Rotating atomic quantum gases with light-induced azimuthal gauge potentials and the observation of Hess-Fairbank effect

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We demonstrate synthetic azimuthal gauge potentials for Bose-Einstein condensates using atomic Raman coupling. The gauge potential is created by adiabatically loading the condensate into the lowest energy Raman-dressed state— a coreless vortex. Here one Raman beam carries orbital-angular-momentum. The azimuthal gauge potentials act as effective rotations and are tunable by the Raman coupling and detuning. We characterize the spin texture of the dressed states, which agree with the theory. The lowest energy dressed state is stable with the Raman beams on, and the half-atom-number lifetime is about 4.5 s. Finally, we employ the azimuthal gauge potential to demonstrate the Hess-Fairbank effect, i.e., we produce dressed atoms in the absolute ground state which have zero quasi-angular momentum when the synthetic magnetic flux is below a critical value. Above the critical flux, a transition into a polar-core vortex is observed. Both types of SO(3) vortices in the $|\vec{F}| = 1$ manifold are created, the coreless vortex and the polar-core vortex. We thus present a paradigm of creating topological excitations by tailoring atom-light interactions. The gauge field in the stationary Hamiltonian opens the door to investigate rotation properties of atomic superfluid under thermal equilibrium.

Ultracold atoms provide a tunable and flexible platform to realize new classes of quantum matter and topological states, and to reach intriguing regimes which are not accessible to solid state systems. One of the milestones is introducing synthetic gauge potentials for neutral atoms as if they were charged particles. A key experimental technique for the realization is optical Raman coupling between different internal spin states where photon momentum is transferred to the atoms as if the spin state changes. This leads to the spin-linear-momentum coupling, which is a type of “general spin-orbit-coupling” (SOC), referring to coupling between the atomic spin and the center-of-mass motion of the atoms. Another class of SOC where the atomic spin is coupled to the OAM of the atoms has been demonstrated using one Laguerre-Gaussian (LG) Raman beam carrying orbital-angular-momentum (OAM) of light: this is the spin-orbital-angular-momentum coupling (SOAMC)

In SOAMC systems, the atoms dressed by LG Raman beams experience azimuthal gauge potentials. The azimuthal dispersion at radial position $r$ is $(\ell - \ell_{\text{min}})^2/2mr^2$ (see Eq. [1]), where $\ell$ is the eigenvalue of the angular momentum $L_z$. The shifted single minimum at $\ell_{\text{min}} = rB(r)$ corresponds to an azimuthal gauge potential $\vec{A} = A(r)\vec{e}_\phi$ which is tunable by the Raman coupling strength $\Omega(r)$ and Raman detuning $\delta$ in our experiment. The azimuthal gauge potential is thus equivalent to an effective rotation in the stationary Hamiltonian. This gauge potential can be used to measure superfluid fractions from the spin population imbalance, where the spin imbalance is gauge-dependent and vanishes in the Landau gauge. Superfluid fraction corresponds to a nonclassical rotational inertia (NCRI), which is a property under thermal equilibrium, instead of metastable. In Refs. [20, 21], NCRI is manifested in the Hess-Fairbank effect: a cylindrical annulus filled with $^4$He rotates at a sufficiently low speed, after $^4$He is cooled below the superfluid transition temperature, the atoms stop rotating and become out-of-equilibrium with the container. NCRI is also an analog of the Meissner effect in superconductors.

Several attempts have been made to realize such light-induced azimuthal gauge potential. In Ref. [24], Bose-Einstein condensates (BEC) in a ring trap under LG Raman beams with the OAM transfer between spin states $(|m_F\rangle \rightarrow |m_F + 1\rangle)$ are studied. A lifetime of $\sim 0.1$ s of the lowest energy dressed state is reported. $F = 1$ SOAMC BECs is demonstrated in Ref. [14] with the gauge potential $A = 0$, given that the atoms are in the middle-energy dressed state in the $<\vec{F}> = 0$ polar phase. In this work, we present the first experimental realization of such azimuthal gauge potential with SOAMC BECs in the lowest energy dressed state, with a nonzero $A$.

SOAMC BECs can be implemented with an atom-light coupling $\vec{\Omega}_{\text{eff}} \cdot \vec{F}$ where the direction of $\vec{\Omega}_{\text{eff}}$ winds by a multiple of $2\pi$ as the azimuthal angle $\phi$ increases from 0 to $2\pi$. $\vec{\Omega}_{\text{eff}}$ is a light-induced effective magnetic field $\vec{B}$ and $\vec{F}$ is the atomic spin. This $\vec{\Omega}_{\text{eff}}$ is typically realized by using LG Raman beams. Versatile design of $\vec{\Omega}_{\text{eff}}$ allows creation of topological excitations in spinor BECs, where the rich variety of order parameters accommodate various types of topological defects. This is analogous to the works using spin rotation with real magnetic fields $\vec{B}$ with the...
Hamiltonian term $\vec{B} \cdot \vec{F}$ [27], where coreless vortices [23], monopoles [29][32], 2D [33][34] and 3D skyrmions [35] are created. These spin textures with spatially-dependent spin orientations give rise to synthetic gauge potentials and synthetic magnetic fields in the stationary Hamiltonian, where synthetic monopoles [29][32] and the geometric Hall effect [30] are demonstrated. Furthermore, creations of 2D skyrmions [37] and spin monopoles [38] with pulses of Raman LG beams are reported.

