Phonon Pair Creation by Inflating Quantum Fluctuations in an Ion Trap

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Quantum theory predicts intriguing dynamics during drastic changes of external conditions. We switch the trapping field of two ions sufficiently fast to tear apart quantum fluctuations, i.e., create pairs of phonons and, thereby, squeeze the ions' motional state. This process can be interpreted as an experimental analog to cosmological particle creation and is accompanied by the formation of spatial entanglement. Hence, our platform allows one to study the causal connections of squeezing, pair creation, and entanglement and might permit one to cross-fertilize between concepts in cosmology and applications of quantum information processing.

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Introduction.—According to quantum field theory, the vacuum is not just empty space, but it is filled with ubiquitous fluctuations, as implied by the Heisenberg uncertainty principle. Although they are not visible directly, these fluctuations have observable consequences such as the Lamb shift of atomic spectral lines [1], van der Waals and Casimir forces [2], or spontaneous emission [3,4]. Intriguingly, extreme conditions, such as a rapid cosmic expansion, can tear these fluctuations apart and, thereby, turn them into pairs of real particles [5–7]. This process of cosmological particle creation is accompanied by the generation of quantum entanglement at large distances. Similar mechanisms cause the Sauter-Schwinger effect [8] and Hawking radiation [9]; see also Ref. [10].

At present, the rate of expansion of our Universe is too slow to create a measurable amount of particles. However, according to our standard model of cosmology, the process of cosmological particle creation played an important role in the early Universe [29]. Together with other effects, it is considered to be responsible for the reheating of our Universe after the inflationary period. Furthermore, a very similar process based on the tearing apart of quantum vacuum fluctuations by an expanding space-time during cosmic inflation explains the creation of the seeds for structure formation. Even though signatures of these effects can still be observed today in the cosmic microwave background radiation, direct tests remain out of reach. As an alternative, some analog features have been observed in several experimental platforms [30–38]. However, preserving the fragile quantum dynamics described above requires fast control of the system on the level of single quanta, as well as close to ideal isolation from the environment. Trapped atomic ions are well suited to study fundamental quantum dynamics as they feature unique fidelities in preparation, control, and detection of quantum states [39–46].

In this Letter, we report on the creation of pairs of particles, precisely phonons, by inflating quantum fluctuations of the motion of trapped ions. The phonon excitation is accompanied by the creation of spatial entanglement and can be described as a squeezing operation. We explain this process as an experimental analog to the cosmological particle creation [47–49].

Cosmological model.—For simplicity, let us consider the Klein-Gordon equation for a scalar (e.g., inflaton or Higgs) field Φ in $1+1$ dimensions:

$$\left(\Box + \frac{m^2c^2}{\hbar^2}\right)\Phi = 0. \tag{1}$$

Here, $\Box$ denotes the d'Alembert operator, c is the speed of light, and $\hbar$ refers to the reduced Planck constant. Analog descriptions for other fields are given in the Supplemental Material [10]. In terms of the conformal time t, the expanding space-time during cosmic inflation can be described by the Friedmann-Lemaître-Robertson-Walker metric,

$$ds^2 = a^2(t)[c^2dt^2 - dx^2], \tag{2}$$

where the time-dependent scale parameter $a(t)$ governs the cosmic expansion. After a spatial Fourier transform

$\Phi(t,x) \rightarrow \phi_k(t)$, the equation of motion (1) for a specific mode k reads

$$\ddot{\phi}_k + \left(c^2k^2 + a^2(t)\frac{m^2c^4}{\hbar^2}\right)\phi_k = \ddot{\phi}_k + \Omega_k^2(t)\phi_k = 0. \quad (3)$$

The squared frequency $\Omega_k^2(t)$ contains the internal (propagating) contribution c^2k^2 as well as the external dynamics $\propto a^2(t)$ in the second term.

