

Experimental realization of a fermionic spin-momentum lattice

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We experimentally realize a spin-momentum lattice with a homogeneously trapped Fermi gas. The lattice is created via cyclically-rotated atom-laser couplings between three bare atomic spin states, and are such that they form a triangular lattice in a synthetic spin-momentum space. We demonstrate the lattice and explore its dynamics with spin- and momentum-resolved absorption imaging. This platform will provide new opportunities for synthetic spin systems and the engineering of topological bands. In particular, the use of three spin states in two spatial dimensions would allow the simulation of synthetic magnetic fields of high spatial uniformity, which would lead to ultra-narrow Chern bands that support robust fractional quantum Hall states.

Ultracold atoms in optical lattices have been established as an important tool for the quantum emulation of condensed matter models [1], especially those with topological features [2, 3]. The inherent tunability afforded by optical lattices provides access to a variety of parameter regimes, which has proved essential in the seminal realizations of topological phases in ultracold matter [4–7]. Since then, efforts to study topology in other systems have led to the exploration of synthetic dimensions [8, 9], which provide internal degrees of freedom beyond those afforded by the trapping geometry and have enabled a new generation of experiments [10].

Several approaches to synthetic dimensions have been experimentally realized. Real-space lattices augmented with spin-orbit coupling (SOC) connect spin “lattice” sites via momentum exchange, creating Hall cylinders pierced by magnetic flux in a synthetic position-spin space [11–18], or creating Hall ribbons in optical clock experiments [19–22]. Real-space lattices are not always needed; SOC itself can provide synthetic degrees of freedom, which can act as a potent generator of Berry curvature [23–28] or provide control parameters for Hamiltonian engineering [29, 30]. Synthetic lattices entirely in momentum-space [31, 32] have been realized, and, with carefully engineered hopping schemes, have proven topological [33–36]. Recently, a synthetic lattice of Rydberg states has been employed for the study of a Su–Schrieffer–Heeger model [37], and a synthetic dimension of trap states created with patterned light [38].

Lattices composed of spin and momentum states, or spin-momentum (SM) lattices, have been proposed [39] as a platform to exhibit topological features, with some schemes potentially realizing the Laughlin state of the fractional quantum Hall effect [40, 41]. As a step towards this, we realize a fermionic spin-momentum lattice using SOC and three atomic Zeeman spin states. Previous experiments using spin-momentum lattices utilized bosons in a single-dimension [42, 43], and used a real-space lattice with lattice-band pseudospins [44]. Recently, a two-spin bosonic SM lattice has been implemented in an optical cavity [45]. Here, by providing sufficient links between three spin sites, we build a lattice of fermions in

a 2D spin-momentum space, without a traditional scalar optical lattice. This platform increases the flexibility of the synthetic dimension approach. In particular, the use of three spin states in two spatial dimensions allows the simulation of synthetic magnetic fields of high spatial uniformity, which lead to ultra-narrow Chern bands that support robust fractional quantum Hall states [41, 46].

Implementation. The synthetic lattice is composed of three Zeeman spin states in the ${}^1S_0(F = 9/2)$ ground state of ${}^{87}\text{Sr}$, labeled $X \equiv |m_F = -9/2\rangle$, $Y \equiv |m_F = -7/2\rangle$, $Z \equiv |m_F = -5/2\rangle$. In a momentum-dependent manner, the spins are cyclically coupled by up to 9 Raman lasers intersecting at 120° . In the rotating-wave approximation, we describe the atom-laser coupling as

$$\hat{V} = \Omega_{mne} e^{i(\mathbf{k}_R \cdot \mathbf{r} + \varphi_i - \varphi_j)} |m\rangle \langle n| + \text{H.c.} \quad (1)$$

where $m \neq n$ runs over the states X, Y, Z , $|\mathbf{k}_R| = |\mathbf{k}_i - \mathbf{k}_j| = k_L \sin^2 \frac{\theta}{2}$ is the magnitude of the single-photon recoil wavevector with $i \neq j$ denoting the beams driving a particular m - n coupling, Ω_{mn} is the coupling strength, and $\theta = 120^\circ$ is the angle between any pair of beams. The single-photon recoil energy is $E_R = (\hbar k_R)^2/2m = \hbar \times 22.7$ kHz. The phase differences $\varphi_i - \varphi_j$ are set to zero in the experiment, but we note that setting nonzero phases is at the heart of the ultra-narrow-band optical flux lattice experiment [41].

