Quantum states can be characterized by detecting all the correlations between the constituents of the system. However, the amount of information that is available scales exponentially with the size of the system. A crucial challenge in the study of quantum many-body systems is to identify and detect the relevant correlations that efficiently describe a state of matter. For example, it took more than forty years after the discovery of conventional superconductors before Bardeen, Cooper and Schrieffer came up with their explanation in terms of bound electrons or Cooper pairs. These are correlations between pairs of fermionic particles with opposite momentum that are localized at the Fermi surface in momentum space. Although it was quickly understood that pairing is the key ingredient for fermionic superfluidity and superconductivity, for many systems, most prominently high-transition-temperature (high-$T_c$) superconductors, the exact nature of correlations remains unknown.

Ultracold quantum gases are an ideal platform for the simulation and study of strongly correlated Fermi superfluids in this context. They offer a unique setting with full tunability of interactions, particle numbers and single-particle spectra combined with high-fidelity detection methods. Density and spin correlations, for example, can be accessed directly in the atomic noise in an image of an expanding gas, even without single-atom resolution. This method has been applied successfully to both bosonic and fermionic quantum gases. For lattice systems, quantum gas microscopy has become an increasingly powerful tool with which to study spatial correlations at the microscopic level.

In this work we study the emergence of fermionic pair correlations in momentum space in spatially continuous two-dimensional (2D) systems starting from the smallest possible instance. Our fluorescence imaging technique enables us to extract the spin- and single-atom-resolved momentum distribution with particle-detection fidelities comparable to those typically achieved in quantum gas microscopes (see Methods). Previously, we have used this method for small systems of indistinguishable fermions and found strong correlations in their relative positions as a manifestation of Pauli’s principle. Here we study a two-component Fermi gas with 12 particles trapped in a 2D harmonic potential and with freely tunable attractive interactions. The particles are prepared in the closed-shell ground-state configurations of the harmonic oscillator with high fidelity.

Our measurements enable us to extract the pair correlations and paired fraction as a function of attraction strength. This allows us to directly identify Cooper pairs emerging at the Fermi surface as the relevant correlations for small attraction strengths. At much stronger interactions, we also find pair correlations inside the Fermi sea, indicating a transition to molecular pairing. Our work provides a method for the single-particle-resolved study of correlations in momentum space.
It also lays the foundation for future studies in more complex settings, such as those with more particles, imbalance, or higher-temperature states, for example in the strongly correlated region of the Bose–Einstein condensate (BEC)–BCS crossover.

**Experimental set-up**

We start our experiments with a balanced mixture of two hyperfine states of \(^{6}\)Li atoms. An optical tweezer provides an approximately harmonic confinement in the radial direction with frequency \(\omega_r = 2\pi \times 1,101(2) \text{Hz}\) (see Methods section ‘Experimental parameters’). To achieve a quasi-2D confinement, we superimpose the optical tweezer with a single anti-node of an optical lattice (2D optical dipole trap) in the vertical direction, providing an axial confinement with an aspect ratio of 7.1, where the atoms remain in the axial ground state (see Fig. 1a). A spilling technique enables us to deterministically prepare the ground state of up to \(N = 10\) atoms per spin component (denoted as \(10 + 10\)) in this potential and with high fidelities\(^{23}\). For \(6 + 6\) atoms, for example, we prepare the ground state in 76(2)% of the experimental cycles and estimate that the remaining entropy per particle is approximately 0.1\,kB (ref. \(^{28}\)). The degeneracy of the \(n\)th level of the 2D harmonic oscillator is \(n + 1\) (\(n = 0, 1, \ldots\)). This leads to stable closed-shell configurations for the ground states of \(1 + 3 + 3 + 6 + 6\) and \(10 + 10\) atoms (see Fig. 1b) and with Fermi energies of \(E_F = (n + 1)\hbar\omega_r\). Here, \(n = \sqrt{(2N + 1)/4} - 3/2\) denotes the highest completely filled oscillator level, \(\hbar\) is the single-spin atom number, and \(\hbar\) is the reduced Planck’s constant. The tightly focused optical tweezer results in an anharmonicity of the 2D potential and reduced trap frequencies (~10%) for larger shell numbers (\(n \geq 2\)).

We control the interactions between the particles by applying a magnetic offset field and using a Feshbach resonance\(^{30}\). In a 2D geometry the interaction properties are fully determined by a bound state of energy \(E_\text{b}\) that is present for any attractive contact interaction strength\(^{31}\). We express \(E_\text{b}\) in units of the harmonic oscillator frequency in the radial direction, \(\hbar\omega_r\). We ensure that the binding energy is always smaller than the axial confinement \(E_F < \hbar\omega_r\) to remain in the quasi-2D limit.