With Eq. (3) we have effectively mapped a complex and experimentally inaccessible problem to the dynamical description of a single quantum harmonic oscillator with time-dependent frequency $\Omega_k(t)$; see Ref. [10]. As long as the dynamics is adiabatic, i.e., the external rate of change $\dot{\Omega}_k/\Omega_k$ is smaller than the internal frequency Ω_k, the mode ϕ_k oscillates almost freely and, thus, its quantum state (described by the wave function ψ_k) stays close to the corresponding ground state. However, if the external variation becomes too fast (i.e., violates the above adiabaticity condition $\dot{\Omega}_k \ll \Omega_k^2$), ψ_k cannot adapt and the quantum state deviates from the ground state and turns into an excited, squeezed state, cf. Figs. 1(a) and 1(b). Even after the external dynamics stops (or slows down again), this squeezed state contains pairs of particles, i.e., lasting excitations of the quantum field Φ. The lowest order of the squeezing parameter corresponds to a single pair of particles (with opposite momenta), while stronger squeezing can also produce excitations with higher (but even) numbers of particles. Note that since the general definition of a wave function for quantum fields in curved space-times can be a subtle issue, we use a simple mode description where the quantum state (assuming a pure state) of every mode ϕ_k is described by the wave function $\psi_k(x)$ or, more precisely, $\psi_k(\phi_k)$, which is the time-dependent probability amplitude that the mode k adopts the value ϕ_k [50]; for further details, see Refs. [7,51,52].

In the cosmological particle creation, the external variation is generated by the cosmic expansion (or contraction) which stretches the physical wavelength $\lambda(t) = 2\pi a(t)/k$ of ϕ_k. If this stretching process becomes too fast, the cosmic expansion tears apart the initial quantum vacuum fluctuations and turns them into pairs of particles with opposite momenta $\pm\hbar k$, generating quantum entanglement at large distances [53], cf. Fig. 1(c). Note that this tearing apart process goes along with a large and rapid (i.e., nonadiabatic) change of $\Omega_k(t)$. One can also generate squeezing by small changes in $\Omega_k(t)$, but this requires many resonant oscillations in order to obtain a significant effect. Such pair creation via resonant parametric driving is also often considered in relation to the dynamical Casimir effect [54].

Analog.—In an ion trap, the motional (phonon) modes represent the field Φ. In particular, we consider the linearized radial displacements δq_i of two (or more) ions (labeled by $i, j = 1, 2, \ldots$), which obey the equations of motion

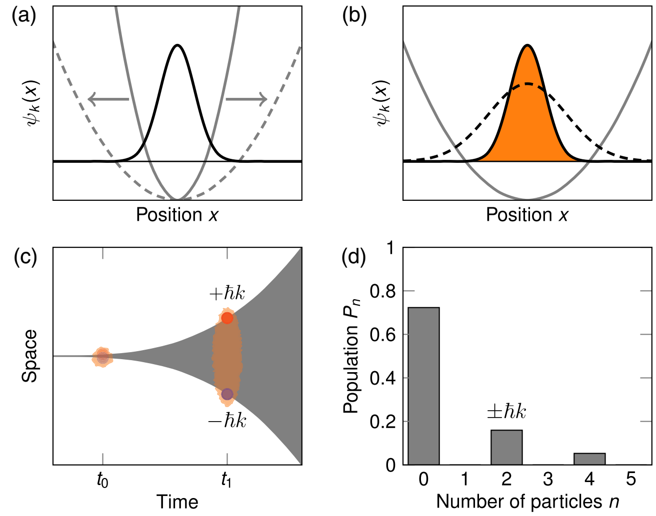

FIG. 1. Particle creation out of quantum fluctuations in harmonic oscillators and the early Universe. (a) The ground state wave function $\psi_k(x)$ (black solid line) of a harmonic oscillator has finite energy and spatial spread (variance) $(\Delta x)^2 \propto 1/\Omega_k$. When Ω_k is changed to a different value (dashed parabola), $\psi_k(x)$ evolves accordingly. (b) If the change of Ω_k is fast, $\psi_k(x)$ cannot follow adiabatically and represents an excited, squeezed state (solid line) within the new potential, that is characterized by a decreased spread $(\Delta x)^2$ compared to the new ground state (dashed line). (c) Schematic of curved space-time during cosmic inflation. Quantum fluctuations are depicted as a pair of virtual particles at some time t_0. Because of inflation the two virtual particles are torn apart until their distance (related to the physical wavelength of the corresponding mode, shaded area) becomes too large for them to recombine and annihilate (at t_1). Thereafter the two virtual particles have become real and move into opposite directions with momenta $\pm\hbar k$. (d) In the harmonic oscillator analog, the creation of multiple particle pairs with momenta $\pm\hbar k$ is evidenced by a squeezed state that contains excitations of even numbers of particles only.