The setup implementing the optical couplings in Eqn. (1) is shown in Fig. 1(a). Up to three running-wave triplets of beams are incident on a degenerate Fermi gas (DFG) spin-polarized mostly into state $|X\rangle$ with $T/T_F = 0.36(5)$, where T_F is the Fermi temperature. Each beam $\hat{\mathbf{k}}_i$ contains up to three frequencies $\omega_i, \omega'_i, \omega''_i$, such that the energy difference between any two frequencies $\omega_i'' - \omega_j''$ matches an energy difference in the X, Y, Z manifold. These beams provide a Raman coupling between states, as in Fig 1(b). The quantization axis is defined by a \hat{z} -oriented magnetic field $B \approx 9.3$ G, along which we align the linear polarization of a beam ω_{Lift} [15, 47, 48] providing a strong ac Stark shift that lifts the degeneracy of the states X, Y, Z . The coupling beam polarizations are linear and angled at $33(1)^\circ$ with respect

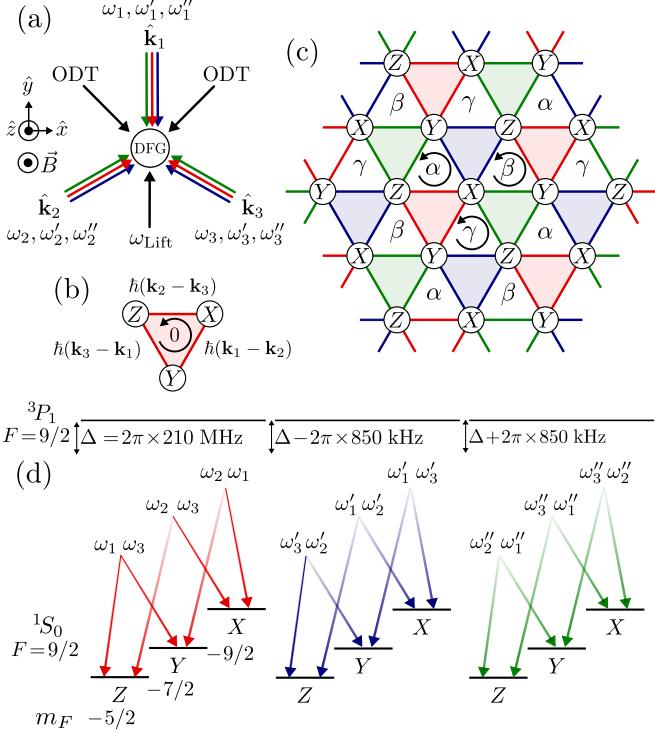


FIG. 1. Experiment schematic and coupling details. (a) A degenerate Fermi gas (DFG) in a crossed optical dipole trap (ODT) is exposed to cyclic Raman couplings between three internal spin states; see text for details. A beam ω_{Lift} provides a nonlinear energy splitting between the states. (b) States X, Y, Z are connected among themselves through spin-momentum exchange in units of $\hbar\delta\mathbf{k}_{ij} = \hbar(\mathbf{k}_i - \mathbf{k}_j)$. With a single set of beams, no net phase pickup is possible, denoted by the 0. (c) The couplings form a lattice in momentum space. Atoms encircling plaquettes labeled by α, β, γ can pick up a net phase. The link color indicates the frequency set in (d) to which the beams belong. (d) Details of the resonant couplings between the three internal states labeled X, Y , and Z , which represent the nuclear angular momentum projections $m_F = -9/2, -7/2, -5/2$, respectively, in the ${}^1S_0(F = 9/2)$ ground state. We circularly exchange the roles of the frequencies colored red (ω_i) in the blue (ω'_i) and green (ω''_i) coupling sets.

