We report on the single atom and single site-resolved detection of the total density in a cold atom realization of the 2D Fermi-Hubbard model. Fluorescence imaging of doublons is achieved by splitting each lattice site into a double well, thereby separating atom pairs. Full density readout yields a direct measurement of the equation of state, including direct thermometry via the fluctuation-dissipation theorem. Site-resolved density correlations reveal the Pauli hole at low filling, and strong doublon-hole correlations near half filling. These are shown to account for the difference between local and non-local density fluctuations in the Mott insulator. Our technique enables the study of atom-resolved charge transport in the Fermi-Hubbard model, the site-resolved observation of molecules, and the creation of bilayer Fermi-Hubbard systems.

Understanding strongly correlated quantum systems poses a major challenge both for theory and experiment. Recent years have seen a significant progress in simulating quantum many-body physics with ultracold atoms [1–4]. In particular, the Fermi-Hubbard model plays a paradigmatic role in the study of strongly correlated fermions, most prominently for understanding high-$T_c$ superconductivity [5]. Quantum gas microscopes [4, 6, 7] of fermionic atoms [8–12] provide the ability to explore fermion correlations with single-atom, single-site resolution. Recent works have demonstrated the metal and Mott insulator crossover [13–16], studied spin and charge correlations [17–22], revealed magnetic polarons [23] and studied spin [24], charge [25] and heat transport [26]. However, most experiments employ fluorescence imaging directly on the lattice used for Hubbard physics. Light-assisted collisions then remove atom pairs residing on the same lattice site from the image [27, 28], leading to parity projection [6, 7] and in particular the appearance of doubly occupied sites (doublons) as holes. Such Fermi gas microscopes thus measure only the density of singly occupied sites (singlons), i.e. the local moment [19]. The full density can be obtained via absorption imaging [13] but without single site resolution, or by selectively imaging either singlons or doublons [29].

Revealing the microscopic correlations giving rise to macroscopic observables of the Fermi-Hubbard model requires single-shot measurements of the full density. As the prime example, the fluctuation-dissipation theorem [30] relates the compressibility to the global number fluctuations of the system via the temperature, requiring measurements of the total density sensitive to atomic shot noise [31–33]. The importance of non-local density fluctuations has been demonstrated [16, 34], but revealing their microscopic origin requires site-resolved density measurements.

Progress in fluorescence imaging of the total, spin-resolved density was made by imposing a superlattice with smaller spacing than the physics lattice [18], thereby spatially separating atom pairs in the 2D plane into distinct wells before imaging. This enabled the study of the interplay between the charge and spin sector [23] in small systems of $\sim 6 \times 6$ sites.

In this Letter, we introduce a bilayer Fermi gas microscope enabling full site-resolved density readout of large
of a 2D Fermi-Hubbard systems in a single fluorescence image. This directly yields the equation of state as pressure, compressibility and doublon density are obtained as a function of density. Site-resolved density correlations reveal the importance of non-local correlations, from the Pauli hole at low filling to strong doublon-hole correlations at half filling. The measured density fluctuation and compressibility directly yield a theory-independent thermometer via the fluctuation-dissipation theorem [30]. In the Mott insulator, we find strongly correlated nearest-neighbor doublon-hole pairs, required to compensate local density fluctuations to yield the near-vanishing compressibility.

To record the full density information, our setup consists of a bilayer optical lattice potential beneath a microscope objective, shown schematically in Fig. 1(a). In the experiment, a 2D Fermi-Hubbard gas is prepared in a single horizontal layer of a 3D optical lattice as reported in [8], with horizontal (vertical) lattice spacing of \( a = 541 \text{ nm} (3 \text{ µm}) \). For imaging, the depth of the horizontal lattices is increased to prevent tunneling in the 2D plane. Some lattice sites will contain doublons. To separate these atom pairs into two vertically separated wells we impose a vertical superlattice (purple arrow in Fig. 1(a)) with 532 nm spacing, created by retro-reflecting a 1064 nm laser beam off the flat surface of the hemispheric microscope objective. Driven by Feshbach enhanced repulsive interactions, two atoms originally in a single lattice site separate vertically into different wells (Fig. 1(b)). There is an energy offset between the two vertical wells caused by the optical setup, so all atoms in singly occupied sites transfer into the same layer. After the splitting process, Raman sideband cooling is performed as in [8] and emitted optical pumping photons are collected through the microscope objective. In contrast to previous work with bosons [35], the layer separation is within the depth of focus of the microscope, allowing atoms in both layers to be simultaneously imaged onto the same diffraction limited spot on the camera. This removes the necessity of mechanically moving the objective’s focus to two different locations, taking two images and referencing them correctly.

