

Generation of Motional Squeezed States for Neutral Atoms in Optical Tweezers

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Optical tweezers are a powerful tool for the precise positioning of a variety of small objects, including single neutral atoms. Once trapped, atoms can be cooled to the motional ground state of the tweezers. For a more advanced control of their spatial wave function, we report here a simple method to squeeze their motion, and the protocol to measure the squeezing factor based on momentum spreading estimation. We explore the limitations set by the technical imperfections of the tweezers, as well as the more fundamental limit set by their anharmonicity, and finally demonstrate a squeezing of 5.4 dB. The implementation of motional squeezing allows us to push back the limit set by the position quantum noise and thus to explore more extreme situations requiring atoms positioned with nanometric precision.

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Single neutral atoms trapped in optical tweezers [1–3] are widely used in numerous quantum technology applications, ranging from quantum sensing [4,5], quantum information networks with cavities [6,7], or waveguides [8,9]; to quantum simulation [10] and quantum computation [11–13]. Part of the experimental effort has been dedicated to reducing the residual motion of the atoms in the tweezers, which can be seen as a harmonic oscillator, for more accurate position control and higher fidelity of quantum operations. This is achieved by cooling the single atom to the motional ground state by Raman sideband cooling, as first reported in Refs. [14,15]. The atom is then described by a wave function with quantum uncertainties along the position x or momentum p_x quadratures saturating the Heisenberg inequality as $\Delta x^{\text{gs}} = \sqrt{\hbar/(2m\omega_r)}$ and $\Delta p_x^{\text{gs}} = \sqrt{\hbar m\omega_r/2}$, where ω_r is the angular frequency of the oscillator and m the mass of the atom. For ^{87}Rb atoms and typical trapping frequency of $\omega = 2\pi \times 100$ kHz, one obtains $\Delta x^{\text{gs}} = 24$ nm and $\Delta p_x^{\text{gs}} = m 0.015$ m.s⁻¹, which can be compared to a typical thermal uncertainty

$\Delta x^{\text{th}} = \sqrt{k_B T / (m\omega_r^2)} \approx 110$ nm for a temperature $T = 50$ μK , or to the typical size of the tweezers, the waist, on the order of 0.5–1 μm .

A natural way to increase the positioning precision, beyond the standard quantum limit described above, consists in squeezing the state: decreasing the uncertainty along one quadrature at the expense of increasing it on the other. We then define the squeezing factor β as $\Delta x^{\text{sq}} = \Delta x^{\text{gs}}/\beta$ and $\Delta p_x^{\text{sq}} = \Delta p_x^{\text{gs}} \times \beta$. Squeezing was first developed in the context of quantum optics [16–18], for metrological purposes, and is now a common tool widely implemented on a large diversity of quantum fields and platforms: motion of trapped ions [19,20], atoms in optical lattices [21], and levitating nanoparticles [22,23]; or electromagnetic waves in superconducting circuits [24,25] and in gravitational wave detectors [26,27]. As a few examples, displacement sensing through quantum amplification has been achieved by squeezing trapped ions by up to 20 dB ($\beta = 10$) [28], a 8 dB squeezed microwave field was produced in a superconducting cavity [29], and 8.5 dB squeezed light has been injected in gravitational wave interferometers [30].

In this Letter, we report a simple method to generate position-squeezed states for a collection of atoms individually trapped in optical tweezers. We first describe the generation and probing of squeezed states. We then show that optimizing the shape of the tweezers by holography [31] allows us to push the achievable squeezing level in an array, by compensating the tweezers radial anisotropy and

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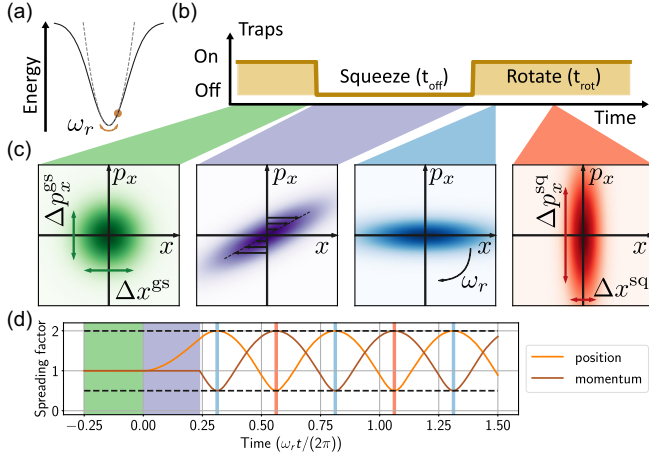