to the xy -plane, projecting approximately equal intensity among the possible Raman transition types π, σ^\pm . When using all nine frequencies, the beams form an infinite lattice in spin-momentum space, as in Fig. 1(c). Since the average starting atomic momentum $\langle p \rangle \ll 4\hbar k_R$, the fermions initially occupy only a small spread of states $|X/Y/Z, \mathbf{q} \approx 0\rangle$, where \mathbf{q} is the quasimomentum. After engaging the coupling beams, atoms with differing spin or \mathbf{q} occupy adjacent spin-momentum lattices, shifted by their corresponding quantum numbers. These lattices are not tight-binding, in the sense that in-situ motion is constrained only by the overall harmonic trap, allowing particles of different spin and \mathbf{q} to mix. Unlike a square lattice—which does not readily allow for nonzero Chern number—this triangular lattice naturally breaks inver-

sion symmetry and allows for magnetic flux. We note that nonzero gauge flux is possible only on the upward-pointing triangles, corresponding to momentum transfers involving all three frequency sets.

As shown in Fig. 1(d), the coupling beams utilize the dipole-forbidden transition ${}^1S_0(F = 9/2) \rightarrow {}^3P_1(F = 9/2)$, detuned below resonance by $\Delta/2\pi = 210$ MHz. The transition's narrow linewidth $\Gamma/2\pi = 7.4$ kHz allows coherent manipulation with minimal spontaneous emission, and no significant destructive interference [49, 50] arises from the THz-separated fine structure states 3P_0 and 3P_2 . In order to make each triplet unique, the upper-state detunings of the blue and green couplings are shifted by $\mp 37 E_R/\hbar$, larger than the $\approx 7.5 E_R/\hbar$ energy splittings. The role of the frequency ω_i in beam \mathbf{k}_i is circularly rotated amongst the three triplets, such that all frequencies resonantly couple all spin states.

In order to realize the lattice, careful attention must be paid to the energy levels of X, Y, Z , which are naturally degenerate. Since the ground states have $J = 0$ —rendering Zeeman shifts insignificant at our bias field—we use an ac Stark shift approach. Further leveraging the narrow-line intercombination transition, the lift beam ω_{Lift} is operated at 434.830043(5) THz, midway between the hyperfine resonance lines ${}^1S_0(F = 9/2) \rightarrow {}^3P_1(F = 7/2)$ and ${}^1S_0(F = 9/2) \rightarrow {}^3P_1(F = 9/2)$. This light was designed to produce a strong tensor shift $\epsilon = 1.76 E_R/\hbar$ of the X state, allowing each pair of Raman beams to uniquely couple two spin-momentum states. A necessary condition of the spin-momentum lattice model is that the coupling strengths $\Omega \ll \epsilon$. The coupling strengths here are $\Omega \approx 0.5 E_R/\hbar$ [51].

Experimental sequence. In our newly-built apparatus, we source ${}^{87}\text{Sr}$ from a commercial atomic oven from AOSense, which includes an integrated Zeeman slower and 2D MOT optics. After two MOT loading and cooling stages lasting 7 s [52, 53], the atoms are loaded into a crossed 1064 nm optical dipole trap (ODT) with an initial temperature of $\approx 2 \mu\text{K}$. The vertical (horizontal) trapping frequency is ramped up to 2.160(5) kHz (313(1), 397(2) Hz), at which point we spin-polarize the sample with a series of pulses resonant with the different m_F states via the ${}^1S_0(F = 9/2) \rightarrow {}^3P_1(F = 9/2)$ transition [51, 54, 55]. The ODT frequencies are then lowered back to ≈ 1 kHz vertically, and forced evaporation proceeds over the next 10 s, finally reaching a quantum-degenerate sample. Without spin-polarization, we routinely achieve $T/T_F = 0.20$, where T_F is the Fermi temperature, rising to $T/T_F = 0.36(5)$ when spin-polarized. Evaporation ends at mean geometric trap frequency $\bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3} = 71.4(1)$ Hz, yielding a 50 nK Fermi gas. Immediately following evaporation, the sample is spin-polarized in the state $X(80 \pm 7\%)$, and the ac Stark shifting beam ω_{Lift} is ramped on in 0.5 ms. Via optical Stern-Gerlach imaging [56], we verify that this timescale does not alter the spin polarization. We