$$\delta\ddot{q}_i + \omega_{\mathrm{rad}}^2(t)\delta q_i + \sum_j M_{ij}\delta q_j = 0, \quad (4)$$

where $\omega_{\mathrm{rad}}^2(t)$ encodes the radial confinement controlled by the time-dependent trapping potential, while the matrix M_{ij} describes the Coulomb interactions between the ions. Identifying $\delta q_i(t)$ with $\Phi(t,x)$, where i corresponds to the discretized position x, this Eq. (4) is analog to Eq. (1).

Diagonalizing the matrix M_{ij} with the eigenvalues λ_I (not to be confused with the wavelength) labeled by the index I is then analog to the Fourier transform $\Phi(t,x) \rightarrow \phi_k(t)$ above and we obtain the mode equations

$$\delta\ddot{q}_I + \omega_{\mathrm{rad}}^2(t)\delta q_I + \lambda_I\delta q_I = 0. \quad (5)$$

This equation is analog to Eq. (3), where the motional modes δq_I (such as the center-of-mass or the rocking mode) correspond to the Fourier modes ϕ_k. The Coulomb interaction strength λ_I then corresponds to the internal (propagating) dynamics c^2k^2. In both cases, they are responsible for generating entanglement. The (external) trapping potential ω_{rad}^2 simulates the mass term $m^2c^4/\hbar^2$. Its time dependence $\omega_{\mathrm{rad}}^2(t)$ represents the external influence and is analog to the cosmic expansion $a(t)$ in Eq. (2). Comparing Eqs. (3) and (5) we identify an expanding universe with an increasing radial confinement $\omega_{\mathrm{rad}}^2(t)$. In both cases, the relative impact of the internal dynamics (λ_I or c^2k^2) decreases, which then causes the tearing apart of quantum fluctuations.

Analog to the cosmological particle creation we consider the creation of pairs of phonons (as vibrational excitation quanta of the ions) under the external variation $\omega_{\mathrm{rad}}^2(t)$. Identifying $\omega_I^2(t) = \omega_{\mathrm{rad}}^2(t) + \lambda_I$ in Eq. (5) with $\Omega_k^2(t)$ in Eq. (3), we may observe particle creation for a nonadiabatic evolution of ω_I. The associated quantum dynamics corresponds to a squeezing operation and, thus, the harmonic oscillator wave function remains an even (Gaussian) function. Therefore, in our analog, the particle pair creation is witnessed by a squeezed state in the ions' motional modes; see Fig. 1(d). Finally, the entanglement of distant space points created by the cosmic expansion corresponds to the motional entanglement of the two ions. After separating the ions and reading out the motional states individually, the entanglement would manifest itself as a thermal (i.e., mixed) state with a temperature set by the squeezing parameter—in complete analogy to Hawking radiation; see Ref. [10].

Moreover, we find that a relative change of the harmonic oscillator potential, such as a certain number of e-foldings, implies an even larger relative change of $a(t)$, i.e., more e-foldings, because the remaining constant term c^2k^2 is positive. Thus, the reported 1.6 e-foldings (see below) imply, in principle, more than 1.6 e-foldings of space $a(t)$. However, we note that in an experimental realization (see below), several nonadiabatic changes of ω_I can be cascaded. Choosing specific timings between individual changes permits constructive accumulation of the squeezing excitations.