In this letter, we first load a BEC into the lowest energy Raman-dressed state, creating a coreless vortex. Here one Raman beam is LG and carries OAM. The dressed atoms experience azimuthal gauge potentials as effective rotations, which we exploit to demonstrate the Hess-Fairbank effect. As the synthetic magnetic flux arising from the effective rotation is below a critical value, the dressed atoms in the absolute ground state in the thermal equilibrium have zero quasi-angular momentum $\ell^\pm$ (see later texts and Eq. (4)). A polar-core vortex [23][39][40] with $\ell^\pm = \mp \hbar$ is also achieved. We demonstrate the capability to create both types of SO(3) vortices, the $Z_2 = 0$ coreless vortex and $Z_2 = 1$ polar-core vortex, in a unified and controlled scheme. This opens a path for creating topological excitations by tailoring atom-light interactions.

For atoms under sufficiently large $\tilde{\Omega}_{\text{eff}} \cdot \vec{F}$ such that the motional kinetic energy $-(\hbar^2/2m)\nabla^2$ is negligible, the energy eigenstates of the overall Hamiltonian are well approximated by the eigenstates $|\xi_n\rangle$ of $\tilde{\Omega}_{\text{eff}} \cdot \vec{F}$. Under this approximation, the atom’s spinor wave function locally follows the local dressed eigenstate $|\xi_n\rangle$. The Raman beams in our experiment give $\tilde{\Omega}_{\text{eff}} = \Omega(r) \cos \phi \mathbf{e}_x - \Omega(r) \sin \phi \mathbf{e}_y + \mathbf{e}_z$, being the quantization axis of $|\xi_n\rangle$. The state of dressed atoms is $\langle \vec{r}|\Psi\rangle = \varphi_n(\vec{r})|\xi_n(\vec{r})\rangle$, where $\varphi_n$ is the external part and $|\xi_n\rangle$ is the normalized spin part of the wave function. $|\varphi_n\rangle = \sqrt{n_a}$ where $n_a$ is the density. The effective Hamiltonian for atoms projected to $|\xi_n\rangle$, which governs the evolution of $\varphi_n$, is [14]

$$H_{\text{eff}}^{(n)} = -\frac{\hbar^2}{2m} \nabla^2 (r, z) + \frac{(L_z - r A_z)^2}{2mr^2} + V(r) + \varepsilon_n + W_n,$$ (1)

Here $L_z = -i\hbar \partial_\phi$ is the angular momentum operator for $\varphi_n$ with the eigenvalue $\ell$, and $A_z = A_z(r)e_z = (i\hbar/r)|\xi_n| \theta_e |\xi_n\rangle e_\phi$ is the azimuthal gauge potential. $V(r)$ is the spin-independent trap, $\varepsilon_n = n\sqrt{\Omega(r)^2 + \delta^2}$ is the eigenenergy of $\tilde{\Omega}_{\text{eff}} \cdot \vec{F}$, and $W_n \approx \hbar^2/2m r^2$ is the geometric scalar potential. We label the lowest, middle, and highest energy dressed state as $|\xi_{-1}\rangle, |\xi_0\rangle, |\xi_1\rangle$, respectively. The lowest energy dressed state is given by Euler rotations [11]

$$|\xi_{-1}\rangle = e^{i(\theta + \gamma)} \left( e^{i\frac{1}{2} - \cos \beta} \sin \frac{\beta}{\sqrt{\Omega}}, e^{-i\frac{1}{2} - \cos \beta} \right)^T,$$ (2)

where $\beta(r) = \tan^{-1}(\Omega(r)/\delta)$ is the polar angle of $\tilde{\Omega}_{\text{eff}}$, and $\theta + \gamma$ is the phase from a gauge transformation. By choosing $\theta + \gamma = 0$ for all $\delta$, it leads to

$$A_{-1} = (\hbar/r) \cos \beta,$$ (3)

where the angular momentum of $\varphi_{-1}$ is $\ell$ in this gauge, and $\ell, \ell \pm m \hbar$ is the mechanical angular momentum of the bare spin $|m_F = 0, \pm 1\rangle$ component of the state $\varphi_{-1}|\xi_{-1}\rangle$, respectively. In order to avoid a singularity at $r = 0$ in the synthetic magnetic field $\vec{B} = \nabla \times \vec{A}$ [42] as $A_{-1}(r \rightarrow 0) \rightarrow \pm \hbar/r$ for $\delta > < 0$, we use alternative gauges with $\theta + \gamma = \pm \phi$, leading to

$$A\pm_{-1} = \frac{\hbar}{r} (\cos \beta \mp 1), \quad \Phi_{\text{eff}}^{-}\pm = h[\cos \beta (r_0) \mp 1], \quad \ell^\pm = \ell \mp \hbar(4)$$

for $\delta > < 0$, where $\ell^\pm$ is the angular momentum of the external wave function $\varphi_{\pm 1}$ in these gauges. $\Phi_{\text{eff}}^{-}\pm$ is the magnetic flux enclosed by $r = r_0$. Note that $\ell - r A_{-1} = \ell^\pm - r A_{\pm 1}$ is gauge invariant.