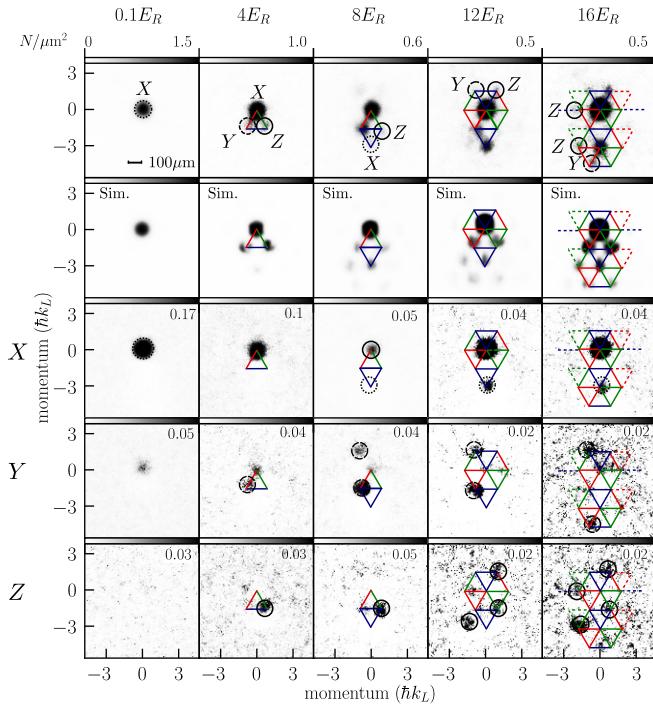


FIG. 2. Demonstration of the spin-momentum lattice. A DFG is exposed to the nine-beam coupling scheme, with the \mathbf{k}_1 triplet frequency-swept in order to populate more lattice sites; see text for details. The sweep progress is indicated in units of recoil energy E_R . The top row shows spin-unresolved momentum-space images, in which the lattice is most clearly evident. The second row shows the predicted dynamics from simulations (Sim.). Subsequent rows show spin-resolved images. In each row, we draw a link connecting each lattice site with a color corresponding to the involved beams, as detailed in Fig. 1(d). In the top row, we encircle and label each new lattice site as it becomes apparent during the sweep, and in the spin-resolved rows the expected populations are also encircled. As a guide to the eye, the links are also drawn across all rows. Each image is an average of ≈ 100 experimental runs taken at 12 ms time of flight. The peak atomic density is indicated in the top right of each image, in units of N atoms per μm^2 .

then introduce the coupling beams with a turn-on time of $< 1\mu\text{s}$.

We demonstrate the spin-momentum lattice in Fig. 2. In order to fill more sites, we emulate motion along a single dimension by subjecting the atoms to an inertial force [17] along $\hat{\mathbf{k}}_1$, ramping all three of that beam's frequencies at a rate $\hbar\partial_t(\omega_1, \omega'_1, \omega''_1) = 16.607 E_R/\text{ms}$ [57]. Hopping to neighboring sites is made favorable when the frequency difference between two coupling beams matches the energy and recoil shifts between states, providing enhanced state transfer between initial state $|X, \mathbf{q}\rangle$ and $|X, \mathbf{q} - \mathbf{K}\rangle, |Y, \mathbf{q} - \mathbf{K}\rangle, |Z, \mathbf{q} - \mathbf{K}\rangle$ for some reciprocal lattice vectors $\mathbf{K} = n_1\mathbf{q}_1 + n_2\mathbf{q}_2$ with integers n_1, n_2 . After a varying sweep time, all optical fields are quenched off, releasing the atoms from the harmonic trap. Atoms

that have tunneled to different lattice sites acquire a concomitant increase in momentum, in discrete units of the two-photon Raman momentum $\hbar k_R = \sqrt{3}/2\hbar k_L$. The lattice sites become spatially resolved after 12 ms time of flight, since the starting momentum distribution's full-width half-max width is $1.05(1)\hbar k_L$ and external heating by spontaneous emission from 3P_1 is minimal [51]. The atoms are then absorption-imaged in the xy plane using the ${}^1S_0 \rightarrow {}^1P_1$ transition at 461 nm ($\Gamma_{461}/2\pi = 30\text{MHz}$), which images all spins with approximately equal efficiency [58]. The bias magnetic field is kept on at all times, in order to maintain the spin quantization axis.

