

Accelerated state expansion of a nanoparticle in a dark inverted potential

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While the wave packet of a massive particle grows linearly under free dynamics, it grows exponentially in an inverted harmonic potential, offering a pathway to rapidly increase quantum fluctuations to macroscopic dimensions. In this work, we experimentally demonstrate this principle by expanding the center-of-mass thermal state of a 125 nm silica nanoparticle to a position uncertainty of 43.4 nm within 260 μ s. This expansion, achieved using an inverted dark potential to minimize decoherence from photon recoil, represents a 952-fold increase, reaching a scale comparable to the nanoparticle's physical size. This work represents a key advancement toward preparing macroscopic quantum superpositions at unprecedented mass and length scales.

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Quantum mechanics predicts that a massive particle cannot exist in a motional state exhibiting no uncertainty in phase space, even at zero temperature. Specifically, the Heisenberg principle states that due to the noncommutation of a pair of canonically conjugated variables, such as position and momentum, the product of their standard deviations is bounded from below by $\hbar/2$ [1]. These Heisenberg-limited states can be prepared by reducing the entropy of their motional state, namely, through cooling. This has been experimentally achieved with massive objects confined in harmonic potentials, including a single ion [2–4], a nanoparticle containing billions of atoms [5–10], and various micromechanical resonators [11–15], with some resonators having masses as large as $\sim 10^{19}$ atomic mass units [16]. In harmonic potentials of frequency Ω , the position standard deviation of these Heisenberg-limited states is on the order of the so-called zero-point motion $\sqrt{\hbar/(2m\Omega)}$, where m is the mass of the object. Consequently, the more massive the object, the smaller the position uncertainty for these Heisenberg-limited states prepared in a harmonic potential. Thus, it appears that the more massive an object is, the more microscopic its quantum effects become.

However, quantum mechanics allows, in principle, the amplification of quantum fluctuations by increasing the standard deviation of one canonical variable while reducing the conjugate one, such that their product remains close to $\hbar/2$. These motionally squeezed states can exhibit position uncer-

tainties much larger than $\sqrt{\hbar/(2m\Omega)}$, potentially extending beyond the size of the particle like in matter-wave experiments [17–20]. Such macroscopic quantum states of massive objects become extremely sensitive to external signals and decoherence, and can be used to test quantum mechanics in regimes where collapse models predict the breakdown of the superposition principle [21,22]. Additionally, they can be transformed into non-Gaussian states exhibiting negativities in their Wigner function [23,24]. In clamped oscillators, where the tight harmonic potential is fixed, moderate motional squeezing up to 4.7 dB [25] was reported. Recently, the possibility of quenching the harmonic potential of a levitated nanoparticle and exploiting motional dynamics in the absence of laser light—which would generate decoherence due to photon scattering—has been proposed as a promising platform for achieving unprecedented levels of motional squeezing [26–29]. A particularly promising method involves using a dark (nonoptical) inverted harmonic potential since, in this case, the generation of motional squeezing is exponential in time [26,30,31]. This stands in stark contrast to the linear regime induced by dynamics in shallower harmonic potentials or even free dynamics. Exponentially fast expansion of states is relevant to overcome challenges imposed by decoherence and stability, affecting both the dynamics during the protocol and the uncertainties between experimental runs [26,28].

In recent years, experiments aiming at the ambitious goal of preparing macroscopic quantum superposition states of nanoparticles have been proposed [21,26–30,32–34]. The often required ground-state cooling of an optically levitated nanoparticle was achieved using either an optical cavity [5,10,35,36] or active feedback [6,7,9,37]. Furthermore, motional state expansion, linear in time, was reported using either shallower optical potential or free dynamics, starting from both the ground state [38,39] and thermal state [40,41] of a nanoparticle. Recently, optically inverted potentials have also

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been realized [42]. However, squeezing protocols utilizing optical means are limited by decoherence induced by photon scattering [43–45]. Recently, hybrid schemes have been reported in the classical regime [46–48], where an optically levitated nanoparticle is released into a shallower harmonic electrical potential [49]. In this Letter, we report a significant progress: the controlled release of an optically trapped and cooled nanoparticle into an inverted dark potential, resulting in an exponentially fast expansion of its motional state [31]. This is achieved using a hybrid electro-optical platform integrating state-of-the-art motion detection and cooling to low phonon occupations with engineered on-chip electrostatic [50] and radio-frequency (rf) potentials. Our approach enables rapid squeezing protocols while effectively suppressing photon recoil heating.

