

Transverse Polarization Gradient Entangling Gates for Trapped-Ion Quantum Computation

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The construction of entangling gates with individual addressing capability represents a crucial approach for implementing quantum computation in trapped-ion crystals. Conventional entangling gate schemes typically rely on laser beam wave vectors to couple the ions' spin and motional degrees of freedom. Here, we experimentally demonstrate an alternative method that employs a polarization gradient field generated by a tightly focused laser beam—an approach theoretically proposed as the Magnus effect for quantum logic gate design [Mazzanti *et al.*, *Phys. Rev. Res.* **5**, 033036 (2023)]. Using this technique, we perform Raman operations on hyperfine qubits encoded in $^{171}\text{Yb}^+$ ions, generating spin-dependent forces along axial motional modes in a linear trap. By utilizing an acousto-optic deflector to create arbitrary spot pairs for individual ion addressing in two-ion (four-ion) chains, we achieve Mølmer-Sørensen gates with Bell-state fidelities exceeding 98.7(1)% [and 97.2(4)%]. Further improvements in numerical aperture (NA) and laser power could reduce gate durations while enhancing fidelity by orders of magnitude. This method is compatible with—and can significantly simplify—optical tweezer gate proposals, where motional mode engineering enables scalable trapped-ion quantum computation. The technique can be readily extended to two-dimensional ion crystals, representing a key advancement toward large-scale trapped-ion quantum processors.

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Introduction—Trapped-ion systems represent a leading platform for quantum computation, offering long coherence times [1–3], high-fidelity gate operations [4–8], and reconfigurable qubit connectivity [9,10]. Architectures based on trapped-ion crystals have emerged as a promising approach for building quantum processors, with hundreds of ions successfully trapped and quantum simulations performed in two-dimensional crystals [11–13]. Beyond quantum simulation demonstrations, a key requirement for quantum computation in this architecture is the implementation of entangling gates with individual addressing (IA) capability. Optical IA enables not only site-resolved qubit control but also facilitates parallel quantum gate operations—essential for quantum error correction and large-scale quantum algorithms [14,15], making optically addressed entangling gates a critical component of this approach.

Although entangling gates have been demonstrated in one-dimensional chains of tens of ions [16,17] and two-dimensional arrays of four ions [18], scaling quantum

computation to hundreds of qubits in such systems remains challenging due to the difficulty of optimizing gates with large, dense motional modes in extended ion crystals. To address this challenge, optical tweezer-based trapped-ion quantum computation has been proposed as a promising approach [19–23]. However, combining optical tweezer (OT) gates with IA-based entangling gates presents significant experimental challenges, as it requires multiple high-NA objective lenses along different axes. This limitation occurs because current IA methods couple to the laser propagation direction, while OT gates primarily modify motional modes perpendicular to the propagation axis—making single-lens OT gate implementations currently unfeasible.

In this Letter, we experimentally demonstrate Mølmer-Sørensen (MS) gates with a different IA coupling method, which is based on a polarization gradient generated by a tightly focused optical beam. This method enables coupling to motional modes perpendicular to the laser propagation direction, providing flexible IA capabilities—including entangling gates with axial modes in linear traps—while maintaining the simplified optical layout advantages of requiring only a single beam and high numerical aperture (NA) objective. Using this technique, we perform carrier

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suppressed Raman operations on hyperfine qubits encoded in $^{171}\text{Yb}^+$ ions and implement spin-dependent forces for entangling gates. By dynamically steering the laser beam via an acousto-optic deflector, we realize IA based entangling gates on two-ion and four-ion chains, achieving entangling gates with Bell-state fidelities exceeding 98.7(1)% and 97.2(4)%, respectively.

Our method provides a promising path toward scalable trapped-ion quantum computation. Its compatibility with optical tweezer architectures and potential extension to two-dimensional ion arrays make it a versatile tool for implementing parallel, high-fidelity gate operations in fault-tolerant quantum computing schemes [24–26]. This scheme, which leverages the transverse gradient of the optical field, provides a more flexible and simplified optical configuration, and may also play an important role in QCCD architectures [27,28].