FIG. 1. Generation of squeezed states. (a) Optical tweezer potential (solid line), and its harmonic approximation (dashed line). (b) Experimental sequence and (c) sketches of the Wigner quasiprobability distribution of the atomic state in the x - p_x phase space through the sequence. (d) Theoretical standard deviations of x and p_x , normalized to the motional ground-state fluctuations, as a function of time. The atom is released from a harmonic trap during a time t_{off} , set here to $\omega_r t_{\text{off}} \approx 1.5$ [giving $\beta = 2$, see Eq. (3)].

minimizing the shape inhomogeneity over the ensemble of traps. Finally, we discuss the limit set to the squeezing factor by the inherent anharmonicity of optical tweezers.

Experimental platform—Our setup, presented in [31,32], is based on an array of ^{87}Rb atoms individually trapped in holographic optical tweezers. These microtraps are created by tightly focusing a 852 nm laser beam via an objective of numerical aperture $\text{NA} = 0.75$. The geometry of the optical tweezers array is controlled via a Spatial Light Modulator (SLM), impinging a versatile phase pattern (hologram) on the trap beam. Here, the chosen configuration is a 6×12 array with typical distance of $5 \mu\text{m}$ between atoms. The atoms are loaded from a magneto-optical trap (MOT), at a temperature of $50 \mu\text{K}$, into the optical tweezers (tweezers depth $U_0/k_B = 0.5 \text{ mK}$), thus occupying the bottom of the trap [see Fig. 1(a)]. In a first approximation, we write the trapping potential as harmonic and radially symmetric:

$$U(x, y, z) = -U_0 + \frac{1}{2}m\omega_r^2(x^2 + y^2) + \frac{1}{2}m\omega_z^2z^2, \quad (1)$$

where the radial (motion along x and y) and the longitudinal (motion along z) trapping angular frequencies are given by

$$\omega_r = \sqrt{\frac{4U_0}{mw_0^2}}, \quad \omega_z = \sqrt{\frac{2U_0}{mz_R^2}}, \quad (2)$$

with w_0 and z_R , respectively, the waist and the Rayleigh length of the tweezers. The typical trapping frequencies measured on our setup are $\omega_r \approx 2\pi \times 105 \text{ kHz}$ and

$\omega_z \approx 2\pi \times 25 \text{ kHz}$. Atoms loaded from the MOT are initially occupying high vibrational levels $\langle n_r \rangle \sim 10$.

Squeezing protocol—First, the atoms are cooled down to the motional ground state via Raman sideband cooling [14,15]. We observe a preparation efficiency of the radial motional ground state of more than 90%, more precisely $\langle n_{x,y} \rangle \approx 0.05$ and $\langle n_z \rangle \approx 0.3$ [32]. The squeezing protocol is then as follows [see Fig. 1(b)]: we switch off the traps for a duration t_{off} , letting the atom wave functions expand, and switch the traps on again to make the atoms wobble inside them during a time t_{rot} . The phase-space representations of the motional state of the atom along one of the two radial directions (say x) is displayed on Fig. 1(c) and illustrates how the state is squeezed. At first, it is a circular distribution with quantum uncertainties at the standard quantum limit (green). Then, during the free-fly step (purple), the distribution is sheared along the horizontal axis, as points higher on the p_x axis fly faster along the x axis. Third, the traps are switched back on, making the distribution rotate with angular frequency ω_r between a momentum-squeezed state (blue) and a position-squeezed state (red). Consequently, the uncertainties along x and p_x , normalized by Δx^{gs} and Δp_x^{gs} , show out-of-phase oscillations between a maximal value β and a minimal value $1/\beta$ [black horizontal dashed lines in Fig. 1(d)], at twice the trapping frequency.