The individual columns of Fig. 2 demonstrate spin- and momentum-resolved imaging at various quench times. Sweep time is indicated by the final frequency deviation of the swept beam, in units of E_R . Intuitively, one would not expect stationary atoms ($\langle p \rangle \approx 0$) to tunnel before at least overcoming the recoil shift $4E_R$, and we observe this in the experiment. Denoting transferred momentum by $\hbar\delta k_{ij} = \hbar(|\mathbf{k}_i - \mathbf{k}_j|)$, and referring to beam triplets by their colors in Fig. 1(d), when the sweep reaches $4E_R$ we see beams ω_1, ω_2 from the red triplet driving the corresponding Raman transition $|X, 0\rangle \rightarrow |Y, \hbar\delta k_{12}\rangle$; similarly, the green beams ω'_1, ω'_3 allow $|X, 0\rangle \rightarrow |Z, \hbar\delta k_{13}\rangle$. By $8E_R$, atoms have firmly populated sites $|Y, \hbar\delta k_{12}\rangle, |Z, \hbar\delta k_{13}\rangle$, with initially-faster-moving atoms beginning to populate the site $|X, p = -3\hbar k_L\rangle$, completing a traversal of the first Brillouin zone. By $12E_R$, more atoms have tunneled through the Brillouin zone, and the momentum center-of-mass proceeds downward at $16E_R$; imaging becomes increasingly difficult due to the lower atom density, so we terminate here. As a consistency check, we also demonstrate spin-resolved imaging using spin blasts [51, 59] in order to verify that sites on the SM lattice are of the expected spin projection, m_F . We observe good consistency with the SM lattice model as drawn in Fig. 1(d), although mechanical effects of the spin-blasts can mask some lattice sites; notably, the X -site at $\mathbf{p} = -3\hbar k_L \hat{\mathbf{y}}$ (see Appendix). Our model shows qualitative agreement with the data for a scaled value of the measured Rabi coupling strengths. Some disagreement is evident, especially at the lattice sites with momenta $\mathbf{p} = \sqrt{3}/2\hbar k_L$, which are predicted to have a stronger amplitude than is observed. We attribute these mismatches to off-resonant effects not included in our effective Hamiltonian [51].

The lattice scheme presented here is readily tunable. Although the full spin-momentum lattice is composed of nine frequencies, we can remove links between lattice sites at will. We explore this flexibility in Fig. 3, where we show the driven dynamics experiment of Fig. 2, but now with all images taken at a common sweep time $12E_R$. In the two-beam scheme, composed of a single frequency in each of two beams $\hat{\mathbf{k}}_1, \hat{\mathbf{k}}_2$, we have reduced the system to a 1D SOC model between an effective spin up $|\uparrow\rangle = |X, \mathbf{q}\rangle$ and spin down $|\downarrow\rangle = |Y, \mathbf{q}\rangle$ [59, 60]. In the 3-