More specifically, we study the expansion of a nanoparticle's state in a dark, inverted, and electrostatic potential $U(z) = -m\omega_z^2 z^2/2$ parametrized by the frequency ω_z . The levitated nanoparticle is released from a tighter optical harmonic trap $U_0(z) = m\Omega_z^2 z^2/2$ at the center of the inverted potential, leading to an exponentially accelerated state expansion [26,28,31]. Throughout the entire protocol described here, the particle occupies a Gaussian state fully described by the mean, variances, and covariance of the position $\hat{z}$ and momentum $\hat{p}$ quadratures. Initially, the particle occupies the symmetric phase space of a thermal Gaussian state at temperature T with an average phonon occupation of the center-of-mass (CoM) motion given by the Bose-Einstein distribution $\bar{n} = [\exp(\hbar\Omega_z/(k_B T)) - 1]^{-1}$, before the state expands in phase space following the inverted harmonic potential, as depicted in Figs. 1(a) and 1(b). In the zeroth-order secular approximation [51] neglecting micromotion, the position variance $\sigma_z(t)^2 = \langle \hat{z}(t)^2 \rangle - \langle \hat{z}(t) \rangle^2$ follows [52]

$$\sigma_z(t)^2 = \sigma_z(0)^2 \left[\cosh^2(\omega_z t) + \frac{\Omega_z^2}{\omega_z^2} \sinh^2(\omega_z t) \right] - \frac{\hbar\Omega_z\Gamma_z^1}{m\omega_z^2} \left[t - \frac{\sinh(2\omega_z t)}{2\omega_z} \right], \quad (1)$$

where the first term describes the coherent dynamics and the second term describes the incoherent dynamics. The displacement noise Γ_z^1 links to the heating rate $\dot{E}_z = \hbar\Omega_z\Gamma_z^1$, dominated by residual gas and electric field noise. The expansion coefficient η_z after the release time t_r is defined as the ratio of the initial and final delocalization lengths $\eta_z = \sigma_z(t_r)/\sigma_z(0)$. Note that our setup allows us to compare the frequency jump protocol [48,52,53] simultaneously in the transversal directions.

We consider a charged, spherical silica (SiO_2) nanoparticle of mass $m \approx 1.95$ fg corresponding to a radius $R \approx 63$ nm trapped by an optical tweezer in the z direction, leading to a 3D harmonic potential with mechanical eigenfrequencies $(\Omega_x, \Omega_y, \Omega_z)/(2\pi) \approx [185, 171, 43.5]$ kHz. The optical trap is interfaced to a six-layer printed circuit board (PCB), as shown in Figs. 1(c) and 1(d), which integrates a Paul trap [54–57] for dark confinement with segmented virtual ground electrodes and electrodes for cold damping [50,58–61]. Individual PCB layers along the optical propagation axis z allow for the generation of an electrostatic, inverted harmonic