The simple method described here allows for the preparation of a motional squeezed state. It was first demonstrated for neutral atoms in optical lattices in Ref. [21], where the expression of β was derived:

$$\beta^2 = 1 + \frac{(\omega_r t_{\text{off}})^2}{2} + \omega_r t_{\text{off}} \sqrt{1 + \frac{(\omega_r t_{\text{off}})^2}{4}}, \quad (3)$$

with the long-time limit, $\beta \approx \omega_r t_{\text{off}}$, being simply the ratio of the spatial expansion of the distribution due to its speed $\Delta p_x^{\text{gs}} t_{\text{off}}/m$ to its initial spread Δx^{gs} .

Switching off the traps is equivalent to adding a squeezing term $-\hat{x}^2 \propto (\hat{a} + \hat{a}^\dagger)^2$, where $\hat{a}$ ($\hat{a}^\dagger$) are the canonical lowering (raising) operators. Squeezing can also be obtained by modulating the trap depth (instead of switching it off fully) at twice the trapping frequency [28], or by performing a two-photon Raman drive [20]. The advantage of the on-off approach chosen here is its simplicity: in the case of optical tweezers, it is possible to switch off in 10 ns (3 orders of magnitude faster than the trap period). It also squeezes much faster, thus minimizing the effect of imperfections from the traps (see later).

Observation of squeezed states—Squeezed states can be characterized either in the n basis, or in the x - p_x basis. For the first case, the populations and relative phases of the different vibrational levels are measured, as done in Ref. [20] by probing a motion-sensitive electronic transition. The other way consists of reconstructing the phase-space distribution. A direct method applicable for the

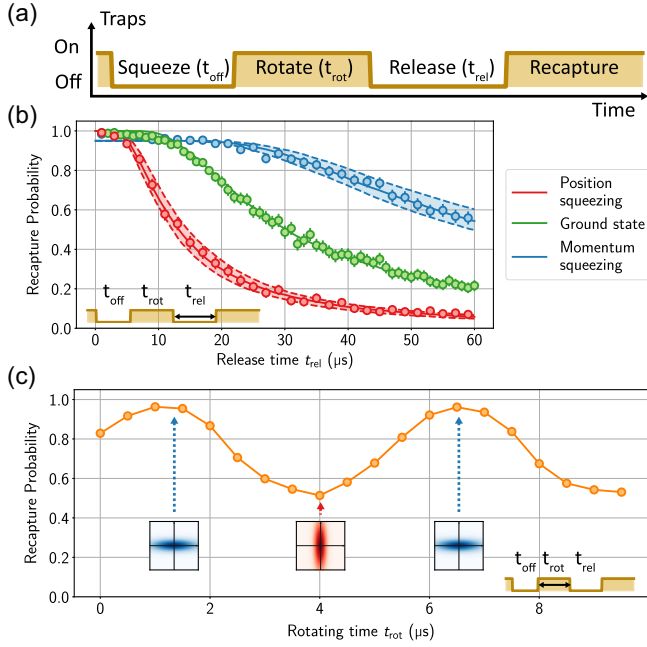


FIG. 2. Probing the squeezed states. (a) Experimental sequence for generating a squeezed state, followed by a release-and-recapture probing. (b) Release-and-recapture results for three different motional states. The two squeezed states were prepared for $t_{\text{off}} = 2.0 \mu\text{s}$, and a trapping frequency $\omega_r \approx 2\pi \times 110 \text{ kHz}$. Solid lines are simulations, the dashed lines and colored area indicate the error bar on the fitted β (see main text). (c) Oscillation of the squeezed state in the trap, performed for a slightly different trapping frequency $\omega_r \approx 2\pi \times 95 \text{ kHz}$ ($t_{\text{off}} = 2.0 \mu\text{s}$ and $t_{\text{rel}} = 20 \mu\text{s}$). Data in (b) [(c)] are obtained by averaging the signals of 72 (50) tweezers over 15 (30) experimental runs. Error bars, which are almost always smaller than the data symbol, represent the 2σ standard error on the mean.

optical tweezers platform was proposed in Ref. [33], while the more standard tomography approach was realized using time-of-flight measurements [34].