We now demonstrate that separated atoms continue to fluoresce without light-induced loss. By raising the harmonic trapping potential, we create a band insulator at the center of the cloud and perform the vertical separation of atom pairs before imaging. Fig. 1(c) shows a typical histogram of an image, with fluorescence counts from singly occupied sites clearly distinguishable from those for originally doubly occupied sites. The fluorescence obtained from atoms in each layer can be tuned via the intensity of Raman light [36]. A typical image is shown in Fig. 1(d). The tell-tale “wedding cake” structure of the central band insulator at high fluorescence, surrounded by the Mott insulating gas at lower fluorescence, is clearly observed. Singly and doubly occupied sites are clearly distinguished [36], leading to the digitized image in Fig. 1(e). After this first image, and within the same experimental run, we can recombine the atoms into a single well again by ramping down the superlattice, and then performing Raman imaging. The resulting image is shown in Fig. 1(f), with a dark central region in place of the band insulator in Fig. 1(d), reflecting light-assisted collisions ejecting overlapping atom pairs [27, 28].

We first explore the effects of interactions on the total density in the crossover from a metal to a Mott insulator. Fig. 2(a) shows examples of radially averaged density \( n \) (circles) and doublon density \( d \) (squares) profiles in a Fermi-Hubbard gas at \( U/t = 7.1(4) \) (blue), 11.8(5) (black), and 25.3(6) (red). (b) Measured normalized compressibility \( \kappa n^2 t \). (c) Local density fluctuations \( \langle \hat{n}^2 \rangle - \langle \hat{n} \rangle^2 \) (circles) and total atom number fluctuations per area \( \langle (N^2) - (N)^2 \rangle / \text{Area} \) in a 5×5 box (triangles). (d-f) Thermodynamic variables vs. density: (d) normalized pressure \( P/U \), (e) compressibility \( \kappa n^2 t \), and (f) doublon density \( d \). All lines show Monte Carlo predictions [37] for \( T/t = 1.4 \) (blue), \( T/t = 1.6 \) (black), and \( T/t = 2.25 \) (red) for the same \( U/t \) as the data.

**FIG. 2. Equation of State of the 2D Fermi-Hubbard model.** (a) Radially averaged profiles of total density (circles) and doublon density (squares) in a Fermi-Hubbard gas at \( U/t = 7.1(4) \) (blue), 11.8(5) (black), and 25.3(6) (red). (b) Measured normalized compressibility \( \kappa n^2 t \). (c) Local density fluctuations \( \langle \hat{n}^2 \rangle - \langle \hat{n} \rangle^2 \) (circles) and total atom number fluctuations per area \( \langle (N^2) - (N)^2 \rangle / \text{Area} \) in a 5×5 box (triangles). (d-f) Thermodynamic variables vs. density: (d) normalized pressure \( P/U \), (e) compressibility \( \kappa n^2 t \), and (f) doublon density \( d \). All lines show Monte Carlo predictions [37] for \( T/t = 1.4 \) (blue), \( T/t = 1.6 \) (black), and \( T/t = 2.25 \) (red) for the same \( U/t \) as the data.
Access to the total density directly yields a measurement of the equation of state of the Fermi-Hubbard model. The canonical equation of state relates pressure $P = P(n, T, U, t)$ to density, temperature $T$ and interaction parameters $U$ and $t$. However, one is free to replace e.g. temperature by any other thermodynamic variable like the doublon fraction, and e.g. $t$ by compressibility $\kappa$, thereby obtaining an equation of state of directly and locally observable quantities [38, 39]. From the variation of density with potential $n(V)$ one obtains the pressure