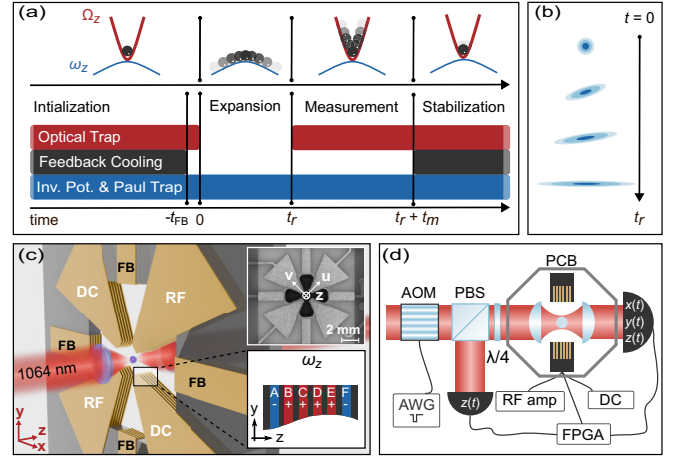


FIG. 1. State expansion in dark inverted potentials. (a) Experimental sequence for state expansion in a dark inverted potential. The particle state is released from a tight optical trap at Ω_z (red) into an inverted, electrostatic potential at ω_z ($t = 0$). The particle delocalizes in the dark and is retrapped at $t = t_r$. Its position at t_r maps into the oscillation amplitude in the optical tweezers before reinitialization ($t_r + t_m$). The feedback is turned off at $-t_{FB} = -5$ μ s. Retrodiction on a $t_m = 500$ μ s time trace in the optical trap reconstructs the particle state at $t = t_r$. The sequence is repeated 400 times to reconstruct the phase-space distribution. (b) Phase-space evolution in an inverted potential. The position variance $\sigma_z(t)^2$ in the inverted harmonic potential grows until the particle is retrapped at $t = t_r$. (c) Hybrid electro-optical chip. A nanoparticle is trapped at the laser focus ($\lambda = 1064$ nm, NA = 0.77) and surrounded by a six-layer PCB. The radio-frequency electrodes create Paul trap confinement in the x, y plane ($V_{RF} = 800$ V, $\Omega_{RF}/(2\pi) = 25$ kHz), cold damping FB electrodes cool the nanoparticle CoM to $\bar{n}_{x,y,z} \approx [721, 3763, 10]$, and dc electrodes generate an electrostatic potential along z ($V_{A,F} = -14$ V, $V_{B,-E} = +14$ V). Top inset: PCB with the Paul trap axes (u, v). Bottom inset: Voltage configuration for inverted harmonic potential. (d) Optical setup. The particle position is measured and feedbacked using homodyne back-detection (z) and balanced forward detection (x, y). The laser is switched on and off with an AOM controlled by an AWG. The Paul trap is controlled by an rf amplifier and the feedback signals are processed by an FPGA.

potential at the center of the PCB by applying tailored dc voltages to each layer [see lower inset in Fig. 1(c)]. From an independent measurement, the inverted harmonic potential of $\omega_z/(2\pi) = 1.4$ kHz is characterized for a particle with $n_q = +84$ elementary charges at the maximum voltage output of the low noise dc source (see Ref. [62]). The Paul trap exhibits secular frequencies $(\omega_u, \omega_v)/(2\pi) \approx [2.7, 2.5]$ kHz, where the uv plane is rotated by $\theta_t \approx 45^\circ$ with respect to the xy plane [see top inset in Fig. 1(c)] [48].

The particle's CoM motion is detected in forward balanced detection in x, y, z [63] and in backward homodyne detection for z [7]. We initialize the CoM state by applying 3D active feedback with cold damping reducing the phonon occupation to $\bar{n}_{x,y,z} \approx [721, 3763, 10]$ phonons, at pressure $p_g \approx 7.7 \times 10^{-7}$ mbar. Motional cooling enables to decouple the CoM modes of the particle, such that we can treat their dynamics independently. Here, we primarily focus on the inverted dynamics along the z axis.