Consider atoms in the lowest energy dressed state $|\xi_{-1}\rangle$ prepared by loading a ground state BEC in $|m_F = -1\rangle$ with a Thomas-Fermi (TF) wave function $\varphi_{\text{TF}} = \sqrt{\rho_{\text{TF}}}$, where $|\vec{r}|\Psi\rangle_{i=0} = \varphi_{\text{TF}} (0, 0, 1)^T$. The atoms are loaded to $|\xi_{-1}\rangle$ as

$$|\vec{r}|\Psi\rangle = \varphi_{\text{TF}} \left( e^{i2\phi \frac{1 - \cos \beta}{2}}, -e^{i\phi \sin \frac{\beta}{\sqrt{2}} \frac{1 + \cos \beta}{2}} \right)^T,$$ (5)

where the phase winding of the $m_F = -1$ component remains zero during loading, since the potential $V(r)$ is
FIG. 2: (a) Optical densities on the $xy$ plane of the lowest energy dressed state projected onto bare spin $|m_F\rangle$ states with various detuning $\delta$ after 24 ms TOF. (b) Experimental spin fractions (solid colored lines) versus the radial position are compared to the spin texture from $|\xi_{-1}\rangle$ (gray lines) and the TOF simulations from TDGPE (dashed colored lines). The green, orange and blue curve represents the $|1\rangle$, $|0\rangle$, $|\rangle$ components, respectively, where the shaded area indicates the standard deviation along $e_y$. The in-situ spin textures Eq. (5) after a magnification of $r \rightarrow r/13.0$ after TOF [41, 43] are shown as gray-dashed, -dotted, -solid lines for $|\rangle$, $|1\rangle$, $|0\rangle$ states.

cylindrically symmetric. Eq. 3 has $\ell = \hbar$ using the gauge in Eq. 3, as shown in our data in Fig. 2. Here, the atom’s spinor wave function follows $|\xi_{-1}\rangle$ only at $r \gg r_c$, where the kinetic energy is negligible and the loading speed $\dot{r}$ is sufficiently slow, achieving the adiabatic loading. Here $r_c$ is the adiabatic radius.

We perform 3D time-dependent-Gross-Pitaevskii equation (TDGPE) to simulate the loading of BEC from $|m_F = -1\rangle$ into the dressed state $|\xi_{-1}\rangle$. This includes the kinetic energies, quadratic Zeeman energy, mean field interaction parameters $c_0 = 4\pi\hbar^2(a_0 + 2a_2)/3m$ and $c_2 = 4\pi\hbar^2(a_2 - a_0)/3m < 0$, where $a_f$ is the s-wave scattering length in the total spin $f$ channel [22]. We also solve the absolute ground state using imaginary time propagation at given $\Omega(r)$ and $\delta$.

Our experiment begins with $N \approx 1.2 \times 10^5$ atoms in a $^{87}$Rb BEC in the $|F = 1, m_F = -1\rangle$ state in a crossed dipole trap. The TF radius of the condensate is $R \approx 5.9 \mu m$ along $e_x$. With a bias magnetic field $B_0$ along $e_x$, the atoms experience a linear Zeeman shift $\omega_Z/2\pi = 0.57$ MHz and a quadratic Zeeman shift $\hbar\omega_Z F_x^2$ with $\omega_Z/2\pi = 50$ Hz. A Gaussian Raman beam (G) and a Laguerre-Gaussian Raman beam (LG) with phase winding $m_r = 1$ co-propagate along $e_z$, coupling atoms in the $F = 1$ manifold and transferring OAM of $\Delta f = \hbar$ to the atoms (Fig. 1b). The two Raman laser beams have wavelengths $\lambda = 790$ nm with a frequency difference $\Delta\omega_L$ and a Raman detuning $\delta = \Delta\omega_L - \omega_Z$ (Fig. 1ab). We adiabatically load the atoms into the lowest energy dressed state $|\xi_{-1}\rangle$ with final Raman coupling strength $\Omega(r) = \Omega_M \sqrt{7(r/r_M)}e^{-r^2/2r_M^2}$. Here, $\Omega_M/2\pi = 3.0$ kHz and $r_M = 17 \mu m$ is the peak intensity radius.

We adiabatically load the BEC into $|\xi_{-1}\rangle$ by first turning on $\Omega(r)$ in 7 ms while holding the detuning at $\delta/2\pi = \delta_f/2\pi + 1.25$ kHz. Next, we sweep the detuning to the final value $\delta_f$ in 7 ms, where $\delta_f$ ranges between 500 Hz and -500 Hz. After preparing the atoms in the dressed state and holding the atoms for a time of $t_h$, we probe the atoms by switching off the Raman beams and dipole trap simultaneously. Then we adiabatically rotate the bias field to the direction of the imaging beam. The atoms of all $|m_F\rangle$ components expand together. We take absorption images after a 24 ms time-of-flight (TOF) along $e_y$ for each $m_F$ state using microwave spectroscopy for state-selective imaging. Here, single $|m_F\rangle$ images are taken in single experimental realizations. Images with different final detuning $\delta$ and $t_h = 1$ ms are shown in Fig. 2a. For $\delta > 0$, the $|m_F = -1\rangle$ component has no hole, and $|0\rangle$ carries a smaller hole than that of $|1\rangle$, consistent with that $|\rangle$, $|0\rangle$, $|1\rangle$ has angular momentum of $0, 2\hbar, 2\hbar$, respectively. The radial cross sections of the spin texture $D_{m_F}/(D_1 + D_0 + D_{-1})$ are shown (Fig. 2b), which average over the azimuthal angles. They are compared to the local dressed state $|\xi_{-1}\rangle$ in Eq. (5) and TOF simulations from TDGPE. Here, $D_{m_F}$ is the optical density of $|m_F\rangle$. With $\delta = -\delta_0 < 0$ and at large $r$ the spin fraction of $|\pm 1\rangle$ is about the same as that of $|\mp 1\rangle$ for $\delta = \delta_0$. The TDGPE simulation gives an $r_c \approx 1.8 \mu m$ at $\delta = 0$, corresponding to the spin texture agreeing with Eq. (5) at $r_{TOF} \gtrsim 23 \mu m$.