Principle and experiment—The interaction Hamiltonian between an optical field and a trapped ion qubit under the rotating wave approximation is [29]

$$H_{I,k} = \frac{\hbar}{2} \Omega_0 e^{i(\mathbf{k} \cdot \hat{\mathbf{x}} - \delta t + \phi)} \sigma_+ + \text{H.c.}, \quad (1)$$

where Ω_0 is the Rabi frequency, $\mathbf{k}$ the laser wave vector, $\hat{\mathbf{x}}$ the ion position operator (coupling to motional states), σ_+ the spin raising operator, and δ, ϕ the laser detuning and phase, respectively. For spin-motion coupling, $\mathbf{k}$ must be nonzero. In a Raman scheme driving hyperfine qubits with two beams, the effective wave vector $\mathbf{k} = \mathbf{k}_1 - \mathbf{k}_2$ requires $|\mathbf{k}_1| \approx |\mathbf{k}_2|$ with different propagation directions to ensure $\mathbf{k} \neq 0$.

For one-dimensional motion along x under Lamb-Dicke approximation,

$$H_{I,k} \approx \frac{\hbar}{2} \Omega_0 [1 + i\eta_k(a + a^\dagger)] \sigma_+ e^{-i\delta t + i\phi} + \text{H.c.}, \quad (2)$$

where $\eta_k = k_x x_{\text{zpf}}$ is the Lamb-Dicke parameter with $x_{\text{zpf}} = \sqrt{\hbar/(2m\omega)}$ the ground state wave packet size. This wave vector coupling inherently maintains a dominant carrier transition ($\eta_k < 1$), limiting sideband operation speeds due to off-resonant carrier excitation.

Constructing a position-dependent Rabi frequency $\Omega_0 \rightarrow \Omega(x)$ (in the laser spot's local coordinate system) enables more flexible spin-motion coupling. Notably, this approach allows spin-motion coupling even when $\mathbf{k} \approx 0$ (with copropagating Raman beams):

$$H_{I,\Omega} = \frac{\hbar}{2} \Omega(x_0 + \hat{x}) \sigma_+ e^{-i\delta t + i\phi} + \text{H.c.}, \quad (3)$$

where x_0 represents the ion's equilibrium position. When the motion amplitude is much smaller than $\Omega(x)$'s spatial profile, the Hamiltonian expands as

$$\begin{aligned} H_{I,\Omega} &\approx \frac{\hbar}{2} [\Omega(x_0) + \Omega'(x_0)\hat{x}] \sigma_+ e^{-i\delta t + i\phi} + \text{H.c.} \\ &= \frac{\hbar}{2} [\Omega(x_0) + \Omega_s(a + a^\dagger)] \sigma_+ e^{-i\delta t + i\phi} + \text{H.c.}, \end{aligned} \quad (4)$$

where $\Omega_s = \Omega'(x_0)x_\omega$ denotes the sideband Rabi frequency and $\Omega'(x_0)$ the Rabi gradient. The $\Omega(x)$ profile originates from tightly focused optical fields (individual addressing spots). Effective spin-motion coupling requires large Rabi gradients, while selecting $\Omega(x_0) \rightarrow 0$ suppresses carrier transitions. This carrier suppression, previously demonstrated in standing wave fields without individual addressing capability [30], now combines with IA for faster entanglement gates. The system's lateral motion coupling also enables an optical Magnus effect analogy [25], suggesting promising quantum logic gate implementations [26].

By engineering the structure of the light field entering the focusing lens—specifically, the spatial mode and polarization of the laser beam [31,32], the profile of $\Omega(x)$ can be controlled. One approach to generating this profile is by utilizing the polarization gradient of the tightly focused addressing spot. As shown in Fig. 1(b), when linearly polarized (x -direction) light passes through a high-NA lens, a longitudinal E_z component emerges at the focal plane. If we approximate the E_z component as resulting from combining two side paths through the lens in ray optics within the diffraction volume, the amplitude and phase of the longitudinal polarization clearly exhibit position dependence, creating a polarization gradient with elliptical components at the laser focus. This effect has been used in experiments such as high-fidelity manipulation of the motional states of neutral atoms [33] and the synthesis of molecules [34].