**Single-particle imaging**

High-resolution fluorescence imaging allows us to extract the in situ momentum distribution of any quantum state we can prepare and at any interaction strength. First, we map the in situ momentum of each particle onto its position by a non-interacting time-of-flight expansion in 2D (see Fig. 1c). Subsequently, the position of each particle is recorded on a camera with a high-fidelity fluorescence imaging scheme\(^{26}\) (see Fig. 1d). We achieve single-atom detection fidelities in the complete field of view of 97.8(9)% (see Methods). To ensure that scattering events during the expansion do not alter the measured momentum distribution, even when starting from a strongly interacting state, we switch off all interactions at the beginning of the time-of-flight sequence. To this end, we transfer all the atoms in one of the two spin components into a third hyperfine state that does not interact with any of the initial-state atoms. The projection is driven by two copropagating Raman laser beams and is almost three orders of magnitude faster (\(T_r = 300\) ns) than any other intrinsic timescale of our many-body state and can therefore be considered as instantaneous projection (see Methods). All the momenta are expressed in natural units of the in situ harmonic oscillator potential given by \(p_{\text{HO}} = \sqrt{\hbar\omega_r} m\), where \(m\) is the mass of the \(^{6}\)Li atoms. We define the Fermi momentum as \(p_F = \sqrt{2mE_F} = \sqrt{2(n + 1)} p_{\text{HO}}\) (see Methods).

**Momentum correlations**

Each image represents a single projection of the full many-body wavefunction. We take 1,000 of these snapshots of the \(6 + 6\) atom ground state (see Fig. 1d) for each interaction strength and search for the relevant correlations. To study fermionic pairing, a natural choice is the opposite-spin, density–density correlator \(C^{(2)}\), defined as:

\[
C^{(2)}(p_{\uparrow}, p_{\downarrow}) = \langle (n(p_{\uparrow}) n(p_{\downarrow})) - \langle n(p_{\uparrow})\rangle\langle n(p_{\downarrow})\rangle \rangle.
\]

Here \(n\) denotes the density operator and (...) is the average over all images. The correlation function \(C^{(2)}\) expresses the conditional probability of finding a spin-up particle at momentum \(p_{\uparrow}\) given that a reference particle with spin down was detected at \(p_{\downarrow}\), after subtracting the contribution from single-particle densities \(\langle n(p_{\uparrow})\rangle\langle n(p_{\downarrow})\rangle\).

Because \(C^{(2)}\) depends on four coordinates, additional steps are required to display the relevant correlations. We fix the reference spin-down particle to some momentum \(p_{\downarrow}\) and plot the \(C^{(2)}\) as a function of \(p_{\uparrow}\). A binning procedure is required to extract the correlation function from the discrete experimental measurements. It is convenient to perform the binning in polar coordinates \(p_{\uparrow} + (\rho_{\uparrow}, \phi_{\uparrow})\). In Fig. 2a–j we show:

\[
C^{(2)}(p_{\downarrow}, \Delta\phi) = \int_{p_{\downarrow}} p_{\downarrow}' 2\pi \int_0^{2\pi} C^{(2)}(p_{\uparrow}, \phi_{\uparrow}; p_{\downarrow}', \phi_{\downarrow}') d\phi_{\uparrow} d\phi_{\downarrow} d\phi_{\uparrow}' d\phi_{\downarrow}'.
\]
where we integrate the momentum of the reference spin-down particle \( \mathbf{p}_o \) over a bin size of \( \Delta \mathbf{p} = \mathbf{p}_1 - \mathbf{p}_2 = \mathbf{p} \) (indicated by the horizontal bars) and \( \delta \) is the Dirac delta distribution. The black crosses in Fig. 2a–j indicate the mean momentum of the spin-down atom \( \mathbf{p}_o \) averaged over all measurements that contribute to the integral. The integrals over the angle take advantage of the radial symmetry of the system, and average over all points in the correlator with the same relative angle \( \Delta \mathbf{p} = \mathbf{p}_\perp \) (for more details, see Methods). The density plots of \( \mathcal{C}^{(2)} \) visualize at what momentum \( \mathbf{p} \), the probability of detecting spin-up particles is enhanced when a spin-down particle is present at the reference spin-down momentum. The measurements reveal how pairing emerges in the ground state of the mesoscopic gas as the attraction strength is increased. Inside the Fermi sea, correlations are strongly suppressed when \( E_B \leq E_F \) (see Fig. 2a–d). However, at the Fermi surface a clear correlation peak appears as soon as the binding energy is of the order of the single-particle gap \( E_B \approx E_{\text{gap}} = \hbar \omega_r \) (Fig. 2f–i). The correlation peak increases with binding energy and is strongest at the Fermi momentum \( \mathbf{p}_F \) and at \( \Delta \phi = \pi \) directly opposite to the reference spin-down particle. This demonstrates that we can observe Cooper pairs directly in a strongly correlated mesoscopic Fermi gas in 2D. By increasing the binding energy much further to \( E_B \gg E_F \), we are able to enter the regime of more tightly bound dimers where pair correlations emerge also inside the Fermi sea (Fig. 2e, j). For the data points in the molecular regime, we have reduced \( \omega_r \) to \( 2\pi \times 343(5) \) Hz to ensure that the condition \( E_B < \hbar \omega_r \) remains fulfilled.