We further study the stability of the lowest energy dressed state with the Raman field on for a hold time $t_h$. With a deloading procedure [14] which maps the dressed states $|\xi_{-1}\rangle, |\xi_0\rangle, |\xi_1\rangle$ back to the bare spin states $|\rangle, |0\rangle, |+1\rangle$, we measure the population in $|\xi_n\rangle$ with $t_h$. We take images along $e_y$ after 14 ms-TOF with Stern-Gerlach gradient at a variable $\delta$ with a $t_h$ up to 7.0 s. Fig. 3a shows the atom number fraction in $|\xi_n\rangle$ over the
The curve (supplement) arising from the bias field noise in our setup. Eq. (3). At each
\(|\ell_g|\) versus detuning with \(t_h=1\) s and \(\Omega_M/2\pi = 3.0\) kHz. Inset in (a) shows the number fraction in \(|\xi_{-1}\rangle\) versus \(t_h\).

The total number \(|\xi_{-1}\rangle + |\xi_0\rangle + |\xi_1\rangle\) at \(\delta = 0\) versus \(t_h\). The atoms slowly populate the excited dressed states \(|\xi_0\rangle\) and \(|\xi_1\rangle\), and the initial rates increase as the peak Raman coupling is reduced from \(\Omega_M/2\pi = 3.0\) kHz to 1.5 kHz. The lifetime of atom number in \(|\xi_{-1}\rangle\) dropping to 50 % is 4.5 s (1.5 s) for \(\Omega_M/2\pi = 3.0\) (1.5) kHz. Fig. 3b displays the atom number fraction versus detuning at \(t_h = 1\) s and \(\Omega_M/2\pi = 3.0\) kHz, where the population rates decrease with \(\delta\). The results suggest the population rate into excited states increases with a reduction of energy gap \(\approx \sqrt{\Omega(r)^2 + \Delta^2}\), which can be explained by the effects of thermal atoms. Our estimated temperature \(T \geq 50\) nK\(= h \times 1\) kHz/\(k_B\) is compatible to the energy gap in the experiment.

Finally we demonstrate the Hess-Fairbank effect by studying the absolute ground state of dressed atoms in the gauge fields. We begin with thermal atoms right above the BEC transition temperature, load to the lowest energy dressed state with detuning \(\delta\), and evaporatively cool the atoms to achieve BECs. After holding a time \(t_{h1}\) for free evaporation, we probe the atoms after a 24 ms TOF. We note a related experiment [44], where normal bosonic atoms are evaporatively spun up and cooled to reach the condensation [77].

From GPE ground state simulations, it shows that the \(\ell\) of the ground state \(\ell_g\) is \(h (-h)\) at \(\delta/2\pi > 210\) Hz \((-210\) Hz), and \(\ell_g = 0\) at \(-210\) Hz\(< \delta/2\pi < 210\) Hz (Fig. 4). We obtain the averaged absolute value of winding number \(\langle|\ell_g|/h\rangle\) versus detuning in Fig. 4, by taking TOF images of the bare spin \(|m_F = 0\rangle\) component, whose mechanical angular momentum is \(\ell_g\) under the gauge in Eq. [3]. At each \(\delta\) we repeat the experiment for 20 times, observing \(|m_F = 0\rangle\) has either no hole or a hole, indicating \(\ell_g = 0\) or \(\ell_g = \pm h\), respectively. Such variations are due to the presence of a detuning noise of \(\approx h \times 70\) Hz rms (as we repeat the experiment with the same \(t_{h1}\), see supplement) arising from the bias field noise in our setup. The curve \(\langle|\ell_g|\rangle\) vs. \(\delta/2\pi\) is broadened and can become slightly asymmetric at \(\pm \delta\) since the detuning can slowly vary during which the data is taken. A simulation including a Gaussian-distributed detuning noise with rms value of \(h \times 70\) Hz is plotted in Fig. 4, with a center shift of \(-70\) Hz. The observed ground state with small \(\delta\) has \(\langle|\ell_g|/h\rangle \approx 0\), corresponding to \(\ell_g = 0\) if there were no detuning noise. This \(\ell_g = 0\) dressed state is a polar-core vortex; see the example in Fig. 4 at \(\delta/2\pi = 50\) Hz.

