the online version of the paper. Correspondence and requests for materials should be addressed to A.G.T. (andrew.truscott@anu.edu.au).

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METHODS

Experimental apparatus. The He* BEC is initially trapped in a bi-planar quadrupole Ioffe-configuration magnetic trap with harmonic frequencies of $(\omega_x, \omega_y, \omega_z)\approx (15, 25, 25)$ Hz and a bias field of $B_0 = 1.31 \pm 0.01$ G along the $x$ axis. The 80-mm-diameter multi-channel plate and delay-line detector are located about 850 mm below the trap centre (corresponding to a full time of about 416 ms), with a spatial resolution in $x$ and $y$ of approximately 120 m, a temporal resolution along $z$ of approximately 2 ns and a quantum efficiency of around 10%.

Diffraction. Similarly to previous work, we use the same laser beams for both Raman and Kapitza–Dirac pulses, changing only the relative frequency detuning of the waveforms, which is set by the bias $B_0$ and the geometric angle between the beams (90°). The laser is blue-detuned by 2 GHz from the $2S_{1/2}^1 \rightarrow 2P_{3/2}$ transition. The transition in the beginning of the sequence also results in the cloud acquiring a momentum change of $\mathbf{\Delta} \mathbf{p} = g_2 \varepsilon_2 \mathbf{k}$, which sets the centre-of-mass of the cloud in downward motion. The duration of the out-coupling Raman pulse is $\tau = 1.8$ μs, with a transfer efficiency of about 95%. This pulse is immediately followed by a 1.4-μs Kapitza–Dirac pulse. The maximum intensity of each laser beam is approximately 450 mW mm$^{-2}$ for the Raman pulse and 30 mW mm$^{-2}$ for the Kapitza–Dirac pulse. Each of the pulse sequences is modulated with an overall Gaussian envelope to control Fourier broadening. Thus, the broadening of the Raman pulse was optimized to maximize the transfer to the magnetically insensible sublevel $m_f = 0$, and the Kapitza–Dirac pulse was set to maximize the number of diffraction orders generated ($L$).

In our experiment we achieved $L = 12$ orders, which consequently produced $L−1=11$ s-wave halos between the adjacent orders, as pictured in Fig. 2b. In principle, by increasing the laser power of the diffraction beams, we could raise $L$ even further, generating more halos with lower density and therefore possessing higher correlations. However, it is optimal to have $L = 12$ for the ghost imaging configuration, as we experiment because of the relatively high absolute velocity of the $^4$He atoms in the higher diffraction orders: for $\ell = \{6, 6, 6\}$, $v_z = \pm 78$ cm s$^{-1}$. There are two problems that arise for $[15, 25, 25]$ Hz: (i) fast travelling, downwards-kicked halos do not expand enough to cover the object mask before falling onto the detector (Fig. 2c), and (ii) upwards-kicked halos will be lost by the falling through the top of the vacuum chamber.

We also observe larger-diameter halos from the non-adjacent orders $(l, l + 2)$, which are reconstructed only partially because they expand beyond the detector size. In principle, $g^{(2)}(\Delta k)$ peaks should be higher for these larger halos, because the scattering modes overlap less! However, the population of these halos is too low to make them usable here.

The non-zero halo thickness $b_k$ (r.m.s.) is due to uncertainty-limited broadening from the finite size of the condensates, as well as mean-field interactions in the BEC during expansion of the halo. In our experiments, we observe $b_k/\kappa = 0.034 \pm 0.003$.