$$P(V) = \int_{-\infty}^{\mu} n(\mu') \, d\mu' = \int_{V}^{\infty} n(V') \, dV',$$

for which knowledge of the central chemical potential $\mu_0$ at $V=0$ is not necessary [34, 38, 41–43]. Together with $P$, one has the compressibility $\kappa n = n \, \partial n / \partial P|_{T}$, and the dimensionless doublon fraction $d$, all as a function of density $n$ (Fig. 2(d-f), respectively). Moreover, for the Hubbard model, due to particle-hole symmetry, the chemical potential is $\mu = U/2$ at $n=1$, which fixes $\mu_0$ and the chemical potential $\mu = \mu_0 - V$ throughout the cloud. For the strongest interactions it can be observed how the pressure needs to rise above $U$ before breakdown of the Mott insulator occurs and the density can grow above $n=1$. Further tell-tale signatures of the Mott insulator are observed in the vanishing of compressibility (Fig. 2(e)) and the reduction in doublon density (Fig. 2(f)) at $n=1$. Finally, the compressibility, together with the total density fluctuations in Fig. 2(c) directly yield the temperature $T$ via the fluctuation-dissipation theorem. To this end, in the following we will investigate density correlations.

The density correlations of a non-interacting Fermi gas are determined by Pauli exclusion, which forbids two identical fermions to share the same phase-space cell. At non-degenerate temperatures, the probability to find two like fermions near each other is suppressed for distances smaller than the thermal de Broglie wavelength

$$\lambda_{dB} \sim a/\sqrt{T},$$

As the phase space density $n \lambda_{dB}^2 / a^2 \gtrsim 1$, i.e. $T \lesssim n t$ the size of this Pauli exclusion hole saturates to the spacing $a/\sqrt{n}$ between identical fermions. In a two-state mixture of fermions and at low filling, repulsion between unlike spins further deepens the correlation hole between particles. These non-local anti-correlations have the effect of reducing the total atom number fluctuations in a given region. Any local upward density fluctuation will be partially compensated by a reduction in nearby density. In Fig. 2(c) we demonstrate that density fluctuations are reduced in a $5 \times 5$ site box (squares) compared to onsite fluctuations (circles), indicating the presence of non-local anti-correlations between fermions.

We now use the full site-resolved density read-out of our microscope to directly measure the correlation hole in an interacting Fermi-Hubbard lattice gas. The Pauli hole has been inferred from antibunching of the parity-projected density in previous work [19]. The connected density-density correlation $\langle \hat{n}_i \hat{n}_{i+\delta} \rangle_{C} = \langle \hat{n}_i \hat{n}_{i+\delta} \rangle - \langle \hat{n}_i \rangle \langle \hat{n}_{i+\delta} \rangle$ characterizes the non-trivial correlation of finding two particles a distance of $\delta$ lattice sites apart, beyond that for uncorrelated particles at the same density. Fig. 3(a) shows the spatial dependence of $\langle \hat{n}_i \hat{n}_{i+\delta} \rangle_{C}$ at various densities. Strong non-local anti-correlations are clearly visible. Fig. 3(b) reports the total, local, and non-local density fluctuations. Significant negative non-local correlations indicate a de Broglie wavelength which extends over multiple lattice sites, requiring $T \sim t$ [44, 45]. We note that non-local correlations were inferred but not directly measured in [16]. The magnitude of local and non-local fluctuations is maximal at $n \approx 0.5$, a direct consequence of strong on-site repulsion between unlike spins. This effectively reduces the available area for each species by half. On-site density fluctuations are thus equal to that of a single spin species in half the area, of density $n$ and binomial fluctuation $(\hat{n}^2) - \langle \hat{n} \rangle^2 \approx n(1-n)$, peaking at $n=0.5$. Pauli exclusion requires a corresponding anti-correlation in the area surrounding a given local fluctuation, so non-local fluctuations peak near the same filling.