For a qubit encoded in the hyperfine clock transition of $^2S_{1/2}$ in $^{171}\text{Yb}^+$ with quantization axis magnetic field B along the y axis, the Raman laser configuration shown in Fig. 1(c) generates a Rabi profile [Fig. 1(d)] approximated by

$$\Omega(x) = \Omega_w \frac{x}{w_0} e^{-2x^2/w_0^2}, \quad (5)$$

where $w_0 \propto \lambda/\text{NA}$ is the beam waist. At spot center ($x = 0$), maximum Rabi frequencies occur at $x_m = \pm w_0/2$ with $\Omega_m = \pm \Omega_w/(2\sqrt{e})$, where $\Omega'(x_m) = 0$ decouples motion and spin. Positioning the ion at x_m via laser spot shifting enables single-qubit rotations.

Our experiments utilize a linear chain of $^{171}\text{Yb}^+$ ions confined in a segmented blade trap [35]. The qubit is encoded in the hyperfine clock states $|0\rangle \equiv |F=0, m_F=0\rangle$ and $|1\rangle \equiv |F=1, m_F=0\rangle$ of the $^2S_{1/2}$ manifold, forming a magnetic-field-insensitive qubit with a 12.64 GHz splitting. The quantization magnetic field (7.94 Gauss) aligns along the z axis. The axial center-of-mass (COM) phonon mode frequency measures 287 kHz via DC tickling. After electromagnetically induced transparency (EIT)