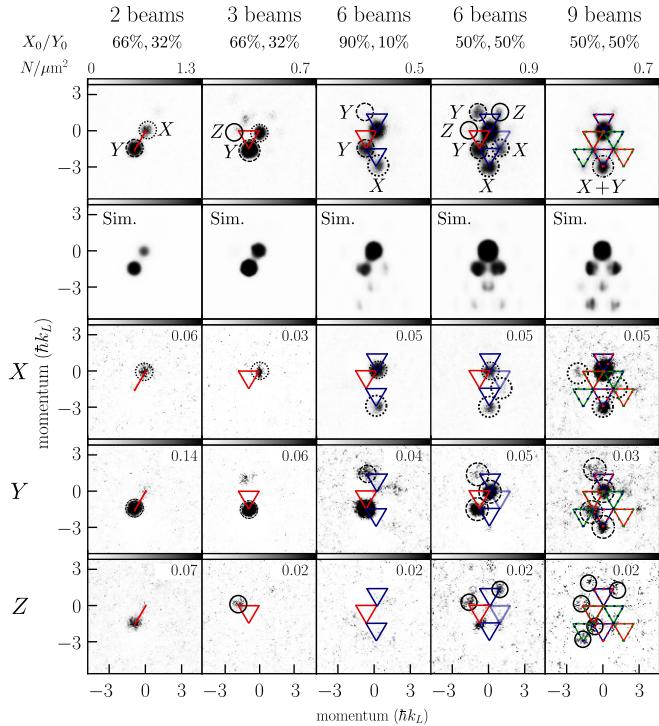


FIG. 3. Building the spin-momentum lattice; all images taken at a common sweep time of $12E_R$. The drawn links and circles are as described in Fig. 2, with the second row showing the predicted dynamics (Sim.). Each removed frequency corresponds to a missing link in the full lattice setup, allowing the exploration of quasi-1D SOC in the first column, to 2D SOC in the second column, culminating in the full lattice shown in the last column. The multiply-colored links in the last two columns indicate two SM lattices simultaneously overlaid: a two-state configuration, in which the experiment began with equal populations of X, Y .

beam case, with a single frequency in each of the $\hat{k}_1, \hat{k}_2, \hat{k}_3$ beams, we have a 2D spin-orbit coupling [25, 61, 62] cyclically linking the three states X, Y, Z . The 6-beam case consists of beams $\hat{k}_1, \hat{k}_2, \hat{k}_3$ each possessing two frequencies, labeled by their colors red and blue as labeled in Fig. 1. In the last two columns, we explore the dynamics starting from an even spin mix of states X and Y , which, in the 9-beam experiment, can be visualized as two SM lattices overlapped on $\mathbf{p} = 0$.

Conclusion and outlook. We have demonstrated a two-dimensional fermionic spin-momentum lattice without the use of standing waves. This adds to the wealth of cold atom synthetic dimension platforms available to study topological materials. The system's 15 ms lifetime exceeds our current experimental duration by a factor of 10, and could be further improved with increased Raman detuning Δ . The current lift beam strength imposes a 30 ms limit, which can be relaxed under appropriate conditions [51]. The number of visible lattice sites can be increased with larger Rabi coupling strengths, or by slow-

ing the sweep rate, which would couple more atoms out of the $p \approx 0$ momentum class. The spin-resolved imaging presented here could be improved by using stronger blast pulses to overcome the Doppler shifts among the lattice's numerous momentum states.

This work launches a novel platform for exploring topological physics with optical flux lattices. The natural extension of this work would be to load the atoms adiabatically into the lowest band and set nonzero coupling phases such that a gauge flux appears on the plaquettes labeled α, β, γ in Fig. 1. The topology of the band structure could then be probed using established anomalous velocity techniques [6, 17, 63], which involve accelerating the dressed atoms in the same manner as done here. Demonstrating this topology would enable the exploration of many-body fractional Hall states [41]. In the present experiment, spin-contact interactions are not expected to play a significant role; but these interactions may be increased through the use of a vertical real-space lattice, as detailed further in [41] for ^{87}Rb , which we note has the same scattering length $\approx 100a_0$ as ^{87}Sr .

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Appendix: Spin-resolved imaging

Atoms released from the SM lattice naturally spatially resolve according to their spin and momentum; however, in cases where the starting state is not pure or for verification purposes, it is necessary to have a means to resolve the spins. Directly imaging the spins using narrow-line absorption imaging [56] is impractical due to the low scattering cross-section and relatively low atom density of Fermi gases. Furthermore, since the atoms are moving after they are released from the SM lattice, the use of optical Stern-Gerlach separation is not desirable because it will (necessarily) disturb the momentum distribution.

We instead visualize the spin dynamics with blast pulses [59], propagating in the xy plane, resonant with the narrow $^1S_0(F = 9/2, m_F) \rightarrow ^3P_1(F = 9/2, m_F)$ transition at 689 nm. To be effective, the blast beam strength needs to be similar to the worst-case atomic Doppler shifts in the SM lattice $\Omega_{\text{Doppler}} \approx n \times 2\pi \times 4E_R = n \times 91$ kHz, where n is the number of photons absorbed in the SM lattice. Simultaneously, to avoid exciting neighboring m_F states, the blast strength must be smaller than the typical Zeeman splitting of the upper $^3P_1(F = 9/2)$ states, here $2\pi \times 790$ kHz.