Correlation function. For independent s-wave halos originating from different diffraction orders (Fig. 2b), we construct a second-order normalized cross-correlation function in k-space:

$$g^{(2)}(\Delta k) = \frac{\langle n(\mathbf{k}) n(\mathbf{k} + \Delta \mathbf{k}) \rangle}{\langle n(\mathbf{k}) \rangle \langle n(\mathbf{k} + \Delta \mathbf{k}) \rangle}$$

where $g^{(2)}(\mathbf{k} + \Delta \mathbf{k})$ is the un-normalized two-particle cross-correlation function. Because of the spherical symmetry of the s-wave halo in momentum space, it is convenient to perform the analysis with atomic momenta. However, to allow a spatial resolution to be extracted, the final results are converted to spatial units as $g^{(2)}(\Delta \mathbf{r})$. Because we know the time-of-flight to the detector $\mathcal{T}_2$, we can change the coordinate representations of the detected atom position, from velocity (momentum) space to position space or position and time: $(v_x, v_y, v_z)$ or $(k_x, k_y, k_z, t)$. For the spatial coordinate representation for $g^{(2)}(\Delta \mathbf{r})$ in Fig. 3, we convert atom velocities $\Delta \mathbf{v} = h/\Delta m/\mathcal{T}_2$, where $m$ is the mass of a $^4$He* atom, to spatial coordinates in the $x$-$y$ plane of the detector: $\Delta \mathbf{r} = \mathcal{T}_2 \Delta \mathbf{v}$. Similar to previous work, the normalization in Equation (1) is performed by calculating $g^{(2)}(\mathbf{k} + \Delta \mathbf{k})$ across different experimental runs, because the atoms will be uncorrelated between runs.

We find that the r.m.s. widths of $g^{(2)}$ are approximately constant across all diffraction orders. Note that the correlation lengths we measure are actually a convolution of the real correlation length with the detector resolution (approximately 120 μm).

The error estimate in the numerical evaluation of $g^{(2)}$ can be expressed as a standard deviation of the atom count frequencies (which represent the numerator of Equation (1) normalized to the same denominator). The error bars in Fig. 3 represent a 2σ confidence interval. The data used to produce Fig. 3 were taken from 54,473 experimental runs under identical experimental conditions to those used to produce the ghost image, but without the object mask in front of the detector because the mask blocks a large portion of the halos.

Ghost image. The target to be imaged was made out of a laser-cut 0.4 mm-thick stainless steel sheet. A microscope image of the object used for the ghost imaging is shown in Extended Data Fig. 1.

The ghost image is reconstructed in k-space using coincidence-counting between the atoms detected in the multi-pixel port $n_{\mathcal{Q}}$ and the bucket port $n_{\mathcal{B}}$:

$$C(k_x, k_y) = n_{\mathcal{Q}}(k_x, k_y) \otimes n_{\mathcal{B}}(-k_x - k_y).$$

(2)

The correlator in Equation (2) is calculated in momentum space and accepts two counts from $n_{\mathcal{Q}}$ and $n_{\mathcal{B}}$ as a valid coincidence if they occur within the correlation length along the $z$ (time) direction, which in spatial coordinates corresponds to $|\Delta z| \leq \sigma_z$. The correlation length $\sigma_z$ is given by the $g^{(2)}(\Delta r = 0.037$ mm (see Fig. 3). If the two counts are determined to be correlated, then the coordinates $(k_x, k_y)$ on the multi-pixel detector are recorded and added to the image. Next, analogously to $g^{(2)}$, we convert the image $C(k_x, k_y)$ to spatial coordinates $C(x, y)$ in the detector plane $z = 0$. This image is shown in Fig. 2c.

Image resolution and visibility. We estimate the resolution of our ghost imaging configuration by taking the ghost image of the vertical bars from the ‘U’ of the ‘ANU’ mask (Extended Data Fig. 1), and fitting it with the convolution $Q(x) = b \ast h$ of the real object shape $h(x)$ and the Gaussian PSF $h(x) = A + C \exp \left(-x^2/(2\sigma^2)\right)$. Therefore, using the r.m.s. width $\sigma$ of the PSF as one of three free parameters of the fit, we can compare the resolution of the ghost image along the $x$ direction with the correlation length along that axis. As stated in the main text, we find $\sigma = 0.40 \pm 0.03$ mm, which is in good agreement with the r.m.s. width of $\sigma = 0.43 \pm 0.01$ mm extracted from $g^{(2)}$.

We characterize the visibility of the ghost image by comparing the average intensity $I$ within the mask region to the average intensity $B$ outside the mask:

$$V = \frac{I - B}{I + B}$$

where $I$ is calculated within the region that matches the object mask. This region is determined by manually overlapping an ‘N’-shaped region with the exact dimensions of the mask on top of the ghost image.