Here, we chose to simply measure only the momentum distribution uncertainty, and by observing its oscillation as the distribution rotates [see Fig. 1(d)], we will infer the squeezing factor for the position distribution [35]. The rms (root-mean-square) velocity of the atoms is estimated with a standard release-and-recapture experiment and classical Monte Carlo simulation [36]. The experimental sequence is shown in Fig. 2(a): after having squeezed and rotated the state, we release the atom again for a varying time t_{rel} before trying to recapture it. Figure 2(b) shows the measured recapture probability as a function of t_{rel} for three different states: the motional ground state (green), and two squeezed states obtained for $t_{\text{off}} = 2.0 \mu\text{s}$ corresponding to an expected squeezing factor $\beta = 1.9$. The squeezed state with minimal momentum uncertainty is obtained for $t_{\text{rot}} = 1.4 \mu\text{s}$ (blue), and the one with minimal position uncertainty for $t_{\text{rot}} = 3.6 \mu\text{s}$ (red), i.e., after an additional $\pi/2$ rotation in phase space. The data for the ground state (green) serves as a reference for fixing the parameters of the

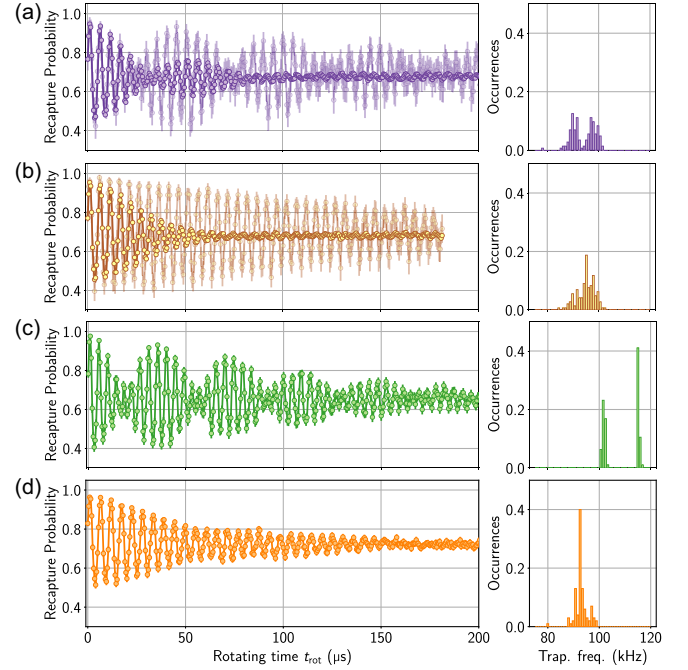


FIG. 3. Oscillations of squeezed states averaged over an array of optical tweezers. The displayed data correspond to the experimental parameters $t_{\text{off}} = 2.0 \mu\text{s}$ and $t_{\text{rel}} = 20 \mu\text{s}$, while varying the rotating time t_{rot} . The subfigures show the recapture probability for different corrections: (a) no corrections, (b) anisotropy compensation, (c) inhomogeneity compensation and (d) both corrections. The trapping frequencies, extracted from Fourier transforms of the individual atomic signals [shown in the lighter curve in (a) and (b)], are plotted in the right panels. For (c) and (d), we used an atomic array with 50 instead of 72 atoms, averaged over 30 iterations compared to 200 iterations for (a) and (b).

Monte Carlo simulation (solid line). For the blue data, the slower decay indicates a momentum squeezing, which is well captured by the simulation using a reduced momentum uncertainty by a best fitting factor $\beta = 1.96$. Whereas for the red data, the decay is faster and matches an increased momentum spread by the same factor. The best fitting factor, and the associated error bar, are estimated via the χ^2 method, and gives $\beta = 1.96 \pm 0.13$, in agreement with the expected value 1.9.

So far, we have reported the generation of a state whose momentum spreading is reduced and then increased by a factor 1.96. To further infer that the position uncertainty is reduced by this same factor, we need the distribution to rotate unperturbed, which is true only for an ideal harmonic potential. We can check this experimentally by observing if the momentum uncertainty periodically recovers its minimal value as the distribution rotates with increasing t_{rot} . For example, we see that it barely decreases over half a trapping period in Fig. 2(c).