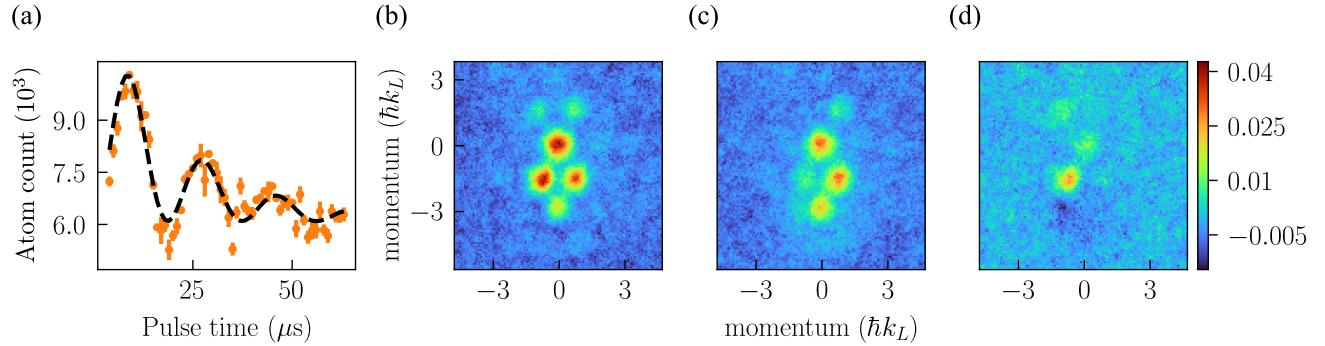


FIG. 4. Spin-resolved imaging process. (a) Spin-blast Rabi frequency calibrated by electron shelving of the $m_F = -7/2$ (Y) state. The fitted single-photon Rabi frequency is $\Omega = 335(1)$ kHz. (b) Averaged atom shot of the SM lattice at sweep time $8E_R$. (c) Spin-blast beam resonant with $m_F = -7/2$ applied to the SM lattice. (d) Subtraction of images (b) and (c), removing the common-mode background and revealing the locations of $m_F = -7/2$ atoms. A negative-OD region forms at the locations of the blast-scattered atoms.

The beams are calibrated as in Fig. 4(a) with electron-shelving [64]: atoms are excited into 3P_1 with a resonant 689 nm pulse for a time t , after which we apply a 4 μ s pulse of 461 nm MOT light to blow away all remaining ground-state atoms. Atoms so “shelved” in the 3P_1 state do not interact with the 461 nm light; instead, the fraction projected onto the ground state is measured with absorption imaging after a brief 2 ms time of flight. The overall decay curve matches the 3P_1 natural lifetime of 22 μ s. The SM lattice in Fig. 4(b) is subjected to this beam in Fig. 4(c), yielding the spin-resolved picture in Fig. 4(d) after subtracting the two images. Since the blast process relies on spontaneous emission, an atom can scatter (at most) a few photons before decaying to an adjacent m_F state, which is then transparent to the narrow-linewidth blast beam. This means the momentum acquired by the targeted atoms is comparable to the momentum of the atoms in the SM lattice, so some overlap is inevitable—as the negative optical density region in Fig. 4(d) shows. This effect can be particularly noticeable if the scattered atoms have strong geometric overlap with a region of interest, as seen in the X row of main text Fig. 2, particularly the $8E_R$ panel, which entirely masks the cloud at $\mathbf{p} = -3\hbar k_L \hat{y}$.

Furthermore, although the beam strengths are chosen to be comparable to the Doppler shifts, the efficiency of these single-tone blast pulses still decreases with increasing SM lattice photon absorption. To overcome this, a larger bias magnetic field would enable the use of stronger blast pulses. To overcome the geometric overlap issue, tilting the blast pulses out-of-plane would push targeted atoms out of the imaging focus region, more effectively removing them from the images.

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