**Fig. 2** | Density–density correlation function. a–j. The density plots show \( \mathcal{C}^{(2)} \), the normalized probability of detecting a spin-up particle at momentum \( \mathbf{p}_+ \), given that a spin-down particle was detected with the momentum \( \mathbf{p}_- \). The cross shows the mean momentum of the spin-down atom \( \mathbf{p}_- \) (indicated by the horizontal bars). The reference particle is located inside the Fermi sea (\( \mathbf{p}_o \) between \( \mathbf{p}_1 = 0.5\mathbf{p}_F \) and \( \mathbf{p}_2 = 1.5\mathbf{p}_F \)). Notable correlations appear only at large binding energies (\( E_B/E_F > 1 \)). f–j. Positive correlations are visible already at \( E_B/\hbar \omega_r = 1 \), when the reference particle is placed at the Fermi surface instead (\( \mathbf{p}_1 = 2\mathbf{p}_F \) to \( \mathbf{p}_2 = 3\mathbf{p}_F \)). They are located opposite to the spin-down particle in momentum space \( \Delta \phi = \pi \) and, in the weak binding regime, present only at the Fermi surface (b, g and c, h). This identifies them as Cooper pairs. k–o. The opposite momentum pair density \( \mathcal{C}^{(2)}(\mathbf{p}) \) is shown. BCS theory (solid blue line) correctly predicts the appearance of a correlation peak at \( \mathbf{p}_F \). For a system of non-interacting 2D molecules in the ground state, we expect a very strong correlation peak centred around zero momentum (solid black line). Note the different vertical scale for o. The error bars represent the standard error of the mean.

**Emergence of pairing**

Between different spins there are considerable second-order correlations only between particles with opposite momenta in our system (see Fig. 2a–j). To get a more quantitative picture of how pairing emerges, we therefore extract the opposite momentum pair density \( \mathcal{C}^{(2)}(\mathbf{p}_1 \to \mathbf{p}, \mathbf{p}_2 \to -\mathbf{p}) \), as defined by equation (1) (see Fig. 2k–o). Owing to the radial symmetry of the system, \( \mathcal{C}^{(2)}(\mathbf{p} \to -\mathbf{p}) = \mathcal{C}^{(2)}(\mathbf{p}) \) depends only on the magnitude of \( \mathbf{p} \). The total number of pairs in the ground state can then be extracted by integrating over \( \mathcal{C}^{(2)}(\mathbf{p}) \) (see Fig. 3). The measurements reveal how the ground state transforms from the non-interacting, completely unpaired state to a paired system.

We identify three different regimes of pair correlations. The weakly paired regime \( E_B \leq \hbar \omega_r \), the regime of intermediate interaction strength \( E_B = E_F \geq \hbar \omega_r \), and the limit of strong binding \( E_B \gg E_F \). In the regime of weak interactions only a small fraction of the system shows pair correlations. For the largest accessible binding energies of \( E_B = 15.9 \hbar \omega_r \), the number of total pairs is 3.4(1), closer to the maximum possible value of 6 for
the $6+6$ particle system (see Fig. 3). Here, the interactions between the bosonic pairs are still very large and this measurement is still in the strongly correlated regime of intermediate interactions. To reach the strong binding limit described by point-like molecules without interactions we would have to increase $E_B$ even further.

We compare the measurements to standard BCS theory and to a model system of $N = 6$ non-interacting 2D dimers in the harmonic oscillator ground state. We expect these mean-field descriptions to become accurate when $N \rightarrow \infty$ and in the limits of weak and strong binding, respectively. BCS theory can qualitatively explain the presence of a correlation peak at the Fermi surface that we find as the main feature of our system in the experiment (see Fig. 2–1). For binding energies much larger than the single-particle gap $E_B \gg \hbar \omega$, we find that the correlations become much stronger and their maximum shifts towards smaller momenta (see Fig. 20). This qualitatively agrees with the expectation for a system of tightly bound 2D dimers in the ground state of the harmonic oscillator where the correlation peak is centred around zero momentum (solid black line). Both mean-field descriptions generally fail to produce accurate quantitative predictions, indicating that both beyond-mean-field and finite-size effects are present in our experiment.

The mesoscopic Fermi gas in a 2D harmonic oscillator is closely related to superconducting grains, quantum dots and systems from nuclear and atomic physics\textsuperscript{3,5,32}. When the coherence length approaches the system size, quantum confinement effects become important and lead to a discrete single-particle spectrum. As soon as the level spacing becomes of the order of the many-body gap $\Delta$, superfluidity breaks down and the system remains in the normal phase, even at zero temperature\textsuperscript{5}. Owing to the small and fixed particle number and discrete spectrum, mean-field approaches generally break down and descriptions in terms of local quantities, such as the conductance, become impossible. Instead, the sample has to be treated as a whole. In the closed-shell configurations we study in our experiment, all the levels up to the Fermi energy $E_F$ are already occupied and there is a gap of $\hbar \omega_p$ to the next unoccupied levels. In the thermodynamic limit, when $N \rightarrow \infty$, this leads to a phase transition from a normal to a superfluid phase at some critical value for the binding energy $E_B^c$. A precursor of this phase transition can be observed at the mesoscopic scale\textsuperscript{28,34}. The critical value for $N=6$ particles is predicted as $E_B^c = 0.78 \hbar \omega$ from an exact diagonalization of the Hamiltonian\textsuperscript{34}. In Fig. 3 we plot BCS theory shifted by the critical binding energy $E_B^c$ as a first-order approximation of the finite-size effects (dotted line). In the weakly paired regime this explains the increase of the pair number of the closed-shell ground state as a function of the binding energy $E_B$ (Fig. 3, inset). Owing to the small particle number the transition is much smoother than the sharp increase at $E_B^c$ that is expected for larger systems. The large single-particle gap enables us to study the weakly paired regime at much larger absolute binding energies ($E_B = \hbar \omega$) and temperatures than what would be required for macroscopic samples. When the attraction strength is increased further, we enter the strongly correlated regime and the measured number of pairs increases above the mean-field prediction (see Fig. 3). Here, fluctuations of the many-body gap beyond the mean-field value $\Delta$ have to be considered for a more accurate quantitative prediction\textsuperscript{5}.