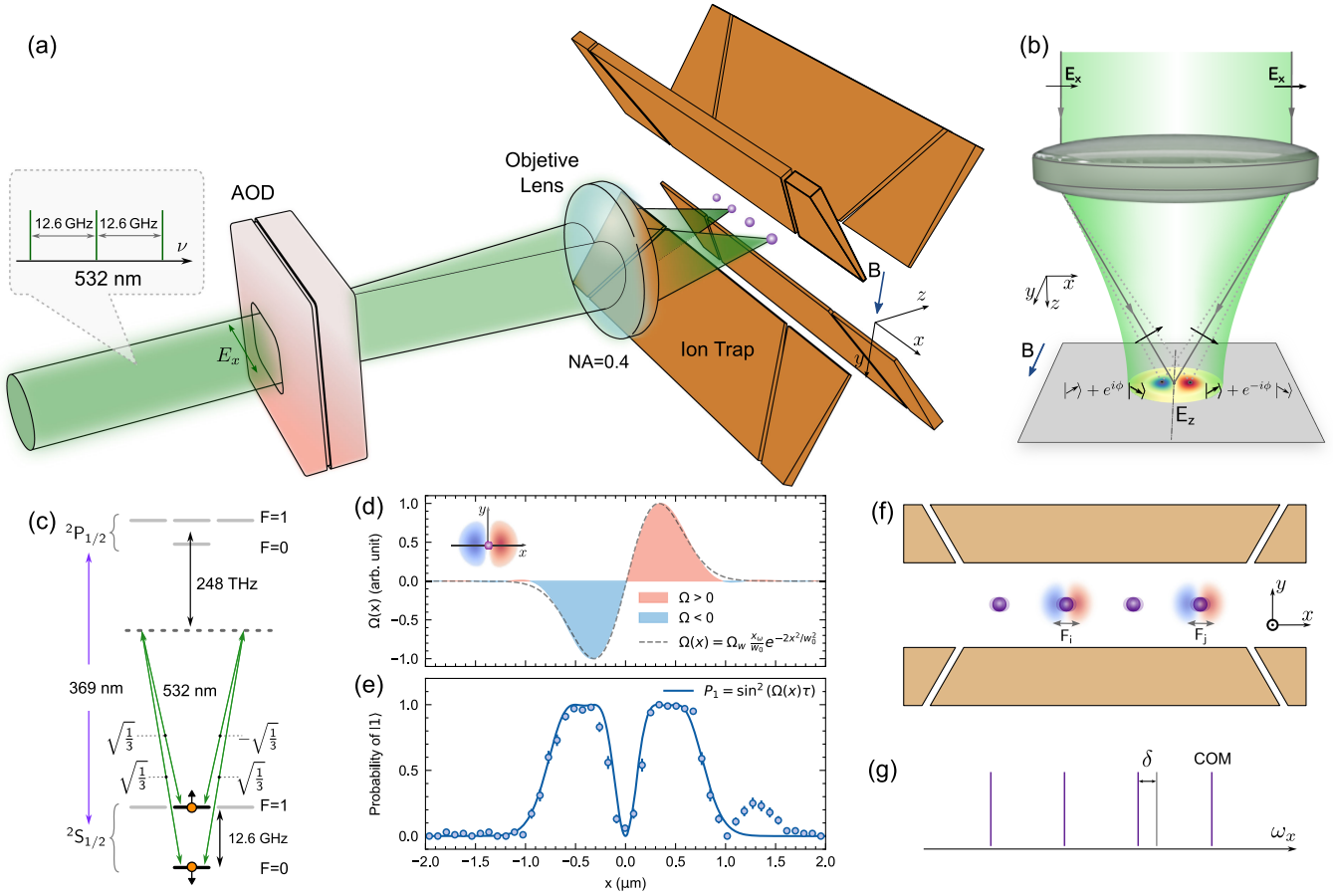


FIG. 1. Principle of polarization-gradient-based individual addressing. (a) An acousto-optic deflector (AOD) is employed in the addressing system to generate multiple laser spots after the objective lens with high numerical aperture (NA). A multitone laser is used to couple to the Raman transition. (b) x -polarized-focused Gaussian beam creates longitudinal E_z and elliptical polarization near focus. (c) For $^{171}\text{Yb}^+$, $\Omega(x) \propto \Omega_+^2(x) - \Omega_-^2(x)$, with Ω_+ and Ω_- corresponding to right- and left-circular polarization components, respectively. Although the $2P_{3/2}$ levels are not emphasized here, they contribute to $\Omega(x)$ through a similar mechanism, resulting in a reduction of the effective Rabi frequency. (d) Simulated Rabi profile (shaded) for NA = 0.4, fit to $\Omega_w(x/w_0)e^{-2x^2/w_0^2}$ (dashed), showing maximal gradient at center. (e) Spin-up probability measured by scanning the focus across the ion with $\tau = 7 \mu\text{s}$ pulses, from which the Rabi profile in (d) is extracted via fitting. The measured peak around $x \sim 1.2 \mu\text{m}$ is mainly caused by the coma aberration. (f) Multiqubit gates are implemented using aligned laser spots to address ions, coupling to an axial motional mode (g) to realize Mølmer-Sørensen gates.

cooling, the axial mode used for two qubit gate is cooled to below 1 phonon. Cooling, initialization, and state detection use 369 nm lasers as described in [36]. Fluorescence collected by a 0.64 NA objective lens and detected by a multichannel PMT achieves state-preparation-and-measurement (SPAM) fidelity $> 99.3\%$.

The experimental setup uses a multitone 532 nm Raman laser propagating along the z axis with linear polarization along the x axis, enabling the drive of axial motional modes. The Raman beam is generated through a two-step modulation process: first, a 1064 nm seed laser is phase modulated using an electro-optic modulator (EOM) to generate sidebands at 13.04 GHz. Subsequently, an unequal-arm Mach-Zehnder interferometer constructed with a pair of acousto-optic modulators (AOMs) driven at 200 MHz is used [37]. The $+1^{\text{st}}$ and -1^{st} order diffracted beams produce a Raman

beam with 12.64 GHz intensity modulation [41,42]. The spatial Rabi frequency variation under a 0.4 NA objective lens follows Eq. (5). Figure 1(e) shows state $|1\rangle$ population evolution under fixed-duration ($\tau = 7 \mu\text{s}$) Raman pulses versus beam position, with the fitted curve giving $\Omega_m = 2\pi \times 75.3 \text{ kHz}$ and $w_0 = 0.813(6) \mu\text{m}$. Ion motion introduces residual excitation ($P_1 = 0.05$) at beam center, corresponding to $\beta = \Omega_m/\Omega(0) \approx 15$ carrier suppression. This ratio should improve with lower ion temperatures and higher phonon frequencies.