The spatial Pauli hole is directly visualized through the density-density correlation function $g_{nn}(\delta) = \langle \hat{n}_i \hat{n}_{i+\delta} \rangle / n_i n_{i+\delta}$. Fig. 3(c) shows the measured $g_{nn}(\delta)$ for nearest-neighbor and next-nearest neighbor displacements $\delta$ versus $n^2$, which normalizes distance by the Fermi wavelength. The strong reduction of $g_{nn}(\delta)$ within one interparticle spacing (blue shaded region) represents the direct observation of the correlation hole

\begin{align*}
\text{FIG. 3. Measurement of non-local density correlations in the 2D Fermi-Hubbard model.} & \\
\text{(a) Connected density-density correlations at various densities at } U/t = 11.8(5). & \\
\text{(b) Density fluctuations } \sum_{i} \langle \hat{n}_i \hat{n}_{i+\delta} \rangle_{C} & \text{ (total, black circles), } \langle \hat{n}^2 \rangle - \langle \hat{n} \rangle^2 \text{ (local, red triangles), and } \sum_{\delta \neq 0} \langle \hat{n}_i \hat{n}_{i+\delta} \rangle_{C} \text{ (non-local, blue diamonds). } & \\
\text{(c) Density-density correlation function } g_{nn}(\delta) & \text{ for displacements (1,0) (blue circles) and (1,1) (black squares) vs. } n^2, \text{ and theory for a non-interacting single-component Fermi lattice gas for displacement (1,0) at } T/t = 0.73 \text{ (blue dashed line).} & \\
\end{align*}
due to Pauli exclusion of like spins, and repulsion of unlike spins. The $g_{nn}^{(2)}$ for a single, non-interacting fermionic species at the full density $n$ shows good agreement, highlighting again that strong inter-spin repulsion reduces the available area for a given spin species by half.

With access to both the measured microscopic density fluctuations (Fig. 3) and the macroscopic compressibility (Fig. 2), we are now in the position to probe the fundamental correspondence between fluctuations and response in thermal quantum systems [46]. The general density fluctuation-dissipation theorem

$$\kappa n^2 = \left. \frac{\partial n_i}{\partial \mu} \right|_T = \beta \sum_\delta \langle \hat{n}_i \hat{n}_{i+\delta}\rangle_C,$$

where $\beta = 1/k_B T$, relates directly measurable macroscopic and microscopic quantities without reference to any theoretical model [30]. Significantly, non-local density correlations will remain a sensitive thermometer down to $T=0$ for any compressible system, such as e.g. the metallic regions away from $n=1$ [32]. Moreover, by averaging over the system’s area, Eqn. (1) relates compressibility to the global atom number fluctuations: $\kappa n^2 = \beta((\bar{N}^2) - (\bar{N})^2)/\text{Area}$. We note however that number fluctuations per area in a small subsystem will - for metallic states - always be larger than the total number fluctuations per area, precisely due to non-local correlations extending beyond the boundaries of the subsystem.

This is the origin of the violation of the area law for entanglement entropy already present for non-interacting fermions [47–49].

Fig. 4 shows the total connected density-density correlation (black circles) versus the normalized compressibility $\kappa n^2 t$ for the data in Fig. 3. A linear fit results in a temperature of the cloud of $T/t = 0.73(3)$, consistent with a measurement of spin correlations under the same conditions [36]. For comparison, local fluctuations (red triangles) are consistently larger than total fluctuations, highlighting again the importance of negative non-local correlations, inferred in [16]. To demonstrate sensitivity as a thermometer, we heat the system by scattering of lattice photons and measure an increasing temperature (Fig. 4 inset). We have thus established a self-calibrated and sensitive thermometer for lattice fermions.