Collective oscillation of the squeezed states—Observing the squeezed state evolution for longer times, up to 20 trapping periods ($t_{\text{rot}} \simeq 200 \mu\text{s}$), reveals striking features. In Fig. 3(a), the oscillation displays a beat note and an overall

damping. This can be explained from the measured distribution of trapping frequencies over the array of tweezers (shown on the right panel): the distribution is bimodal (inducing beating), and with a finite spread (inducing damping). The bimodality comes from a slight radial anisotropy (10%–20%) between the x and y axis, which naturally occurs at high NA [37,38]. The inhomogeneities over the array come from imperfections of the hologram. We also display typical single-atom signals, where the damping is much reduced confirming that the damping of the averaged signal originates from inhomogeneities.

Counteracting on the SLM phase pattern allows us to compensate for the anisotropy and the shape inhomogeneities, as we recently demonstrated [31]. After anisotropy compensation, the resulting oscillation [Fig. 3(b)] shows no more beatnote and only damping (again, much faster than for a single atom), indicating that the two radial modes now have the same trapping frequency, which is confirmed by the distribution of individually measured trapping frequencies (side panel). Alternatively, if we modify the phase pattern to decrease the frequency spread over the array, decreasing it by one order of magnitude to $\Delta\omega/\omega < 0.01$, we now observe a much reduced damping and a clearer beat note [Fig. 3(c)]. Finally, when applying both anisotropy and inhomogeneity corrections [Fig. 3(d)], we minimize the degradation of radial squeezing over time.

If the position squeezing is only needed along one of the two radial directions is needed [32], then correcting the anisotropy is not required and we can instead optimize only the inhomogeneity as in (c). However, since our release-and-recapture measurement protocol does not distinguish between the x and y velocities, we prefer to synchronize both radial modes, at the cost of a slightly worse homogeneity, which, however, will not be limiting anymore as explored more quantitatively below.

Squeezing limitations—In this final part, we calculate how much position squeezing can be achieved given the imperfect rotation of the squeezed phase-space distribution while the atoms oscillate in the traps. As just observed, a first limitation comes from different frequencies for the radial mode. In Fig. 4(a), we consider two radial modes (for example, from two different traps), whose relative frequency difference is $\Delta\omega/\omega = 0.2$, initially squeezed in momentum by a factor $\beta = 3$. Because of their different angular speeds, the two rotating distributions will not be aligned to the p_x axis at the same time, leading to a broadening of the position distribution averaged over the traps as compared to an ideal position-squeezed state [see inset of Fig. 4(a) showing the projection of the phase-space distribution on x]. For a given value of $\Delta\omega/\omega$, this sets a maximal squeezing factor $\beta_{\max}$ achievable over the atomic array. A scaling law can be obtained by equating the x projection of the imperfectly rotated distribution $\beta_{\max} \sin(\Delta\omega t_{\text{rot}})$ to the ideal squeezing value $1/\beta_{\max}$. Using $t_{\text{rot}} \propto 1/\omega$, we obtain: $\beta_{\max} \propto (\Delta\omega/\omega)^{-1/2}$.