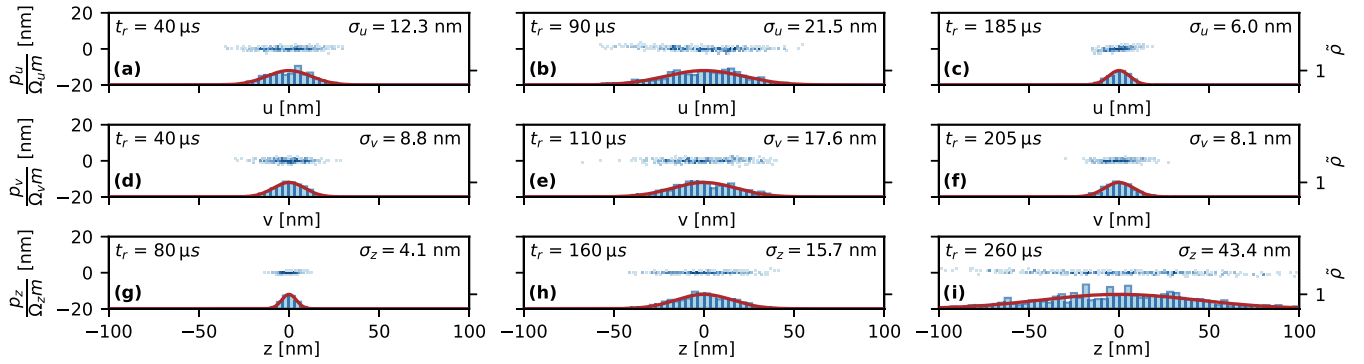


FIG. 2. Phase-space distribution of the particle's CoM in the u [panels (a)–(c)], v [panels (d)–(f)], and z directions panels [(g)–(i)] at selected t_r expanding in electrostatic potentials. The 1D histograms show the amplitude normalized, projected position distribution $\tilde{\rho}$. The red curves represent the estimated Gaussian with standard deviation $\sigma_j(t)$, $j \in \{u, v, z\}$. The distributions are constructed from 400 repetitions of the experimental sequence and plotted as 2D histograms. The frequencies $\Omega_u = \Omega_v = 2\pi \times 173$ kHz are defined assuming the uv system is rotated by $\theta_l \approx 45^\circ$ to the xy basis. The u and v directions are dominated by Paul trap dynamics with maximum position variances $\sigma_{u,v} = [21.5, 17.6]$ nm at $t_r = T_{u,v}/4 = \pi/(2\omega_{u,v})$ [panels (b) and (e)], after which $\sigma_{u,v}(t)$ decreases due to the confining potential, resulting in a recompressed state at $t_r = T_{u,v}/2 = \pi/\omega_{u,v}$ [panels (c) and (f)]. Along the z direction, the particle's position variance grows exponentially as it evolves in a dc inverted potential, reaching a state size of $\sigma_z = 43.4$ nm.

The experimental protocol is illustrated in Fig. 1(a). The precooled nanoparticle is positioned at the Paul trap center by aligning it to the optical trap. The initial motional state sizes are $(\sigma_x(0), \sigma_y(0), \sigma_z(0)) \approx [183, 435, 45.6]$ pm. Following the deactivation of the electrical feedback, the particle is released from the optical trap at a fixed phase ϕ of the applied RF Paul trap voltage V_{RF} . The particle evolves in the dark electrical potential for a duration t_r before being recaptured in the optical trap. Once recaptured, the particle's position is measured for $t_m = 500$ μ s before reinitializing the state with feedback. This sequence is repeated 400 times for each release duration t_r . Throughout the entire protocol, both the Paul trap and the inverted potential remain active. To compensate for time-dependent position shifts caused by micromotion and electrical stray fields, electrical compensation fields are individually adjusted for each release time using the feedback electrodes. The particle's position and velocity at t_r are extracted using retrodiction on the position time trace [38,49] (see Ref. [62]). The position standard variation increases during the measurement due to heating by $(\delta\sigma_x, \delta\sigma_y, \delta\sigma_z) \approx [91, 102, 321]$ pm, corresponding to $\delta\tilde{n}_{x,y,z} \approx [178, 206, 520]$ phonons related to the initial thermal state with $\Omega_x, \Omega_y, \Omega_z$. This effect overestimates the measured state $\sigma_j(t)^2$ with $j \in \{u, v, z\}$, which we add to the error bar in Fig. 3. However, the contribution of $\delta\sigma_j$ to the position uncertainty, overestimating $\sigma_j(t)$, is only significant for small states of $\sigma_j(t) < 1$ nm.