When the particle number is increased $N \rightarrow \infty$, we expect the correlations to become even sharper peaked around the Fermi surface in the weakly interacting limit. The limiting cases of infinite and weak attraction converge against the mean-field description. In the regime of intermediate interactions, new theories are required for a quantitatively accurate prediction of the pair correlations. The precise measurements of correlations, also beyond second order, in our experiment can be used as important benchmarks for new numerical and analytical approaches in the future.

**Outlook**

In conclusion, we have demonstrated that single-particle-resolved correlations can be accessed in continuous systems that are strongly interacting. We directly observe how Cooper pairs emerge at the Fermi surface. The correlation signal as a function of momentum and attraction strength allows us to characterize the ground state and identify different pairing mechanisms. An even more thorough characterization of quantum many-body states will become possible when extending our imaging scheme to additionally detect the single-particle-resolved population of the closed-shell ground state as a function of the binding energy $E_B$ (Fig. 3). Here, fluctuations of the many-body gap beyond the mean-field value $\Delta$ have to be considered for a more accurate quantitative prediction\textsuperscript{5}.

**Online content**

Any methods, additional references, Nature Research reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at https://doi.org/10.1038/s41586-022-04678-1.

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Article

Methods

Preparation sequence
A more detailed explanation of the experimental sequence can be found in ref. 29. We apply the same scheme here to prepare the closed-shell ground-state configurations of up to 10 + 10 particles of the 2D harmonic oscillator. We start by transferring a cold gas of $^6$Li atoms from a magneto-optical trap into a red-detuned optical dipole trap. A radio frequency pulse sequence is applied to create a balanced mixture of atoms in hyperfine states $|1\rangle$, $|3\rangle$, $|5\rangle$, and $|7\rangle$ (see Extended Data Fig. 1a). Next, a tightly focused optical tweezer is loaded from the optical dipole trap and several evaporation stages are used to create a deeply degenerate gas of around 350 atoms in the optical tweezer. A precise spilling method, discussed in detail previously 29, results in about 30 atoms in the optical tweezer with all levels up to the Fermi surface filled with very high probabilities.

The optical tweezer has a quasi-1D aspect ratio of $\omega_z/\omega_r = 5.1$. To create a 2D sample, we perform a continuous crossover to a quasi-2D confine-

 levels up to the Fermi surface filled with very high probabilities.

 Imaging sequence
A detailed sketch of the imaging sequence is shown in Extended Data Fig. 1. It can be separated into two parts. First, the free time-of-flight (TOF) expansion, where the cloud size increases by a factor of approximately 50 and the in situ momenta of the particles are mapped onto their position. Second, the image acquisition itself, where the atoms are excited by resonant laser beams, start to fluoresce and their positions are recorded on a camera.

The TOF sequence begins after the ground state with the desired particle number and interaction strength has been prepared. The first step is to switch off the optical tweezer, which radially confines the cloud. This leads to a quick expansion of the gas in the 2D dipole trap (see Fig. 1a, c). The expansion serves two purposes: First, the increased distance between the single atoms allows us to resolve them in the first place. Second, it maps the in situ momentum of each particle onto its position. For the study of fermionic superfluidity, imaging in momentum space is advantageous over position space because it offers more direct access to the relevant correlations—such as Cooper pairs—of the gas. In the future, we plan to extend our scheme to enable us to additionally take images of the in situ density with single-particle resolution 29.

To ensure that the mapping into momentum space by the TOF is accurate and that we measure genuine in situ correlations, no scattering events may occur during the expansion. To this end, we use two Raman laser beams to quickly switch from the strongly interacting $|1\rangle$–$|3\rangle$ to the almost non-interacting $|1\rangle$–$|4\rangle$ mixture in $T_e = 300$ ns (see Extended Data Fig. 1a, b). Between states $|1\rangle$ and $|4\rangle$ there is no Feshbach resonance present and all our measurements are consistent with their scattering length $a_s$, being very close to zero. By checking for interaction shifts in the spectrum of two particles in a harmonic trap we determined an upper limit for the scattering length $|a_s| < 300 a_0$, where $a_0$ is the Bohr radius. This sets an upper limit of one scattering event between two atoms in 50 experimental runs for our parameters when expanding in the $|1\rangle$–$|4\rangle$ mixture.

To check that the switch into the almost non-interacting state is fast enough, we have studied the effect of its duration on the measured correlation signal (see Extended Data Fig. 3). The measurement was taken at a fixed in situ interaction strength $E_r/\hbar \omega_r = 0.6$. We find very good agreement of the decline in correlation strength with a model that assumes that each scattering event destroys all in situ correlations of the two participating atoms. The model contains no free parameters and depends only on the scattering rate $\lambda_s$ of the $|1\rangle$–$|3\rangle$ mixture at the magnetic offset field ($B_0 = 750$ G) and the in situ density. It enables us to extrapolate to measurements taken at the highest in situ scattering rates (dashed line). We conclude that for $T_e = 300$ ns, scattering during the TOF can be neglected for all interaction strength settings.