To demonstrate the entangling gate, we employ a two-dimensional acousto-optic deflector (AOD) for beam steering, generating multisite addressing spots. Each addressing spot is aligned to its corresponding ion using the measurement method shown in Fig. 1(e), where scanning the beam and curve fitting determine each ion's precise position

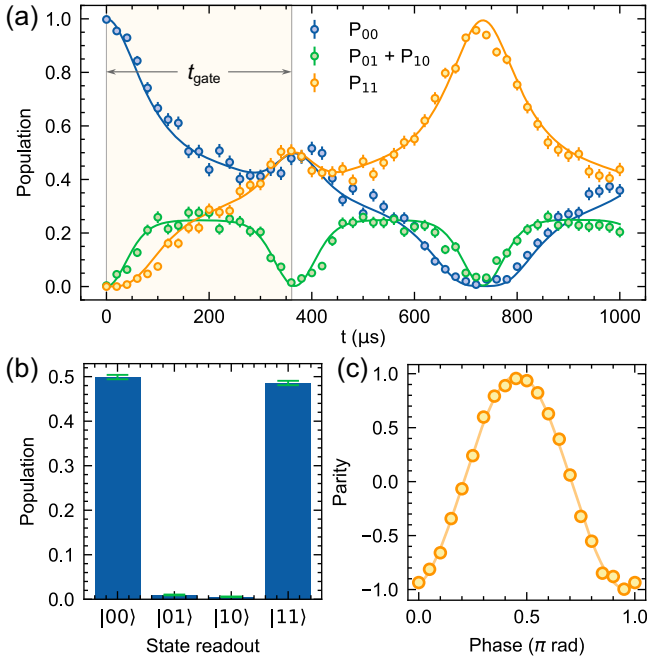


FIG. 2. Mølmer-Sørensen entangling gate in a two-ion chain. (a) Population evolution during gate operation, showing maximal entanglement at $t_{\text{gate}} = 367 \mu\text{s}$. Solid lines indicate theoretical predictions. (b) Bell-state population $P_{00} + P_{11} = 0.985(1)$, averaged over 10,200 repetitions. (c) Parity oscillation with 0.9887(1) contrast after applying $\pi/2$ pulses with variable phase ϕ to each ion. All data represent 400-repetition averages (unless noted), with error bars showing one standard error. State preparation and measurement (SPAM) errors are corrected [43].

relative to the spot center, achieving alignment accuracy below 100 nm [37]. As shown in Figs. 1(f) and 1(g), the configured polarization gradient addressing couples to the axial motional modes, enabling MS gate implementation. Leveraging the spin-motion coupling from Eq. (4), MS gates are implemented through selective addressing of ions i and j as

$$U = \exp\left(-i\frac{\pi}{4}\sigma_i^x \otimes \sigma_j^x\right). \quad (6)$$

First, we demonstrated our scheme on a two-ion chain. After cooling and initialization, a MS type gate is implemented by simultaneously applying red-detuned and blue-detuned Raman laser tones, coupling with the axial stretch mode. Figure 2(a) shows the evolution of spin populations as a function of interaction time. The fidelity of the gate is characterized through parity oscillations and population measurements. Parity is measured by applying a $\pi/2$ pulse to each ion after the MS gate, and is defined as $\text{Parity} = P_{00} + P_{11} - P_{10} - P_{01}$, where P_{ij} represents the population of the state $|ij\rangle$. The parity oscillates as a function of the phase of the $\pi/2$ pulse, and the contrast is defined as the oscillation amplitude. Figure 2(c)

TABLE I. Error budget for two qubit gate.

Error source	Simulated value
Mis-set detuning	1.0×10^{-2}
Motional dephasing	2.4×10^{-3}
Heating	6.4×10^{-4}
Photon scattering	2.8×10^{-4}
Rabi frequency fluctuation	2.5×10^{-4}
Carrier excitation	7.5×10^{-6}
Thermal errors	4.7×10^{-6}
Total	1.3×10^{-2}
Measured infidelity	$1.3(1) \times 10^{-2}$

demonstrates a contrast of 98.87(1)% and $P_{00} + P_{11} = 98.5(1)\%$, leading to a Bell-state fidelity calculated as $F = (P_{00} + P_{11} + \text{Contrast})/2 = 98.7(1)\%$. Additionally, for the initial state of $|01\rangle$ or $|10\rangle$, we also evaluated the fidelity of the antisymmetric Bell-state via measuring the populations, which is found to be $> 98.7(3)\%$ (see Supplemental Material [37]).