Results. Our experimental setup enables a direct comparison between the dynamics in an inverted potential along z ($\Omega_z \rightarrow i\omega_z$) and the dynamics induced by the frequency jump protocol along u, v ($\Omega_{x,y} \rightarrow \omega_{u,v}$) [26,27,38,40,42,49,53]. Figure 2 shows the phase-space distribution for u, v, z at different release times t_r , where the phase-space angle due to experimental delays is adjusted in postprocessing. The phase-space distributions are represented as 2D histograms constructed from 400 repetitions of the experimental protocol. For each t_r , the 1D histogram in the inset shows the projection on the position axis. These are overlaid with a normalized Gaussian curve with the extracted standard deviation $\sigma_j(t)$, visually confirming the Gaussian distribution of the data.

Figures 2(a)–2(f) display the expansion of the initial state, followed by recompression along u and v . In contrast, we observe an ever-growing expansion along the inverted potential in the z direction as seen in Figs. 2(g)–2(i). Note that only the major axis of the phase-space distribution can be resolved reliably, unlike the minor axis due to heating effects during the measurement and insufficient accuracy in the phase-space angle (see Ref. [62]).

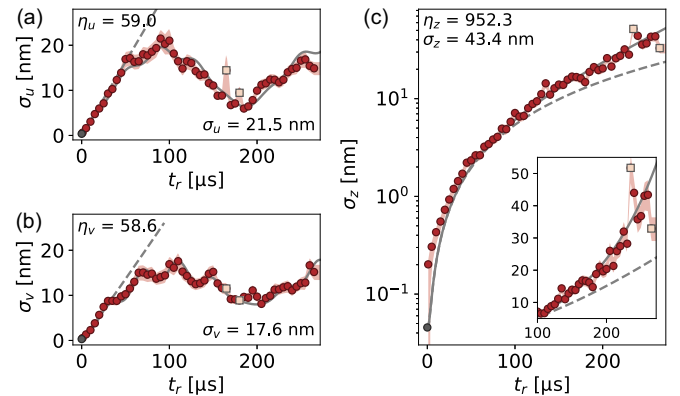


FIG. 3. State expansion of the particle's motional state in a dark electrical potential. (a), (b) In the frequency jump protocol in the Paul trap $\Omega_{x,y} \rightarrow \omega_{u,v}$, the particle state uncertainty $\sigma_{u,v}(t_r)$ expands to a maximum at $t_r = T_{u,v}/4$ followed by a recompression at $t_r = T_{u,v}/2$. Maximum state sizes of $\sigma_{u,v}(T_{u,v}/4) = [21.5, 17.6]$ nm and expansion ratios of $\eta_{u,v} \approx [59, 58.6]$ are reached. (c) In the inverted dc harmonic potential along z , the phase-space expansion is unbound with a maximal state size $\sigma_z(t_r = 260 \mu\text{s}) = 43.4$ nm corresponding to $\eta_z \approx 952.3$. The shaded error bars represent the uncertainty on the position standard deviation and the heating during the measurement. The theory (solid line) fits the data including micromotion (see Ref. [62]) with $\Gamma_j^1, \Omega_j, \omega_j, \sigma_j(0)^2$ and ϕ_j with $j \in \{u, v, z\}$ as fit parameters. The dashed line corresponds to free expansion (see main text). The shaded square points highlight states deviating from a Gaussian distribution due to experimental imprecision.

Figure 3 depicts the position standard deviation $\sigma_j(t)$ of the phase-space distribution in dependence on the release time t_r along u, v, z . The shaded error bars represent the standard deviation of the position uncertainties and the measurement contribution to $\sigma_j(t)$ associated with $\delta\bar{n}_{x,y,z}$ overestimating the state size. For the frequency jump protocol along u, v in Figs. 3(a) and 3(b), we observe a maximum expansion at $T_{u,v}/4$ followed by a recompression at $T_{u,v}/2$, where $T_{u,v} = 2\pi/\omega_{u,v}$ as predicted by the theory. The micromotion due to the Paul trap confinement leads to additional smaller amplitude oscillations dictated by the Paul trap frequency Ω_{RF} . Figure 3(c) displays the exponentially accelerated state expansion along the inverted potential. The micromotion plays only a negligible role in the inverted potential. The solid line is a fit to the theory [Eq. (1) for z and Eq. (46) in Ref. [62] for u, v] with $\Gamma_j^1, \Omega_j, \omega_j, \sigma_j(0)^2$ and ϕ_j as fit parameters (see Table S1 in the Supplemental Material [62]). The decoherence rate Γ_j^1 is only considered in the zeroth-order secular motion [second term in Eq. (1)], since higher-order terms are negligible. For comparison, Fig. 3 displays the state expansion in free dynamics in $j \in \{u, v, z\}$ (dashed line) following