To image the two spin components of our gas on our camera we make use of the free-space imaging scheme discussed in detail in ref. 29. It enables us to image single atoms in free space with high fidelity and without any confining potential or cooling required. Two counterpropagating illumination beams on the D2 line are used to excite the $^6$Li atoms and the emitted fluorescence light is collected through an objective (numerical aperture (NA) = 0.6) on an EMCCD camera (see Fig. 1c). Each image is exposed for 15 μs and we collect around 20 photons per atom on the camera. This leads to single-atom detection efficiencies of the order of 98% (see Extended Data Fig. 4). The different spin components are resolved by taking two images in quick succession (see Extended Data Fig. 1). A sequence of radio frequency transitions is used to transfer atoms from each spin component to state $|3\rangle$ prior to their measurement. The latter has a closed imaging transition and is therefore best suited for high-fidelity imaging. When the first image is taken, the illumination beams are resonant only to atoms in state $|3\rangle$ and the other atoms in state $|4\rangle$ are so far detuned (~2 GHz) that off-resonant scattering is negligible.

Data analysis
From the EMCCD camera we obtain binary images with bright pixels where one or more photons hit the camera chip. To analyse the images we first apply a low-pass filter (see Extended Data Fig. 4c). We search for peaks in these images and count all peaks above an optimized amplitude threshold as atoms (see Extended Data Fig. 4). The position of each atom in pixels is finally mapped to its in situ momentum (see Methods section ‘Experimental parameters’). This results in a list of all momenta $p_{ij}$ of all atoms $i = 1, \ldots, 6$ in each experimental run $r = 1, \ldots, N_{\text{exp}} \geq 1,000$ and for different interaction strengths $E_r$ (see Extended Data Figs. 5, 6). By taking the average over all images, we obtain the single-particle momentum space distribution and the average kinetic energy of the sample (see Extended Data Fig. 7a–c).

To extract the correlation function $C^{(2)}$, we transform to polar coordi

nates $\mathbf{p} \to (p, \phi)$ such that $C^{(2)}(\mathbf{p}, \mathbf{p'}) \to C^{(2)}(p, \phi, \mathbf{p}, \mathbf{p'})$. Owing to the radial symmetry of our system, $C^{(2)}(p, \phi, \mathbf{p}, \mathbf{p'}) \equiv C^{(2)}(p, \phi, \mathbf{p}, \Delta \phi)$ depends only on the difference between the angles of both particles $\Delta \phi = \phi - \phi'$, but not the absolute values of $\phi$ and $\phi'$. We make use of this symmetry and integrate over all measurements with the same $\Delta \phi$ to increase the signal-to-noise ratio. Finally, we bin the data according to the momentum of the reference particle $p_r \to p_r' \in [p_r, p_r]$ in momentum bins with a width of $\Delta p = p_{10}$. Different choices for the momentum of the reference particle $[p_r, p_r]$ correspond to different 2D slices through the 4D correlation function $C^{(2)}$ (compare Fig. 2a–j).

The slices can be expressed as $C^{(2)}(p_r, p_r', \Delta \phi) \equiv C^{(2)}(p_r, \Delta \phi)$, with

$$C^{(2)}(p_r, \Delta \phi) = \frac{1}{\Delta p} \int_{p_r}^{p_r'} \int_{-\Delta \phi}^{\Delta \phi} C^{(2)}(p_r, \phi, p_r', \phi') d\phi d\phi'$$

and are shown in Fig. 2. The unnormalized correlation function $C^{(2)}(p_r, \Delta \phi)$, defined as above but without subtracting the $(n/n)$ term (see equation (1)), is shown in Extended Data Fig. 7d for comparison.
An alternative representation of the correlation function is possible by transforming to relative \( \mathbf{p}_R = \mathbf{p}_0 - \mathbf{p}_1 \) and centre-of-mass \( \mathbf{p}_{\text{CM}} = (\mathbf{p}_0 + \mathbf{p}_1)/2 \) coordinates (see Ref. 1). Integrating over either one of the coordinates, we obtain the pair-correlation function expressed in the relative \( C_{\text{RR}}(\mathbf{p}_R) = \int C_{\text{RR}}(\mathbf{p}_0, \mathbf{p}_1) \, d\mathbf{p}_0 \), or centre-of-mass \( C_{\text{CM}}(\mathbf{p}_{\text{CM}}) = \int C_{\text{RR}}(\mathbf{p}_0, \mathbf{p}_1) \, d\mathbf{p}_1 \), momentum coordinates (see Extended Data Fig. 8). The result is completely equivalent to the data in Fig. 2. In the relative coordinate system, Cooper pair correlations appear at \( |p_R| > 2p \) (see Extended Data Fig. 8a–e). In the correlation function \( C_{\text{RR}}(\mathbf{p}_R) \), the pairing signal is spread out over a much larger area and the function is therefore much more sensitive to noise. In addition, it is not possible to detect pairing inside the Fermi surface as \( E_F, \rho \rightarrow \infty \) (Extended Data Fig. 9d, e). The correlation function is sensitive only to the number of pairs present at a given relative momentum \( |p_R| \) and not to their relative angle \( \Delta \phi \). The centre-of-mass correlation function \( C_{\text{CM}}(\mathbf{p}_{\text{CM}}) \) reveals the emergence of pairing as \( \rho \rightarrow \infty \) (Extended Data Fig. 8f–j). We find a sharp peak centred at zero centre-of-mass momentum of the pairs with a weight that increases with \( E_F \) as in Fig. 3.