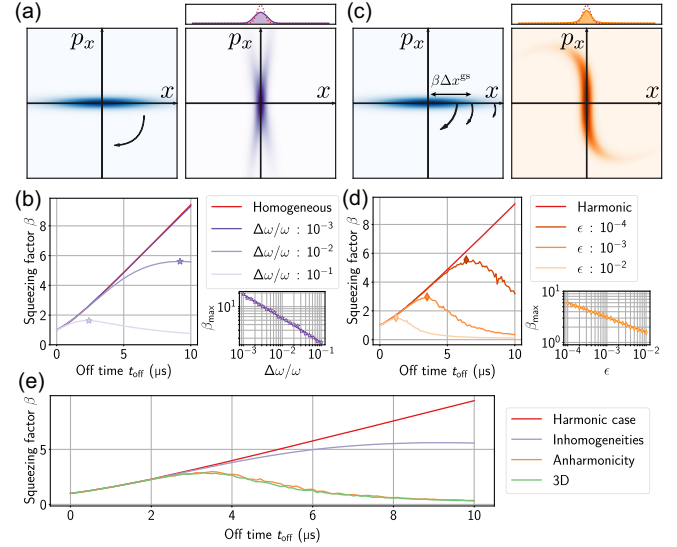


FIG. 4. Squeezing performance. (a) Sketch of the effect of the trapping frequency inhomogeneity on the squeezing factor. Here we show the rotation of two motional states for a trapping frequency difference $\Delta\omega/\omega = 0.2$, and the corresponding projection on the x axis (the red dashed line represents the ideal squeezed state). (b) Squeezing factor as a function of t_{off} , for different trapping frequency inhomogeneities. The inset shows the maximal squeezing factor as a function of $\Delta\omega/\omega$. (c), (d) Same as (a), (b) for the study of the anharmonicity. The anharmonicity is $\epsilon = 10^{-2}$ in (c). (e) Squeezing factor as a function of t_{off} , in an ideal 1D harmonic potential (red), in the presence of our remaining inhomogeneities $\Delta\omega/\omega = 10^{-2}$ (purple), for the non-approximated 1D Gaussian potential (orange) and finally for the actual 3D potential (green).

We then perform numerical calculations by evolving squeezed states in 1D harmonic potentials with trapping frequencies distributed following a Gaussian distribution with rms $\Delta\omega$ [Fig. 4(b)]. We extract a maximal squeezing factor (star symbols) that indeed scales as $\beta_{\max} \propto (\Delta\omega/\omega)^{-0.53}$. For the uncorrected hologram, with $\Delta\omega/\omega \sim 0.1$, the squeezing factor would be limited to $\beta_{\max} = 1.6$. After optimization of the hologram, which decreases the inhomogeneity by an order of magnitude (and was used for Fig. 2 data), we could hope to reach a squeezing of more than 5 (14 dB).

We now consider a second, more inherent, imperfection of the tweezers: its anharmonicity. As seen in Fig. 1(a), the harmonic approximation of the tweezers potential breaks down when the atom explores a too large region. This results in a distortion of the phase-space distribution [Fig. 4(c)], as the atom evolves more slowly in the subharmonic region. This anharmonicity can be captured by expanding a 1D Gaussian potential to fourth order:

$$U(x) = -U_0 + \frac{1}{2}m\omega_f^2 x^2 \left(1 - \epsilon \left(\frac{x}{\Delta x^{\text{gs}}}\right)^2\right), \quad (4)$$

where $\epsilon = \hbar\omega_r/(8U_0) \approx 10^{-3}$ for our tweezers parameters. We can then derive a scaling law for the maximal squeezing factor $\beta_{\max}$ achievable for a given anharmonicity ϵ . We estimate the angular frequency spread $\Delta\omega$ within the phase-space distribution by comparing its value at the center to that at a characteristic distance $x = \beta\Delta x^{\text{gs}}$ [see illustration in Fig. 4(c)]. Using the fourth order approximation gives $\Delta\omega = \omega \left(1 - \sqrt{1 - \epsilon\beta^2}\right)$, leading to $\beta_{\max} \propto \epsilon^{-1/4}$. Numerical results, shown in Fig. 4(d), are obtained by evolving a squeezed state in 1D Gaussian potential with fixed trapping frequency ω but varying anharmonicity ϵ (obtained by adjusting the trap depth U_0 and waist w accordingly). We obtain $\beta_{\max} = 3$ at $\epsilon = 10^{-3}$ and a power law of $\beta_{\max} \propto \epsilon^{-0.28}$.

Finally, we compare these different limitations for our tweezer parameters in Fig. 4(e), illustrating that the anharmonicity becomes the limit once the trap shapes are optimized by holography. We have also performed a simulation with the full 3D potential of the tweezers, showing no difference from the 1D case, justifying *a posteriori* why we have not discussed it earlier. There is also squeezing along the longitudinal direction, but it is almost negligible ($\beta_z = 1.1$ at $t_{\text{off}} = 2 \mu\text{s}$) given the much smaller trapping frequency ω_z in that direction, and would give no observable signature on the release-and-recapture experiment which is mostly sensitive to radial motion.

Discussions—Having analyzed the limitations set by the trap shape, we now conclude on the squeezing level demonstrated in Fig. 2. Taking into account the possible imperfect rotation of the state, and the imperfect preparation of the motional ground state, we claim the generation of a position squeezing level of 5.4 dB. We refer to the End Matter for more details.