$$\sigma_j(t)^2 = \sigma_j(0)^2(1 + \Omega_j^2 t^2) + \frac{2}{3} \frac{\dot{E}_j}{m} t^3, \quad (2)$$

where $\dot{E}_j$ is the heating rate accounting for the expansion due to heating. From Fig. 3, we observe that the free expansion, the frequency jump, and the inverted protocol follow the same dynamics for small release times t_r . Nevertheless, the expansion in u and v is bound by the Paul trap confinement to the ratio $\Omega_{x,y}/\omega_{u,v}$, neglecting the micromotion. In contrast, free expansion is unbound but slower than the inverted dynamics. The standard deviation of the state reaches a maximum $\sigma_z = 43.4 \text{ nm}$ at $t_r = 260 \mu\text{s}$, corresponding to a state size comparable to the particle's radius and an exponentially fast expansion ratio $\eta_z \approx 952.3$. The final state size, only limited by our linear detection range, beats the free expansion and would continue to grow for longer expansion times, which is the main result of this manuscript.

Nonetheless, to reach large quantum states, it is of utmost importance to expand the state coherently. The figure of merit here is the coherence length [21] $\xi_j(t) = \sqrt{8\mathcal{P}_j(t)\sigma_j(t)}$, with the state purity $\mathcal{P}_j(t) = \text{tr}[\hat{\rho}_j(t)^2] = \hbar/(2\sqrt{\sigma_j(t)^2\sigma_{p_j}(t)^2 - \sigma_{jp_j}(t)^4})$, where $\hat{\rho}_j(t)$ is the density matrix, $\sigma_{p_j}(t)^2$ is the momentum variance, and $\sigma_{jp_j}(t)^2$ is the covariance of the Gaussian state. For the case of the optical trap, the purity of the initial thermal state is given by $\mathcal{P}_j = (2\bar{n}_j + 1)^{-1}$. The variances and covariances depend on the very same parameters, namely, $\Gamma_j^1, \Omega_j, \omega_j, \sigma_j(0)^2$ and m (see Ref. [62]). Assuming that our micromotion model captures well the dynamics of the system not only for the measured position variance $\sigma_j(t)^2$ (see Fig. 3) but also for the momentum variance $\sigma_{p_j}(t)^2$ and covariance $\sigma_{jp_j}(t)^2$, we estimate the coherence length $\xi_j(t)$ by using the fitting parameters from Fig. 3. In Fig. 4, we deduce $\xi_j(t)$ for the current experiment (solid red line) in the presence of a fitted total heating rate of $\dot{E}^{u,v,z} = k_B \times [8.47, 11.25, 5.91] \text{ K/s}$, which stands in reasonable agreement with theoretically predicted values (see Ref. [62]). Moreover, theory predicts electrical