The fluctuation-dissipation theorem also provides insight into charge fluctuations in the Mott insulator at half-filling, at temperatures $T \ll U$, where the compressibility vanishes. In any system where either $T \to 0$ or $\kappa n^2 \to 0$, Eqn. (1) implies that local and non-local density fluctuations must cancel. For finite tunneling
t \sim T \ll U$, the system remains insulating, although the local operator $t$ acts as a perturbation that causes charge fluctuations over short distances [34]. The dominant contributions to $\langle \hat{n}_i \hat{n}_j \rangle_C = \langle \hat{d}_i \hat{d}_j \rangle_C + \langle \hat{h}_i \hat{h}_j \rangle_C - 2\langle \hat{d}_i \hat{h}_j \rangle_C$ are nearest neighbor doublon-hole fluctuations which occur with probability $\sim (t/U)^2$ [50]. Their existence has been inferred in [19] by observing bunching of holes after parity projection. For fermions, these nearest neighbor doublon-hole correlations signal spin singlet formation, as Pauli exclusion prevents tunneling for spin triplets.

Armed with full density read-out, in Fig. 5 we now directly detect these doublon-hole fluctuations. At our temperatures $T \ll U$, where thermal fluctuations are frozen out, doublon-hole fluctuations are purely quantum in origin. Fig. 5(a) shows an example raw image and the reconstructed charge distribution of a lattice gas containing a Mott insulator in the strong coupling (atomic) limit ($U/t = 18.8(5)$), featuring many isolated doublon-hole pairs near $n=1$. For weaker interactions $U/t = 7.1(4)$ (Fig. 5(b)), doublon-hole pairs are no longer clearly distinguished. In Fig. 5(c-e) we show the spatial dependence of the connected doublon-hole correlator $\langle \hat{d}_i \hat{h}_j \rangle_C$, the nearest neighbor correlator $\langle \hat{d}_i \hat{h}_{i+1} \rangle_C$, and the doublon-hole distribution function $g^{(2)}_{dh} = \langle \hat{d}_i \hat{h}_{i+1} \rangle / d_i h_{i+1}$ versus density at $U/t = 11.8(5)$, all of which demonstrate strongly enhanced local doublon-hole correlations near $n=1$.

In Fig. 5(f) we report the nearest neighbor doublon-hole pair density $\langle \hat{d}_i \hat{h}_{nn} \rangle = \sum_{j \in nn} \langle \hat{d}_i \hat{h}_j \rangle$ with respect to $(t/U)^2$. The linear relationship highlights the physical origin of doublon-hole pair correlations in a coherent, off-resonant tunneling process of amplitude $\sim t/U$. To demonstrate the strength of bunching, we obtain the conditional probability $P(h_{nn}|d_i) = \langle \hat{d}_i h_{nn} \rangle / d_i$ to find a hole next to a doublon in Fig. 5(g). As a comparison, we also show the conditional probability for a Poisson process at the same hole and doublon density $4d$ (blue shaded area). At small $t/U$, the conditional probability far exceeds random chance, showing that doublons and holes are tightly bound in a Mott insulator.

In conclusion, we demonstrate a robust method to measure the total site-resolved density in a cold-atom realization of the 2D Fermi-Hubbard model. We use this ability to directly detect non-local correlations, in particular the Pauli correlation hole at low filling and doublon-hole correlations in the Mott insulating region. Model-free thermometry is established via the fluctuation-dissipation theorem. Using a magnetic field gradient, we can also perform spin dependent splitting and thereby simultaneously observe both charge and spin [36]. Our superlattice geometry opens up the ability to study bilayer and even multilayer Fermi-Hubbard models, relevant for high-temperature superconductivity [51, 52].

**Note added:** After completion of our experimental work [53], a spin-resolved bilayer imaging technique was realized in [54].