We have analyzed the sources of infidelity and presented them in Table I. In the current experiment, the dominant source of infidelity originates from calibration errors in the red and blue sideband frequencies. Due to the relatively low Rabi frequency, a detuning error of approximately 30 Hz can lead to an infidelity of 1%. The presence of the optical tweezer effects for changing the motional mode frequencies, and the differential AC Stark shifts for the clock states, causing energy levels changing along with the uncontrolled position drift between the laser focuses and ions, which limits more precise calibration, particularly for the sidebands frequencies. The second leading contribution comes

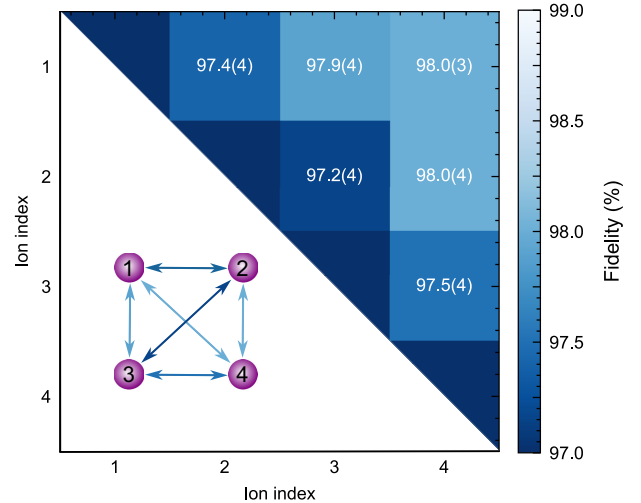


FIG. 3. Bell-state fidelities of the MS gates for arbitrary two ions in the four-ion chain. All reported fidelities are SPAM corrected. The population and parity data, as well as a typical SPAM error matrix, are provided in Supplemental Material [37].

from motional decoherence. For the stretch mode, with a blue sideband coherence time exceeding 100 ms, the calculated infidelity induced by motional decoherence is less than 0.25%. The implementation of far-detuned 532 nm Raman laser and carrier suppression mechanisms in the protocol effectively minimizes nonresonant effects and photon scattering contributions negligible. The gate fidelity can be further improved by simply reducing the gate time and employing modulated drive waveforms [44–48].

To demonstrate scalability, we implement an MS gate on a four-ion chain. As shown in Fig. 3, this implementation achieves a fidelity exceeding 97% for arbitrary pairs of ions. The decrease in fidelity is primarily due to the extended gate duration and crosstalk caused by intermodulation products in the AOD. The gate duration could be shortened by increasing laser power and using a larger NA objective. AOD-induced crosstalk can be further suppressed by adding compensation signals with opposite phases to counteract the intermodulation components [49].

Conclusion and outlook—Our method to construct entangling gates offers several advantages: (1) flexible individual control, enabling axial motional mode coupling as demonstrated; (2) carrier-suppressed coupling permitting large sideband Rabi frequencies for fast gates; (3) robustness against initial thermal motion, since the equivalent Lamb-Dicke parameter x_{zpf}/w_0 is significantly smaller than in standing-wave configurations; (4) simplified optical setup requiring only a single beam and high-NA objective; (5) enhanced Raman phase stability from shared optical paths for all frequency components. Importantly, these advantages enable our scheme’s extension to optical tweezer architectures with trapped ions, offering opportunities to overcome scalability bottlenecks in large-ion crystal quantum computation [19,23].

The quantum gate fidelity is limited by both the operation time and the motional-state coherence time. The Raman sideband Rabi frequency scales approximately as NA^4 , implying that larger numerical apertures can significantly accelerate gate operations. In practice, however, the achievable NA is constrained by the vacuum and optical systems. An experimentally feasible setup with $\text{NA} = 0.64$ could reduce the current gate time by roughly one order of magnitude [35]. Increasing the Raman laser power offers another route (the photon-scattering contribution to the infidelity remains at the level of 10^{-4} , as its rate increases linearly with intensity, while the gate time decreases inversely with intensity), but may induce charging of the vacuum window, leading to ion drift. This effect can be mitigated by applying an indium tin oxide (ITO) coating to the window surface.