We compare our dressed atoms under gauge-induced effective rotations to BECs under mechanical rotation. Fig. 4d shows the flux \(|\Phi_{B mech}^+|\) approaches zero at large \(|\delta|\), and \(\Phi^+\) is along \((-+)e_z\) at \(\delta > (\approx) 0\). Thus, our system is analogous to the transition from zero to single-vortex ground state in rotating BECs as the following: \(\ell^\pm = 0 \rightarrow \ell^\pm = -h\) (i.e., \(\ell = h \rightarrow \ell = 0\)) under \(\Phi_{B mech}^- < 0\) with the gauge in Eq. [4]. And similarly for \(\ell^\pm = 0 \rightarrow \ell^\pm = h\) (i.e., \(\ell = -h \rightarrow \ell = 0\)) under \(\Phi_{B mech}^+ > 0\). At the transition of \(\delta/2\pi = \pm 210\) Hz, \(|\Phi_{B mech}^\pm| = 0.8h\) is much smaller than the critical effective flux \(\Phi_{B mech}^\pm\) for mechanically rotating BECs. The flux is \(\Phi_{B mech} = 2\pi n_m R^2 = h(\omega/\omega_c)\ln(R/r_v)\), where \(R\) is the BEC’s radius, \(\omega\) is the angular frequency, \(\omega_c = h/(2\pi m R^2)\ln(R/r_v)\) is the critical value for the ground state to possess single vortex, and \(r_v\) is the vortex core size. Thus the critical flux is \(\Phi_{B mech}/h \gg 1\), while the critical flux for our dressed state is \(|\Phi_{B mech}^\pm| = 0.8h\). This is due to the different form of \(A(r)e^{i\theta}\) between our dressed atoms and the mechanically rotating BECs (with symmetric gauge); see supplement.

In conclusion, we have demonstrated Raman-coupling-induced azimuthal gauge potentials. They act as effective rotations which we exploit to show Hess-Fairbank effect for the cold atoms. This work paves the way to probe atomic superfluid with the scheme of effective rotations, which can achieve a stationary Hamiltonian and thermal equilibrium. This scheme circumvents the issues from imperfect cylindrical symmetries in mechanical rotations.
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Note added: After we completed the manuscript, we noticed a recent work [46] on pseudo-spin 1/2 SOAMC BECs. They reported observations similar to our demonstration of Hess-Fairbank effect. Our work focuses on the prospective of effective rotations enabled by light-induced azimuthal gauge potentials and characterizing topological spin textures.

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[1] Y.-J. Lin, R. L. Compton, A. R. Perry, W. D. Phillips, J. V. Porto, and I. B. Spielman, Phys. Rev. Lett. 102, 130401 (2009).
[2] M. Aidelsburger, M. Atala, M. Lohse, J. T. Barreiro, B. Paredes, and I. Bloch, Phys. Rev. Lett. 111, 185301 (2013).
[3] H. Miyake, G. A. Siviloglou, C. J. Kennedy, W. C. Burton, and W. Ketterle, Phys. Rev. Lett. 111, 185302 (2013).
[4] J. Struck, C. Ölschläger, M. Weinberg, P. Hauke, J. Simonet, A. Eckardt, M. Lewenstein, K. Sengstock, and P. Windpassinger, Phys. Rev. Lett. 108, 225304 (2012).
[5] C. V. Parker, L.-C. Ha, and C. Chin, Nature Physics 9, 769 (2013).

[8] N. Goldman, G. Juzeliunas, P. Öhberg, and I. B. Spielman, Rep. Prog. Phys. 77, 126401 (2014).
[9] H. Zhai, Reports on Progress in Physics 78, 026001 (2015).
[10] Y.-J. Lin, K. Jimenez-Garcia, and I. B. Spielman, Nature 471, 83 (2011).

[11] Z. Wu, L. Zhang, W. Sun, X.-T. Xu, B.-Z. Wang, S.-C. Ji, Y. Deng, S. Chen, X.-J. Liu, and J.-W. Pan, Science 354, 83 (2016).

[12] L. Huang, Z. Meng, P. Wang, P. Peng, S.-L. Zhang, L. Chen, D. Li, Q. Zhou, and J. Zhang, Nature Physics 12, 540 (2016).

[13] G. Juzeliunas, P. Öhberg, J. Ruseckas, and A. Klein, Physical Review A 71, 053614 (2005).
[14] H. R. Chen, K. Y. Lin, P. K. Chen, N. C. Chiu, J. B. Wang, C. A. Chen, P. P. Huang, S. K. Yip, Y. Kawaguchi, and Y. J. Lin (2018), arXiv:1803.07860.
[15] C. Qu, K. Sun, and C. Zhang, Physical Review A 91, 053630 (2015).

[16] M. DeMarco and H. Pu, Physical Review A 91, 033630 (2015).
[17] Y.-X. Hu, C. Miniatura, and B. Grémaud, Physical Review A 92, 033615 (2015).

[18] L. Chen, H. Pu, and Y. Zhang, Physical Review A 93, 013629 (2016).

[19] N. R. Cooper and Z. Hadzibabic, Physical Review Letters 104, 030401 (2010).
[20] A. J. Leggett, Reviews of Modern Physics 71, S318 (1999).
[21] A. J. Leggett, Reviews of Modern Physics 73, 307 (2001).
[22] G. B. Hess and W. M. Fairbank, Physical Review Letters 19, 216 (1967).
[23] R. Ishiguro, O. Ishikawa, M. Yamashita, Y. Sasaki, K. Fukuda, M. Kubota, H. Ishimoto, R. E. Packard, T. Takagi, T. Ohmi, et al., Physical Review Letters 93, 125301 (2004).
[24] S. Moulder, Ph.D. thesis, University of Cambridge (2013).
[25] Y. Kawaguchi and M. Ueda, Physics Reports 520, 253 (2012).