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**SUPPLEMENTARY MATERIALS**

**A. Point spread function**

In order to measure the full density information of a single-layer Fermi-Hubbard gas, a bilayer quantum gas microscope is used to prevent light-assisted collisions. By using a directly retro-reflected vertical superlattice with wavelength 1064 nm, the distance between the two layers is significantly smaller than similar setups with bosons [33] and fermions [54], allowing us to image both layers simultaneously. The spacing between two superlattice layers (532 nm) is smaller than the wavelength of the emitted light (770 nm) during fluorescence imaging, hence atoms in either layer can be in focus at the same time. In Fig. S1(a), we show the intensity profile of isolated lattice sites with one (blue) or two (red) atoms averaged over 200 isolated sites. For comparison, the horizontal lattice spacing is depicted as a white bar in each image. The radially averaged intensity profile normalized to the peak intensity for each occupation is shown in Fig. S1(b). The point-spread-functions have comparable shape, which is also similar to previous work [8].

![Fig. S1.](https://example.com/supplementary-figure.png)
B. Differential fluorescence imaging

We are able to selectively tune the amount of fluorescence obtained from each layer. The horizontally polarized Raman light ($\sim 767$ nm) used in the imaging scheme [8] is reflected off the substrate at a shallow angle ($\sim 10.8$ degrees), forming an interference lattice with large spacing ($\sim 2 \mu$m) in the vertical direction (schematically shown in Fig. 1(a)). By changing the angle of the beam by less than 0.5 degrees, we are able to place the interference node at either layer used for imaging (532 nm separation), or between the two layers. For a large range of interference node placements, loss rates are not significantly affected. We use a 10 mm glassplate to repeatably translate the Raman beams by $\sim 5$ mm in the Fourier plane, so that within a single experimental run, it is possible to image each layer sequentially while the other layer is spatially independent and depends on the probability of loss converting doublons to singlons $P_{s\leftarrow d}$, doublons to holes $P_{h\leftarrow d}$, and singlons to holes $P_{h\leftarrow s}$. We assume $H_i$ depends on the local densities surrounding a given particle, $\vec{r}_i$, $\vec{s}_i$, and $\vec{d}_i$, and on the conditional probability of exchange with a neighboring particle denoted by $\alpha_{ab}$. Specifically, the probability of particle $a$ at site $i$ transitioning to particle $b$ through hopping is given by $\alpha_{ab} \vec{r}_i$. We assume symmetry $\alpha_{ab} = \alpha_{ba}$ and we include an additional term which converts a doublon-hole pair to a singlon-singlon pair, denoted by $\alpha_{split}$.

Therefore

$$\hat{L} = \begin{pmatrix} 0 & P_{h\leftarrow s} & P_{h\leftarrow d} \\ -P_{h\leftarrow s} & 0 & P_{s\leftarrow d} \\ 0 & 0 & -P_{h\leftarrow d} - P_{s\leftarrow d} \end{pmatrix},$$

and

$$\hat{H}_i = \begin{pmatrix} -\alpha_{hd} \vec{d}_i - \alpha_{hs} \vec{s}_i - \alpha_{split} \vec{d}_i \\ \alpha_{hs} \vec{r}_i + \alpha_{split} \vec{d}_i \\ \alpha_{hd} \vec{d}_i \end{pmatrix},$$

For each experiment, we observe the imaging transition probabilities between the first and second images. By averaging over many images, we measure $\hat{1} + \hat{L} + \hat{H}_i$, $\hat{r}_i$, $\hat{s}_i$, and $\hat{d}_i$. Because the lower triangular elements of $\hat{1} + \hat{L} + \hat{H}_i$ are solely due to hopping, a linear fit to the observed transition rates as a function of the local densities provides $\alpha_{hs}$, $\alpha_{hd}$, $\alpha_{sd}$, and $\alpha_{split}$. Next, we measure the upper triangular elements of $\hat{1} + \hat{L} + \hat{H}_i$, subtract the known local matrix elements of $\hat{H}_i$ at the measured local densities, and average the resulting transition elements to obtain $P_{s\leftarrow d}$, $P_{h\leftarrow d}$, and $P_{h\leftarrow s}$.