For Fig. 2a–j we postselect for only those runs where all atoms were detected in both images. In Figs. 2k–o, 3 we also show data where up to three of the atoms were missed. This greatly reduces the statistical errors of our results (owing to the increased number of images used) and we have checked that we find no qualitative difference in the results for any of our measurements compared to a strict postselection. The postselection rate of ~5% is an order of magnitude smaller for \( 6 + 6 \) molecules than in the weakly interacting limit (see Fig. 2a–j). The centre-of-mass momentum of the pairs fluctuates on the order of \( \rho_{\text{MC}} \), explaining why they are not always detected with exactly opposite momenta (see Extended Data Figs. 5, 8–f). Our definition of \( \rho_{\text{MC}} \) ensures that the correct value is reached in the limit \( N \rightarrow \infty \) in a homogeneous system or when a local density approximation becomes applicable. The ambiguity in the definition of \( \rho_{\text{MC}} \) and the fluctuations of the pair centre-of-mass momentum do not affect the interpretation of our measurements in the mesoscopic system. As the particle number is increased, we expect that the relative momentum uncertainty reduces continuously until it vanishes in the thermodynamic limit and at zero temperature: only pairs with zero centre-of-mass momentum remain.

We estimate that the temperature of our initial state is very low and the entropy per particle is of the order of 0.15. (kB Boltzmann’s constant) When increasing the energy or temperature of the initial state by modulating the radial confinement, the amplitude of the pair correlation reduces considerably (see Extended Data Fig. 9).

### BCS theory

It is straightforward to calculate the density–density correlator \( C_{\text{RR}}(\mathbf{p}_0 - \mathbf{p}_1) \) as defined by equation (1) in BCS theory. We recall the Bogoliubov transformations:

\[
c_{\mathbf{p}_1} = u_{\mathbf{p}_1} v_{\mathbf{p}_1} - v_{\mathbf{p}_1}^* u_{\mathbf{p}_1}^*,
\]

\[
c_{\mathbf{p}_2} = u_{\mathbf{p}_2} v_{\mathbf{p}_2} + v_{\mathbf{p}_2}^* u_{\mathbf{p}_2}^*.
\]

with \( u_{\mathbf{p}}^2 = (1 + \xi_{\mathbf{p}}/E_F)/2 \) and \( v_{\mathbf{p}}^2 = (1 - \xi_{\mathbf{p}}/E_F)/2 \). Here, \( c_{\mathbf{p}, \uparrow} (c_{\mathbf{p}, \downarrow}) \) is the fermionic creation (annihilation) operator of the particle with momentum \( \mathbf{p} \) and spin \( s \). The BCS spectrum is given by \( E_F = (\xi_{\mathbf{p}_0}^2 + \Delta^2)^{1/2} \) with \( \xi_{\mathbf{p}_0} = \hbar \omega_{\mathbf{p}}/2m - E_F \) and for the pairing energy \( \Delta \) we use the mean field value \( \Delta = \langle \xi_{\mathbf{p}_0}^2 - E_F \rangle \). Making use of the fact that the BCS ground state is free of excitations \( \{ c_{\mathbf{p}, \uparrow}, c_{\mathbf{p}, \downarrow} \} \), we arrive at

\[
C_{\text{BCS}}(\mathbf{p}_0 - \mathbf{p}_1) = (c_{\mathbf{p}_1}^+ c_{\mathbf{p}_2} + c_{\mathbf{p}_1} c_{\mathbf{p}_2}^+) - (c_{\mathbf{p}_1}^+ c_{\mathbf{p}_2} + c_{\mathbf{p}_1} c_{\mathbf{p}_2}^+) = N^2 \Delta^2/4(\xi_{\mathbf{p}_0}^2 + \Delta^2).
\]

Here, the normalization factor \( N \) is chosen such that we obtain the correct total particle number at zero interactions (\( \Delta = 0 \)):

\[
N_{\text{tot}} = \int \left| c_{\mathbf{p}_1}^+ c_{\mathbf{p}_2} + c_{\mathbf{p}_1} c_{\mathbf{p}_2}^+ \right| d\mathbf{p}_0 = 2mN \int 0^2 v_{\mathbf{p}}^2 d\mathbf{p}.
\]

### 2D molecules

A simple model for our system in the regime of strongest interactions (\( E_F \gg E_F \)) is to assume that all the particles form bosonic dimers that occupy the \( n = 0 \) ground state of the 2D harmonic potential. Following
ref. 43, an ansatz for the two-body molecular wavefunction outside the scattering potential and in relative coordinates is

$$\Psi_{\text{rel}}(r) = \begin{cases} a_1 e^{-r/r_0}, & \text{for } r > r_0 \\ -\log \frac{r}{r_B} + a_2, & \text{for } r_0 \leq r \leq r_B \\ a_3 e^{-(1/2)B r} + a_4, & \text{for } r < r_0 \end{cases} \quad (7)$$