The visibility will be reduced from the ideal value of one as a result of three main factors. First, any atom pairs in the halo that are uncorrelated, but arrive within the same time window, will result in a false count, which will reduce the visibility; however, given the high degree of correlation present in our halos (as evidenced by the large $g^{(2)}$ peaks), this will be a small contribution. Second, there are stray background counts due to particles not generated by s-wave scattering (about 2.5 counts per halo in the region of interest) and to the dark count rate of the multi-channel plate (about 0.2 counts per halo), which can lead to false correlations. Finally, the major limiting factor to our visibility is the spatial resolution of our imaging procedure ($\sigma = 400$ μm), which is related to the detector resolution and the finite spatial correlation length, and results in blurring around the edges of the image. This blurring has a particularly strong effect on the visibility because the smallest features on the mask are within a factor of two larger than the spatial resolution.

The effect of spatial resolution on visibility is illustrated by an estimate of $V$ in the $x$ dimension, where the mask feature size $d_k$ is around twice the correlation length: $d_k \approx 2 \sigma_z$. The ratio $B/I$ for this regime can be approximated by exploiting the Gaussian shape of the convolving PSF $h(x)$ (width $\sigma_z$) as follows. If calculated over approximately equal areas of the image, both $I$ and $B$ are proportional to $\int h(x) dx$ (because the integral of the convolution here is approximately equal to the product of the integrals). $I$ is integrated over the object region $|x| \leq \sigma_z$, whereas $B$ is calculated outside the object, $|x| > \sigma_z$. Thus, we obtain $B/I = 0.68 \pm 0.02$, yielding $V = 0.36$, which is in good agreement with the visibility of $V \approx 0.35$ for the ‘N’ in Fig. 4.

Because the ghost imaging resolution $\sigma_{\text{image}}$ is a combination of the detector resolution and the spatial correlation length of the halo, reducing either would yield better visibility. In our case these contributions are roughly equal, so making either of them negligible (that is, in the limit of either perfect detector resolution or zero spatial correlation length) would improve our resolution by a factor of about two, which in the above integration would in principle increase the visibility to $V \approx 0.9$. However, in reality we would not achieve such high visibility, because false coincidences due to uncorrelated atoms or dark counts of the detector limit our ultimate visibility to $V \approx 0.8$, even for perfect spatial resolution and a negligible correlation length.

Extended Data Fig. 2 shows the ghost image visibility for images from each individual halo as a function of the average atom number $\langle N_x \rangle$ in that halo. The inset shows the corresponding object (the letter ‘N’) and its ghost image for the halo with the highest number of counts. Because of the different centre-of-mass...
velocities, the halos with the smallest time-of-flight only partially cover the mask. The first two halos to arrive onto the detector, \((+6, +5)\) and \((+5, +4)\), have a small halo population \((N_h)\) and minimal overlap with the mask, which results in a weak ghost image signal relative to the background and hence \(V < 0\). Therefore, these halos were not used to produce the cumulative ghost image shown in Fig. 4. Equation (3) gives high visibility values (many counts within the object region and a small background) even if the halos do not cover the object completely and thus do not generate a full representation of the entire image. However, we used this visibility definition to be consistent with that used in the photon ghost imaging literature\(^{22}\).

**Data availability.** The data that support the findings of this study are available from the corresponding author upon reasonable request.

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Extended Data Figure 1 | The object. Microscope image of the mask used to create the ghost image. The region indicated by the dashed line forms the vertical bars shown in Fig. 4a, which was used to determine the ghost imaging resolution.
Extended Data Figure 2 | Ghost image visibility. Visibilities (dots) for images (insets) reconstructed from each individual halo with different average numbers of atoms $\langle N_a \rangle$. Diffraction orders producing the halos are labelled as $(\ell + 1, \ell)$. The dashed curve is a guide to the eye. Error bars represent the standard error of the mean associated with the variances of the pixel values contributing to $I$ and $B$. 

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