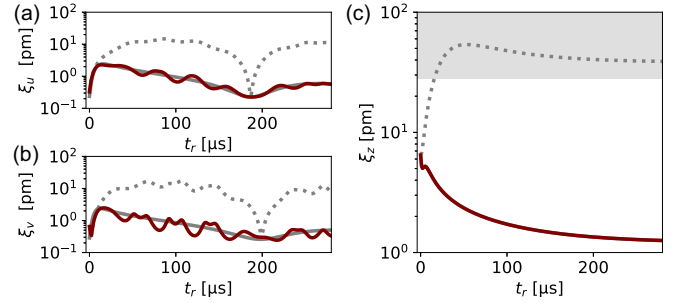


FIG. 4. Theoretical estimation of the coherence length $\xi_j(t_r)$ during state expansion based on fitted parameters $\Gamma_j^1, \Omega_j, \omega_j, \sigma_j(0)^2$ and ϕ_j for $j \in u, v, z$ related to the data shown in Figs. 2 and 3. (a), (b) The frequency jump protocol shows revivals of $\xi_{u,v}(t)$ with an overlaying decay at a rate given by $\Gamma_j^1 \neq 0$. (c) The inverted harmonic stabilizes ξ_z after an initial decay. Red solid line: Current heating rate $\dot{E}_j$. Gray dotted line: Improved heating rate $\dot{E}_j/1000$. Gray solid line: Current heating neglecting micromotion. The gray shaded area indicates $\xi_z(t) > \xi_{|n=0}$.

noise to be the largest heating source along u, v , while gas collisions and electric field noise play an equally important role along z . Starting from $\xi(0)/\xi_{|n=0} \ll 1$ along u, v , the coherence length for the frequency jump protocol increases by a factor of $\max[\xi(t)_{u,v}/\xi_{u,v}(0)] \approx [7.1, 3.9]$. This gain in coherence length is followed by an oscillatory behavior at the secular frequencies $\omega_{u,v}$, combined with a rapid decay due to Γ_j^1 . Interestingly, smaller oscillations due to micromotion at Ω_{RF} can also be observed. In contrast, along z , $\xi_z = 6.5 \text{ pm}$ is larger to begin with due to the higher purity of the initial state $\mathcal{P}_z \approx 5 \times 10^{-2}$. We observe a partial decay down to $\xi_z(t) = 1.16 \text{ pm}$ before $\xi_z(t)$ stabilizes to a nearly constant value, as the expansion nearly compensates for the heating. This behavior demonstrates the advantage of the inverted potential over free expansion. Notably, reducing Γ_j^1 by a factor of 1000 would enable much larger coherence lengths displayed by the gray dotted line in Fig. 4. The dominant source of heating that determines Γ_j^1 in our system is residual gas collisions. Current experiments routinely reach ultrahigh vacuum conditions to suppress this contribution [38,64,65]. The second major source of decoherence in our setup is electrical noise, which, due to misalignment, allows leakage of noise from rf amplifiers into the inverted dc axis. Improved alignment should significantly reduce this effect. Blackbody radiation is expected to play a negligible role in our system due to the use of dark potentials [21,28,66]. Additionally, librational modes can be effectively suppressed with current techniques [39,65].

Summary and outlook. In summary, we have demonstrated the accelerated expansion of a nanoparticle's motional state in a dark, inverted harmonic potential. This approach minimizes decoherence due to light scattering and reduces blackbody radiation while enabling faster protocols. Our results highlight the inverted potential as a powerful tool for the generation of large quantum states. We achieve state sizes comparable to the size of the nanoparticle itself and to the largest reported to date [40], attaining an expansion factor of nearly 3 orders of magnitude within a few hundred microseconds—exceeding previous demonstrations by almost a factor of 40 [49]. This

corresponds to a squeezing level of $\mathcal{S}_z = -10 \log_{10}(\eta_z^{-2}) = 59.6$ dB for a pure state in the absence of decoherence and correlations between position and momentum. Our protocol also surpasses the speed of both free-expansion and frequency-jump schemes [49], while operating entirely in the dark, in contrast to optical protocols [42]. However, for low phonon numbers we do not observe an increase in coherence length, which is consistent with the current heating rates and remains below the state of the art [38]. Furthermore, our method enables the modification of the potential from purely inverted potentials to a double-well configuration [28]. The quartic contribution of the double well introduces nonlinear dynamics, enabling the study of non-Gaussian states. Both inverted

and nonharmonic potentials are essential ingredients for generating macroscopic quantum superpositions.

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Data availability. The data that support the findings of this article are openly available [67].

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