We choose the constants $a_i$ such that the wavefunction and its first derivative are continuous and properly normalized. Here, $r$ is the interparticle distance and $r_B$ is the molecular binding length defined as $r_B = \hbar \sqrt{2mE_B}$, with atomic mass $m$ and two-body binding energy $E_B$. We have inserted a short-distance cut-off $r < r_0 = 0.1 r_B$ to regularize the divergence of the logarithmic part for $r \to 0$. We have checked that this cut-off does not affect the calculated pair correlations at small momenta. The wavefunction for a single molecule is the product of the relative wavefunction $\Psi_{\text{rel}}$ and the centre-of-mass wavefunction $\Psi_{\text{com}}$. $\Psi_{\text{com}}$ is given by the two-dimensional ground-state wavefunction of two particles in a harmonic oscillator. The total wavefunction $\Psi_{\text{total}}(x_1, x_2, y_1, y_2)$ thus depends on four variables, namely the coordinates of the two particles. To calculate the pair-correlation signal that we expect for this trial wavefunction, we first perform a numerical Fourier transform in four dimensions. This allows us to directly calculate $C^{(2)}(\mathbf{p} - \mathbf{p})$ for a single dimer as defined in equation (1). Our model assumes that the molecules are completely independent. The total correlation function $C^{(2)}(\mathbf{p} - \mathbf{p})$ for $N + N$ particles is then just given by the summing over the contributions of $N$ single molecules.

**Data availability**

The data that support the findings of this study are available from the corresponding author upon request. Source data are provided with this paper.

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**Author contributions** M.H., L.B. and K.S. performed the measurements. M.H., L.B., K.S. and S.B. analysed the data. C.H. and M.H. performed the numerical calculations. P.L. set up the phase-locked loop for the Raman beams. P.M.P. and S.J. supervised the project. M.H. wrote the manuscript with input from all authors. All authors contributed to the discussion of the results.