Additionally, while our Letter employs tightly focused Gaussian beams for demonstrations, it is noteworthy that structured light modes (such as Laguerre-Gaussian modes and Hermite-Gaussian modes) offer further versatility: they not only enable flexible transverse spin-motion coupling and effective carrier suppression but also suppress the photon

scattering and differential light shifts between clock states, thereby enhancing overall fidelity and scalability [50–53].

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Data availability—The data that support the findings of this article are openly available [54].

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End Matter

The profile of Rabi frequency—To determine the relative position between the beam center and the ion, we fitted the beam profile using

$$\Omega(x) = \Omega_0 \frac{x}{w_0} e^{-x^2/w_0^2}. \quad (\text{A1})$$

Here, we provide a further explanation of the origin of this formula. When a linearly polarized Gaussian beam is tightly focused by a high-NA objective, the vectorial nature of light becomes significant. We adopt the Richards-Wolf vectorial diffraction formalism. In spherical coordinates, the focused electric field (in vacuum) components in the focal region ($z = 0$) can be written as

$$\begin{aligned} E_x(x, 0, 0) &\propto \int_0^{\theta_{\max}} \exp \left[-\left(\frac{\sin \theta}{\text{NA}} \right)^2 \right] \\ &\quad \times \sqrt{\cos \theta} \sin \theta (1 + \cos \theta) J_0(kx \sin \theta) d\theta \\ E_z(x, 0, 0) &\propto \int_0^{\theta_{\max}} \exp \left[-\left(\frac{\sin \theta}{\text{NA}} \right)^2 \right] \\ &\quad \times \sqrt{\cos \theta} \sin^2 \theta J_1(kx \sin \theta) d\theta, \end{aligned} \quad (\text{A2})$$

where $\theta_{\max} = \arcsin(\text{NA})$. These expressions give the primary electric field components in the $z = 0$ plane, near the focus, along the x direction. E_x and E_z together form a two-dimensional elliptical polarization vector $\mathbf{E}(x, 0, 0) = E_x(x)\hat{\mathbf{x}} + E_z(x)\hat{\mathbf{z}}$. With $\text{NA} < 0.5$, we can get an approximation as

$$\mathbf{E}(x) \approx E_0 \cdot e^{-x^2/w_0^2} \cdot \left(\hat{\mathbf{x}} + i\beta \cdot \frac{x}{w_0} \cdot \hat{\mathbf{z}} \right) \quad (\text{A3})$$

where w_0 represents the beam waist and $\beta \approx (\pi/4)\text{NA}$. For the elliptical polarization $\mathbf{E} = A\hat{\mathbf{x}} + iB\hat{\mathbf{z}}$, with $A = E_0 \cdot e^{-x^2/w_0^2}$, $B = A\beta \cdot (x/w_0)$, we can decompose it into circular polarization basis $\mathbf{E} = E_+\hat{\mathbf{R}} + E_-\hat{\mathbf{L}}$, with $E_+ = (1/\sqrt{2})(A + B)$ and $E_- = (1/\sqrt{2})(A - B)$.

For a qubit encoded in the hyperfine clock transition of $^2S_{1/2}$ in $^{171}\text{Yb}^+$, with a quantization axis magnetic field B along the y axis, the circular polarization basis coupling different transitions from the path in Fig. 1(b) in the main text. As the CG parameter of the transitions in form of $(1/\sqrt{3}, 1/\sqrt{3}, -1/\sqrt{3}, 1/\sqrt{3})$, the Rabi frequency of Raman laser is

$$\begin{aligned} \Omega(x) &= \Omega_+\Omega_+ - \Omega_-\Omega_- \propto E_+E_+ - E_-E_- \\ &\propto (A + B)^2 - (A - B)^2 = 4AB = I_0\beta \cdot \frac{x}{w_0} e^{-2x^2/w_0^2}, \end{aligned} \quad (\text{A4})$$

where $I_0 = E_0^2$ is the intensity of beam center. For a Gaussian beam with power P , the intensity is $I_0 = (2P/w_0^2)$. Since $w_0 \sim (\lambda/\text{NA})$, a rough relationship can be derived $\Omega(x) \propto \text{NA}^4$. Thus, the sideband Rabi frequency $\Omega_s \propto \Omega'(x_0) \cdot x_{\text{zpf}}$ is proportional to NA^4 , as mentioned in the main text.