[26] M. Ueda, Reports on Progress in Physics 77 (2014).
[27] T. Isoshima, M. Nakahara, T. Ohmi, and K. Machida, Physical Review A 61, 063610 (2000).
[28] A. E. Leanhardt, Y. Shin, D. Kielbinski, D. E. Pritchard, and W. Ketterle, Physical Review Letters 90, 140403 (2003).
[29] M. W. Ray, E. Ruokokoski, S. Kandel, M. Möttönen, and D. S. Hall, Nature 505, 657 (2014).
[30] M. W. Ray, E. Ruokokoski, K. Tiurev, M. Möttönen, and D. S. Hall, Science 348 (2015).
[31] D. S. Hall, M. W. Ray, K. Tiurev, E. Ruokokoski, A. H. Gheorghe, and M. Möttönen, Nature Physics 12, 478 (2016).
[32] T. Ollikainen, K. Tiurev, A. Blinova, W. Lee, D. Hall, and M. Möttönen, Physical Review X 7, 021023 (2017).
[33] J.-Y. Choi, W. J. Kwon, and Y.-i. Shin, Physical Review Letters 108, 035301 (2012).
[34] J.-Y. Choi, W. J. Kwon, M. Lee, H. Jeong, K. An, and Y.-i. Shin, New Journal of Physics 14, 053013 (2012).
[35] W. Lee, A. H. Gheorghe, K. Tiurev, T. Ollikainen, M. Möttönen, and D. S. Hall, Science Advances 4, eaa03820 (2018).
[36] J.-Y. Choi, S. Kang, S. W. Seo, W. J. Kwon, and Y.-i. Shin, Physical Review Letters 111, 245301 (2013).

[37] A. J. Leggett, Reviews of Modern Physics 88, 2595 (1996).
[38] A. J. Leggett, Reviews of Modern Physics 87, 210403 (2001).
[39] E. Lundh, C. J. Pethick, and H. Smith, Physical Review A 55, 2126 (1997).
[40] D. Zhang, T. Gao, P. Zou, L. Kong, R. Li, X. Shen, X.-L. Chen, S.-G. Peng, M. Zhan, H. Pu, et al. (2018), arXiv:1806.06263.
Supplemental Materials: Rotating atomic quantum gases with light-induced azimuthal gauge potentials and the observation of Hess-Fairbank effect

TDGPE SIMULATIONS

We numerically simulate the dynamics by solving the three-component 3D time-dependent-Gross-Pitaevskii equation (TDGPE). We use the Crank-Nicolson method and calculate in the system size of $(256)^3$ grid points with grid size 0.16 $\mu$m. During TOF, we solve the full 3D TDGPE for up to $\leq 4$ ms at which the interatomic interaction energy becomes less than 3 percent of the total energy. The further evolution is calculated by neglecting the interaction term. The results for the lowest energy dressed state with a short hold time $t_h = 1$ ms are shown in Fig. S1; our corresponding data is in Fig. 2. Using the loaded atomic state from TDGPE at $\delta = 0$, the probability of projection to the local dressed state $|\xi_{-1}\rangle$ is $\geq 0.99$ at $r \geq r_c = 1.8$ $\mu$m.

SYNTHETIC MAGNETIC FLUX OF DRESSED ATOMS

We compare the thermodynamic ground state of the dressed atoms in $|\xi_{-1}\rangle$ under synthetic magnetic fields to that of mechanically rotating BECs. Taking $\delta > 0$ where $\vec{B}^* \cdot \vec{e}_z < 0$ as the example, we consider the azimuthal kinetic energy in Eq. (1), which is gauge invariant and can be written as $(\ell^+ - rA_{-1}^+)^2/2mr^2$ in the gauge of Eq. (4). The contribution to the total kinetic energy is

$$\int dr 2\pi r (\ell^+ - rA_{-1}^+)^2 2mr^2$$

$$= \int dr 2\pi r (\ell^+)^2 2mr^2 - \int dr 2\pi r (\ell^+)^2 h m^2 r^2 (\cos \beta - 1) + \int dr 2\pi r (rA_{-1}^+)^2 2mr^2.$$  (S1)

At $\delta > 0$, $\cos \beta - 1$ is negative. We thus compare the energy with $\ell^+ = 0$ and with $\ell^+ = -h$, where the lower one is the ground state. The energy for $\ell^+ = -h$ relative to $\ell^+ = 0$ is

$$E_{B^*} = \int_{r_v}^R dr 2\pi r \frac{h^2}{2mr^2} + \int dr 2\pi r \frac{h^2}{mr^2} (\cos \beta - 1),$$  (S2)

where the first term is the vortex kinetic energy $E_v$ and $r_v$ is the vortex core size. At large $\delta$, $\cos \beta - 1 \approx 0$ and $E_{B^*} \approx E_v$, leading to the $\ell^+ = 0$ ground state. As $\delta$ decreases, $|\cos \beta - 1|$ increases, and the ground state makes a transition to $\ell^+ = -h$ when the absolute value of the second term in Eq. S(S2) equals $E_v$.

Then we analogously consider a mechanically rotating BEC with angular frequency $-\omega$, where $\omega > 0$. The kinetic energy for $\ell = -h$ relative to $\ell = 0$ is

$$E_{\text{mech}} = \int_{r_v}^R dr 2\pi r \frac{h^2}{2mr^2} - \int dr 2\pi r \frac{h^2}{mr^2} \frac{m\omega^2 r^2}{h}$$

$$= E_v - \int dr 2\pi r h \omega.$$  (S3)

The critical frequency for the ground state changes from $\ell = 0$ to $\ell = -h$ with increasing $\omega$ is $\omega_c = h/(2\pi m R^2) \ln(R/r_v)$.