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**Additional information**

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Extended Data Fig. 1 | Sketch of the TOF imaging sequence. a, The TOF imaging scheme is initiated by switching off the trapping potential and the interparticle interactions. To this end, we shine in two copropagating Raman laser beams to quasi-instantaneously transfer all atoms from state $\ket{3}$ to state $\ket{4}$ (step 1). After the Raman transition, we ramp the magnetic field from the value $B_0$ that sets the in situ interaction strength to a constant value of $B=750$ where the fidelity of the following spin flips is maximal. Only the imaging transitions of state $\ket{3}$ and $\ket{6}$ are closed. Two successive radio frequency (RF) Landau–Zener sweeps are applied to move all the atoms from state $\ket{1}$ to $\ket{3}$ during the free expansion (steps 2, 3). This is followed by taking the first image with illumination beams that are resonant only to state $\ket{3}$ (step 4). Before taking the second image, a microwave (MW) Landau–Zener sweep transfers the remaining atoms in $\ket{4}$ to state $\ket{3}$ again (steps 5, 6). The microwave pulse leads to higher transfer fidelities than the Raman lasers but is much slower. b, The duration of the initial Raman flip with $T_\pi = 300$ ns is chosen as fast as technically possible to prevent any scattering between atoms from occurring during the expansion. The radio frequency flips are solely optimized for transfer fidelities and are distributed over the remaining time of $T_{\text{TOF}} = 9$ ms. The time of the microwave flip is set by the maximal frame rate of the camera. c, From a single experimental run we obtain two binary images where bright pixels indicate that at least one photon hit the chip at this location. In the first step we apply a low-pass filter to these images. A simple peak detection with an optimized acceptance threshold then allows us to extract the position of all spin-up and spin-down particles.
Extended Data Fig. 2 | Different particle numbers. In our experiment, we are able to prepare closed- and open-shell configurations with different particle numbers. a–c, Single momentum space projections of the three lowest closed-shell ground states with 3 + 3 (a), 6 + 6 (b) and 10 + 10 (c) particles. The dashed circle indicates the corresponding Fermi momentum for each particle number. d, The total weight in the pair-correlation peak of $C^2$ is plotted versus particle number and for the setting where the reference particle is fixed at the Fermi surface (for 6 + 6, see Fig. 2b). We find that the absolute number of paired atoms increases, and their fractions remain approximately constant from $N = 3 + 3$ to $N = 10 + 10$ particles. The error bars represent the standard error of the mean.
Extended Data Fig. 3 | Scanning the interaction switch-off time. $T_\pi$ is the duration of the spin flip from a strongly interacting $|1\rangle \rightarrow |3\rangle$ to an (almost) non-interacting $|1\rangle \rightarrow |4\rangle$ mixture at the beginning of the TOF sequence (inset). When we increase $T_\pi$, the magnitude of the pair correlations reduces greatly above a threshold of $T_\pi = 100 \mu$s (red circles). The reason is that scattering events during the TOF expansion redistribute the momenta between the participating atoms and destroy the correlations that were present between the in situ momenta. We model the effect by assuming that each scattering event between two atoms annihilates all the in situ correlations for those particles. The only parameters that enter the model (solid line) are the scattering rate $\lambda_{sc}$ of the $|1\rangle \rightarrow |3\rangle$ mixture at the magnetic offset field ($B_0 = 750$) and the in situ density. The dashed line shows the same model prediction but at one of the highest scattering rates used in our experiments ($B_0 = 695$, $E_B/\hbar \omega_r = 1.97$). It follows that at the spin flip time of $T_\pi = 300$ ns that we use for our experiments, no scattering is expected to occur during the TOF at any interaction strength setting. The mean kinetic energies $E_{kin}$ (blue squares) show a similar dependence on $T_\pi$. This indicates again that only for $T_\pi \approx 1/\lambda_{sc}$ (with $\lambda_{sc} \approx 50$ kHz) the true in situ momentum distribution is obtained after the TOF sequence. The error bars are obtained from the counts in each bin of the momentum distribution and the correlation function respectively and by assuming Poissonian statistics.
**Extended Data Fig. 4 | Single-atom detection fidelity.** The raw images are analysed by first applying a low-pass filter followed by a simple peak-detection algorithm. A histogram of the amplitudes of all detected peaks in 2,000 images of a single spin component is plotted. We find a bimodal distribution. The maximum at low amplitudes originates from background noise of the camera. The second maximum at higher amplitudes is due to real photon clusters on the chip. Every peak with an amplitude above the threshold (vertical black line) is counted as an atom. This leads to single-atom detection fidelities of 97.8(9)%. There is a probability of 5.0(5)% for a false positive detection of an atom on each image for our chosen region of interest of 320 × 320 px. For 6 + 6 atoms this leads to a rate of false-to-true detections of 0.83(10)%. The solid red and blue lines are Gaussian and exponential fits to the data, respectively, and the dashed line is their sum.
Extended Data Fig. 5 | Collection of 28 single momentum-space projections taken at $E_B/h\omega_r = 1.97$. The images have been postselected for the correct particle number of the 6 + 6 ground state but are otherwise chosen randomly. The dashed circles indicate the Fermi momentum. Atoms pairs with $\Delta \phi < 30^\circ$ and $p_\uparrow$ and $p_\downarrow$ larger than $2/3p_F$ are highlighted. These are the particles that contribute to the pair-correlation peak at the Fermi surface (see Fig. 2i). We find considerably more such pairs in images taken at larger interaction strengths (compare Extended Data Fig. 6).
Extended Data Fig. 6 | Collection of 28 single momentum-space projections taken at $E_B/h\omega_r = 0$. The images have been postselected for the correct particle number of the $6 + 6$ ground state but are otherwise chosen randomly. The dashed circles indicate the Fermi momentum. All detected atom pairs with $\Delta \phi < 30^\circ$ and both $p_\uparrow$ and $p_\downarrow$ larger than $2/3p_F$ are highlighted. These are the particles that would contribute to the pair-correlation peak at the Fermi surface (see Fig. 2f). Without interactions we find no additional pairs other than what is expected already from the single-particle densities. Considerably more pairs are present in images taken at larger interaction strengths (compare Extended Data Fig. 5).
Extended Data Fig. 7 | Average momentum space distributions. a, b, The mean momentum-space distributions $n_{\uparrow}(p_{\uparrow})$ of a single spin component and averaged over 1,000 images for two different binding energies and 6 + 6 atoms are shown. The dashed circle indicates the Fermi momentum. The distributions are to a good approximation radially symmetric. With increasing binding energy, the average momentum increases and we find more particles outside the Fermi momentum. This agrees with the picture that increasing the attraction enables particles to overcome the single-particle gap and form first Cooper pairs that finally turn into tightly bound dimers. c, From the average momentum-space distributions of both spin components it is straightforward to calculate the total mean kinetic energy of the system per spin component. For 6 + 6 non-interaction particles, we find a value very close to the expected ground-state (GS) kinetic energy per spin component of $E_{\text{kin}}^{\text{gs}} = 7\hbar \omega_r$. The kinetic energy increases monotonously as the attraction strength increases. The error bars represent the standard error of the mean. d, The unnormalized correlator $\langle p_{\uparrow} p_{\downarrow} \rangle_{\downarrow}$ is shown for $E_B/\hbar\omega_r = 1.97$. It is defined as in equations (1), (2) but without the $\langle n_{\uparrow} n_{\downarrow} \rangle$ term. It shows that for strong enough binding energies the paired fraction becomes large enough that pair correlations are visible even without subtracting the single-particle density contributions.
Extended Data Fig. 8 | Alternative visualization of the density–density correlation function. The pair-correlation functions in relative (\(C_R^{(2)}\)) and centre-of-mass (\(C_{CM}^{(2)}\)) coordinates as a function of the interaction strength \(E_\alpha\) are shown in (a–e) and (f–j), respectively. The figure represents an alternative method of binning and visualizing the 4D correlation function \(C^{(2)}\), but is otherwise equivalent to the data shown in Fig. 2. The dashed circle indicates twice the Fermi momentum, \(2p_F\). In relative coordinates (a–e), we find a surplus of particles with momenta of \(|p_R| \approx 2p_F\), as expected for the formation of Cooper pairs with atoms located at the opposite ends of the Fermi surface. In the centre-of-mass frame (f–j), emergence of pairing is indicated by a sharp peak at zero momentum that increases in weight with the interaction strength. k–o, Radial integrals of the relative momentum correlation densities in a–e. The error bars represent the standard error of the mean.
Extended Data Fig. 9 | Correlations in a heated sample. We increase the energy of the sample by modulating the radial confinement with a pulse of 50 ms duration that is a square pulse of width 700 Hz in frequency space and with variable amplitude $A$ (inset). We find that the pair correlations reduce with increasing energy of the sample until they vanish completely. This measurement was taken at intermediate binding energies of $E_B/\hbar\omega_r = 0.6$. In the future, we plan to study above-ground-state physics of our mesoscopic Fermi gas in more detail. To this end, we have to develop a precise method to measure temperatures of the sample. The error bars are obtained from the counts in each bin of the correlation function and by assuming Poissonian statistics.