A more straightforward approach to reduce the position uncertainty of the wave function could have been simply to increase the trapping frequency. But this would not be a scalable strategy: as $\Delta x^{\text{gs}} \propto U_0^{-1/4}$, reducing the position spreading by a similar factor of 2 would require a sixteen-fold increase in optical power. We also note that deeper tweezers reduce the anharmonicity parameter ($\epsilon \propto U_0^{-1/2}$), but will also not be a scalable solution to push further the achievable squeezing given the unfavorable scaling $\beta_{\max} \propto \epsilon^{-1/4}$.

Finally, we envision several applications for these position-squeezed states. Given that they are not static, we target applications on timescale of a few hundred nanoseconds at most. Combining them with strong Rydberg interactions, which can reach several hundreds of MHz [32,39], is a first direction. They could enable the engineering of more complex Hamiltonians or the use of motional degrees of freedom in quantum information processing via spin-motion coupling [40–47]. We could exploit squeezed states to reduce decoherence in spin-exchange oscillation [32,48]. We could also use quantum amplification based on squeezing [28] to measure momentum kicks between Rydberg atoms [49]. Quantum

amplification requires unsqueezing, which can be realized on our setup by repeating the squeezing protocol once the distribution reaches the mirror-image position relative to the p axis after rotation [50]. A second direction is to combine the squeezing technique with flying atoms [51] to prepare ordered arrays within the subwavelength regime where fascinating collective light-matter phenomena emerge [52,53]. Finally, one could also exploit the anharmonicity for the generation of nonclassical states [54–56].

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Data availability—The data that support the findings of this article are openly available [57].

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End Matter

In this section, we derive the position squeezing level measured in Fig. 2, by first estimating the effect of the imperfect rotation of the motional state on the timescale of interest, and then by considering the imperfect preparation of the motional ground state.

Perturbation of the rotation—As explained in the main text, our method to determine the position squeezing relies on a measurement of the momentum squeezing and the rotation of the motional state inside the optical tweezers. For the chosen experimental value $t_{\text{off}} = 2 \mu\text{s}$, the numerical results of Fig. 4(e) demonstrates that the anharmonicity is not yet a limit. In addition, the results in Fig. 2(b) were obtained with an optimized hologram, so that trap anisotropy and inhomogeneities do not play a role. Consequently, we have shown that the perturbative terms in x^2 (different radial frequencies) and in x^4 (anharmonicity) do not prevent the motional state to neatly rotate on the timescale of interest, so that the squeezing factor measured on the momentum quadrature is also the one of the position quadrature.

For completeness, we eventually consider a perturbative term in x^3 , which could originate from a coma-type aberration. On short rotation timescale, this term would have the detrimental effect of reducing the squeezing factor on x while leaving the one on p_x unaffected, due to its odd exponent; which prevents us from characterizing it with our momentum-measurement based method presented in Fig. 2. Nevertheless, an x^3 perturbation would lead to a

damping of the momentum standard deviation oscillation on long timescales, like the one we observe on the single-atom trace of Fig. 3(b). We then estimate the upper bound of this x^3 perturbation by finding the coefficient which qualitatively reproduces the single-atom damping (considering anharmonicity only, the best agreement is obtained for an ϵ value close to the expected one $\epsilon \approx 10^{-3}$, indicating that the damping is already well explained by the quartic term, and that we are overestimating the cubic perturbation). For the maximal x^3 perturbation estimated that way, the squeezing factor on position would still be only 10% smaller than the one on momentum.

This allows us to claim that the squeezed distribution oscillates unperturbed, at least on the 90% level, on the timescale of interest.

Thermal excitations—To conclude on the squeezing level with respect to the standard quantum limit, one eventually needs to take into account the remaining occupation of higher vibrational levels of the initial state $\langle n_x \rangle \approx 0.05$, increasing by a factor $\eta = 1.05$ the standard deviations of the motional quadratures above the standard quantum limit. The measured factor $\beta = 1.96$ is a reduction factor as compared to the initial standard deviations, so it needs to be corrected by this η increase, leading to a squeezing level below the standard quantum limit $\beta/\eta = 1.87$ (5.4 dB).