Once $\hat{L}$ is known, we use $(\hat{1} + \hat{L})^{-1}$ to infer the original atomic density distribution given the information in image 1. Specifically, we infer that the real den-
sities at a given lattice site \((\vec{r}_{i,0}, \vec{s}_{i,0}, \vec{d}_{i,0})^t\) with observed densities \((\vec{n}_{i,1}, \vec{s}_{i,1}, \vec{d}_{i,1})^t\) in the first image are \((\vec{n}_{i,0}, \vec{s}_{i,0}, \vec{d}_{i,0})^t = (1 + \hat{L})^{-1}(\vec{n}_{i,1}, \vec{s}_{i,1}, \vec{d}_{i,1})^t\). The total density is then \(\vec{n}_{i,0} = \vec{s}_{i,0} + 2\vec{d}_{i,0}\).

To infer loss corrected atom number variances such as \(n_i^s\), we use the formula \(n_i^s = \sum_a (1 + \hat{L})^{-1}_{aa} \vec{n}_{i,a,1} \vec{n}_{i,a,0}\). We do not average the square of the inferred density at each site in each image, because this would involve two applications of \((1 + \hat{L})^{-1}\), whereas only one probabilistic event has occurred. For correlators involving different sites such as \(d_{i,0}h_{j,0}\), we assume independent loss and correct each site independently:

\[
\frac{d_{i,0}h_{j,0}}{n_{i,0}h_{j,0}} = \frac{d_{i,0}h_{j,0} - \sum_a (1 + \hat{L})^{-1}_{dd} (1 + \hat{L})^{-1}_{hh} \vec{a}_{i,1,1} \vec{a}_{j,1,1}}{n_{i,0}h_{j,0}}.
\]

Lastly, when measuring atom number fluctuations in a box involving \((\sum n_{i,0})^2\), we correct onsite terms separately from cross terms between different sites.

Atoms are not lost during imaging with typical measured fidelities of \(1 - P_{h\rightarrow s} \gtrsim 96\%\) for singlon detection, and \(1 - P_{s\rightarrow d} - P_{h\rightarrow d} \gtrsim 90\%\) for doublon detection.

![Image](http://example.com/image.png)

**FIG. S3.** (a) Image of a single spin species that has been transferred to the second vertical layer. (b) Measured spin-spin correlator \(C_{\text{spin}}(1) = 4(\hat{s}_{z,i}\hat{s}_{z,i+1})_C\) vs. singlon density \(s = n - 2d\) for \(U/t = 11.8(5)\) (red circles), and Monte Carlo theory [37] for \(T/t = 0.65\) (black solid line).

### D. Combined spin and charge read-out

Here we demonstrate how to extend the bilayer microscopy technique to combined spin and charge read-out. The basic principle is to map spin information to spatial information before imaging by transferring all spins of a single species to a specific layer. Two subsequent images of each vertical layer using the differential fluorescence imaging technique in Fig. S2 can then provide the full spin and charge information.

During normal bilayer imaging, all singlons are transferred to a specific layer of the vertical superlattice before imaging, due to residual vertical optical potential gradients. To map each spin to separate layers, a compensating vertical optical potential gradient (as in Fig. S2) and additional vertical magnetic field gradient are applied before imposing the vertical superlattice. The subsequent adiabatic transfer from a vertical single well to a double well with spin-dependent bias transfers all singlons of a given spin species to different, known vertical layers before imaging. We perform an RF transfer to two hyperfine states with large, opposing magnetic moments before performing this splitting process. The presence of only one spin species in either layer is checked by imaging a specific layer and removing one spin species with resonant light, as in [19]. In Fig. S3(a), we show an image of a single spin species that has been transferred to the second vertical layer via the process described. Patches of strong anti-ferromagnetic spin correlations are visible by eye.

### E. Thermometry via spin correlations

In Fig. S3(b), we show the connected spin-spin correlator \(C_{\text{spin}}(1) = 4(\hat{s}_{z,i}\hat{s}_{z,i+1})_C\) vs. singlon density \(s = n - 2d\) for \(U/t = 11.8(5)\). The measured spin correlator is consistent with Monte Carlo theory (black solid line) at a temperature \(T/t = 0.65\).

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