To compare the dressed atoms and the mechanically rotating BEC, we consider the dimensionless gauge potential $rA/h$, which is equal to the flux enclosed by radius $r$ in unit of $h$ (see Eq. (4)). Thus, $\cos \beta(r) - 1$ and $m\omega r^2/h$ versus $r$ are plotted for these two cases in Fig. S2a, respectively. Here $\omega_c/2\pi = 8.0$ Hz for the TF radius $R \approx 6$ $\mu$m and assuming $r_v = 0.5$ $\mu$m. We show curves with various detuning $\delta/2\pi$ for the dressed atoms. We observe the flux enclosed by $r = R$ is $|\Phi_{\text{mech}}|/h \gg 1$ while $|\Phi_{\text{mech}}|/h < 1$ is limited by our $\Delta \ell = h$. Therefore we plot the integrand of the second term in Eq. S(S2) and Eq. S(S3), respectively, versus $r$ in Fig. S2b, which are contributed from the gauge potential. The integrand is scaled as $(\Phi/h)(R/r)$. At the critical detuning for the dressed atoms, the quasi-angular momentum $\ell^+ = \ell - h$ of the ground state changes from 0 to $-h$. In Fig. S2b, the enclosed area by the red curve denoting the critical detuning is the same as that of the red-dashed curve representing the mechanically rotating BEC with critical $\omega_c$, assuming the vortex energy $E_v$ is the same in the two cases. We find this condition is fulfilled.
FIG. S1: 3D TDGPE simulation results for the BEC loaded into the $\ell = \hbar$ lowest energy dressed state at various detuning $\delta$ with a hold time $t_h = 1$ ms. (a) Images of optical densities $D_{mF}$ on the $xy$ plane after a 24 ms TOF (b) Spin texture $D_{mF}/(D_1 + D_0 + D_{-1})$ vs. radial position. Dashed curves denote the in-situ profile after a $r \rightarrow r/13.0$ magnification; solid curves denote those after a 24 ms TOF. Green, orange and blue curves denote $|-1\rangle, |1\rangle, |0\rangle$, respectively. (c) total optical density $(D_1 + D_0 + D_{-1})$ for simulated TOF profiles.

FIG. S2: (a) Gauge potential $-rA/\hbar$, which is equal to the flux $-\Phi/\hbar$, versus radial position $r$. Black, green, red and blue solid lines denote the lowest energy dressed state with $\delta/2\pi = 1000, 500, 245, 125$ Hz, respectively. $\delta/2\pi = 245$ Hz is the critical detuning. Red dashed line represents the mechanically rotating BEC with critical frequency $\omega_c$. (b) $-(\Phi/\hbar)(R/r)$ versus $r$. Symbols represent the same as those in (a).
at the critical detuning $\delta/2\pi \approx 245$ Hz. Fig. S2a shows that at small $r$ the gauge $|rA_{-1}|$ of dressed atoms with $\delta/2\pi \leq 500$ Hz is larger than the $m\omega_t r^2$ for the mechanical rotation. The contribution at small $r$ is further enhanced by the $2\hbar/\sqrt{mr}$ factor in the integrand (see Fig. S2b). Therefore, the critical flux for the symmetric gauge is $|\Phi_{\text{mech}}|/\hbar \gg 1$ while $|\Phi_{B_r}|/\hbar \approx 0.8$ for our dressed atoms.

**EXPERIMENTAL SETUP AND PROCEDURES**

For the coreless vortex data in Fig. 2, we first produce a $^{87}$Rb BEC and then load it into the Raman-dressed state $|\xi_{-1}\rangle$. We achieve a BEC of $N \approx 1.2 \times 10^8$ atoms in a crossed dipole trap in $|F, m_F \rangle = |1, -1\rangle$ [S1]. The dipole trap contains two 1064 nm laser beams propagating along $e_x, e_y = (e_x \pm e_y)/\sqrt{2}$ with beam waists of $\sim 30 \mu m$. After the forced evaporation and a 1.5 s free evaporation, the dipole beam powers are ramped up in 0.2 s in order to reach the final trap frequencies of $(\omega_x', \omega_y', \omega_z)/2\pi = (140, 140, 190)$ Hz. The radial trap frequency is chosen such that it is sufficient for the atoms to sustain a Thomas-Fermi profile in the presence of the anti-trapping light shift potential near Raman resonance, $-\sqrt{\Omega(r)^2 + \delta^2}$.

After the BEC preparation we wait for the external trigger from the 60 Hz line, after which we apply feed-forward current signals into bias coils to cancel the field noise from 60 Hz harmonics (see later discussions). Then we load the BEC into the lowest energy dressed state $|\xi_{-1}\rangle$ with the following procedures. We ramp the detuning to $\delta/2\pi = \delta_f/2\pi + 1.25$ kHz while the Raman beams are off, ramp $\Omega(r, t)$ in 7 ms to the final value of $\Omega_M/2\pi = 3.0$ kHz, and then ramp the detuning to $\delta_f/2\pi$ between 500 Hz and -500 Hz in 7 ms (see Fig. S3), subsequently holding $\Omega_M$ and $\delta$ at constant for $t_h$. The Raman beams are at $\lambda = 790$ nm where their scalar light shifts from the D1 and D2 lines cancel. The Gaussian Raman beam has a waist of 200 $\mu m$, and the LG Raman beam produced by a vortex phase plate has a phase winding number $m_r = 1$ and radial index of 0. The Raman beams are linearly polarized along $e_x$ and $e_y$, respectively.

To show the Hess-Fairbank effect in Fig. 4, we start with thermal atoms right above the BEC transition temperature, load to the dressed state $|\xi_{-1}\rangle$ using the same sequence as that for data in Fig. 2. Next we force evaporatively cool the atoms to achieve BECs and hold 0.2 s for free evaporation, then increase the dipole beam powers in 0.4 s to reach the trap frequencies $(\omega_x', \omega_y', \omega_z)$, subsequently holding for $t_{h1}$ before TOF. To reduce the shot-to-shot field noise when TOF starts, we wait for the 60 Hz line trigger at 20 ms before the TOF. This adds to $t_{h1}$ by a varying time up to 16.7 ms (period of 60 Hz) as the experiment is repeated.

For projection measurements of the spinor state $|\xi_{-1}\rangle$, we abruptly turn off the dipole trap and Raman beams, simultaneously and adiabatically rotate the magnetic bias field from along $e_z$ to that along the imaging beam direction within 0.4 ms. The spinor wave function is then projected to the bare spin $m_F$ basis. For data in Fig. 2 and Fig. 4, the atoms then expand in free space with all $m_F$ components together for a TOF. To perform spin-selective imaging, we apply microwave spectroscopy for imaging each $|m_F\rangle$ [S1]. Each $|m_F\rangle$ image is obtained in an individual experimental
realization. After the $F = 1$ atoms are transferred to $F = 2$ by the microwave pulses, we apply a resonant absorption imaging pulse of $\sim 14$ $\mu$s with $\sigma^+$ polarization at the $|F = 2, m_F = 2\rangle \rightarrow |F' = 3, m_F = 3\rangle$ cycling transition. With our $I/I_{sat}$ parameters, we use the modified Beer-Lamberet law [S2] to derive correct optical densities. For data in Fig. 3, we apply a Stern-Gerlach field gradient to spatially separate individual $|m_F\rangle$ states during TOF, then a $F = 1 \rightarrow F' = 2$ repumping pulse is applied, after which the images of all $|m_F\rangle$ states are taken in a single shot.

We characterize the ambient field noise after the external trigger from the 60 Hz line. The noise is dominated by the 60 Hz line signal with a standard deviation ($\sigma$) of $\sim h \times 300$ Hz. After we apply feed-forward signals in the bias fields to cancel the dominating field noise at 60 Hz, the $1 - \sigma$ residual field noise can be reduced to $h \times 120$ Hz within 0.4 s after the 60 Hz trigger, in the best case. For longer time the 60 Hz line has a phase decoherence and the feed-forward cancelation does not work. For the data in Fig. 2 where the Raman loading takes 14 ms, we prepare the dressed state after the 60 Hz line trigger in order to reduce the shot-to-shot field variation with a fixed time after the trigger, and the measured $1 - \sigma$ field noise from repeated experimental shots is $\sim h \times 70$ Hz.

**DATA ANALYSIS**

For data in Fig. 2, we average over 4 to 5 images for each $\delta$ taken under identical conditions. We post-select images whose vortex positions in $|m_F = \pm 1\rangle$ with respect to the BEC center are $< 0.3$ $\mu$m (converted from TOF position to the in-situ position). We determine $\delta/2\pi$ from the rf-spectroscopy with an uncertainty of $\approx 20$ Hz. The uncertainty of the spin fraction displayed in Fig. 2b is $\sigma/\sqrt{N}$, where $\sigma$ is the standard deviation of pixels along $e_\phi$ at a fixed $r$, and $N = 4$ or 5 is the number of images.

For data in Fig. 4, the variations in the winding number of $|m_F = 0\rangle$ component within 20 identical experimental realizations are most likely due to the detuning noise. Thus, in the images at $\delta/2\pi = 50$ Hz with $\langle |\ell_g|/h \rangle \sim 0$, for $|m_F = 0\rangle$ component we post-select those with no hole corresponding to $\ell_g = 0$. For $|m_F = 1\rangle$ component we select those with smaller holes corresponding to $\ell_g + h = h$ and excluding those with larger holes corresponding to $\ell_g + h = 2h$. Similarly, we post-select those with $\ell_g - h = -h$ in $|m_F = -1\rangle$ with $\delta/2\pi = 50$ Hz, and further select those with $\ell_g = h$ for $|m_F = 0, \pm 1\rangle$ images with $\delta/2\pi = 400$ Hz based on the hole size in the images. For both $\delta/2\pi = 50, 400$ Hz, each $m_F$ state image is a single-shot image.

In the Hess-Fairbank effect experiment, we consider the effects of the dominating detuning noise at 60 Hz with an amplitude of $\approx h \times 420$ Hz. After BEC is reached, the sum of the hold time and the ramp time of dipole beam power is 0.8 s. After the 60 Hz line signal decoheres, the feed-forward cancelation does not function, and thus the 60 Hz detuning noise could drive the atoms slightly out of the equilibrium and out of the absolute ground state. Further reduction of bias field noise is needed in order to improve the current measurements.

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[S1] H. R. Chen, K. Y. Lin, P. K. Chen, N. C. Chiu, J. B. Wang, C. A. Chen, P. P. Huang, S. K. Yip, Y. Kawaguchi, and Y. J. Lin (2018), arXiv:1803.07860.

[S2] G. Reinaudi, T. Lahaye, Z. Wang, and D. Guéry-Odelin, Optics Letters 32, 3143 (2007).