Experimental parameter regime.—In our experiments we study the motional modes of two atomic magnesium ions confined in a linear Paul trap; details on the experimental setup are given in Refs. [10,17]. The two ions align along the axial direction of the trap, and in the following we focus on their harmonic motion along the weaker of the radial directions. We distinguish an in-phase (center-of-mass) mode ω_1 from an out-of-phase (rocking) mode with frequency ω_2. The latter is intrinsically robust against homogeneous noise fields, and thus, well suited to implement the dynamics according to Eq. (1). Our apparatus allows for real-time control of a trapping voltage U that tunes ω_2, as depicted in Fig. 2(a). We define an analog scale parameter, postponing the internal dynamics of Eq. (3), by $a_2(t) = \omega_2(t)/\omega_2(0)$; see Ref. [10]. In our experiments we ramp $\omega_2/(2\pi)$, down and up, spanning 2.50(1) to 0.50 (1) MHz, where each ramp corresponds to an inflation of

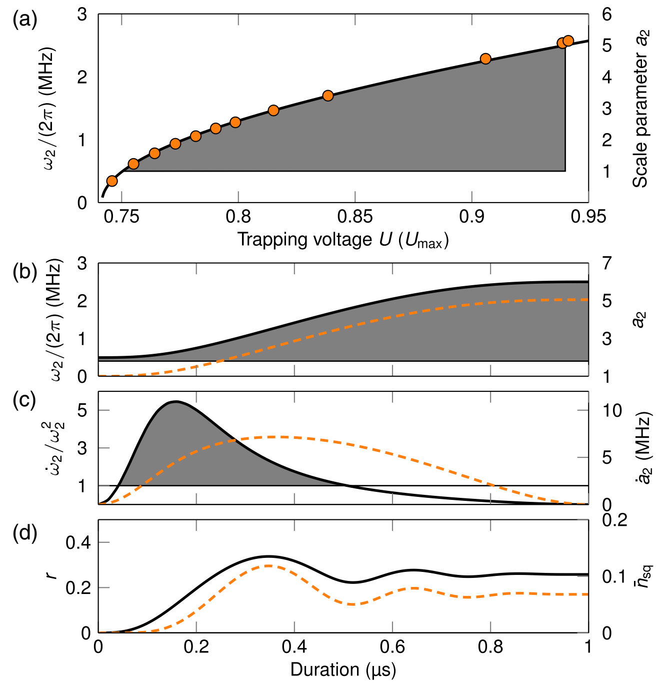

FIG. 2. Accessible experimental parameter regime in our setup. (a) Frequency of the harmonic oscillator mode ω_2 as a function of U; data points depict experimental results (error bars smaller than symbol size), the solid line shows a model fit to the data. The selected frequency range (shaded area) is compared to the analog scale parameter a_2 (right-hand axis), which corresponds to 1.6 e-foldings of space. (b) Real-time switching of ω_2 (solid line) within $t_{\mathrm{ramp}} = 1$ μs and corresponding time evolution of a_2 (dashed line, right-hand axis). (c) The fast variation $\dot{\omega}_2/\omega_2^2$ (solid line) is accompanied by significant time derivatives of the scale parameter $\dot{a}_2$ (dashed line, right-hand axis). (d) Numerical simulation of the quantum dynamics. Quantum fluctuations are inflated and squeezing r (solid line) emerges. This corresponds to the pairwise creation of $\bar{n}_{\mathrm{sq}}$ phonons (dashed line, right-hand axis).

space by about 1.6 e-foldings; see Fig. 2(b). We ramp U within $t_{\mathrm{ramp}} = 1$ μs, yielding a nonadiabatic evolution of the wave function ψ_2, characterized by $\dot{\omega}_2/\omega_2^2 \approx 5$, cf. Fig. 2(c). We note that, in general, deriving adiabaticity criteria and estimating nonadiabatic corrections can be a nontrivial task [55]; see also Refs. [56–61]. Here, we confirm the nonadiabaticity of the implemented dynamics by numerical simulations [10] for a single ramp of ω_2 taking into account the finite ramping duration. The resulting simulated squeezing parameter r is presented in Fig. 2(d) and indicates particle creation with an average particle number $\bar{n}_{\mathrm{sq}} = \sinh^2(r)$. Because of the noninstantaneous switching of ω_2, the Wentzel-Kramers-Brillouin (WKB) phase $\varphi_2(t) = \int_0^t \omega_2(t')dt'$ of ψ_2 evolves significantly during t_{ramp} and leads to an oscillatory accumulation of r and $\bar{n}_{\mathrm{sq}}$.

Experimental results.—In a first experimental realization we cool all modes close to their motional ground state, for the rocking mode ω_2 with a measured, residual thermal

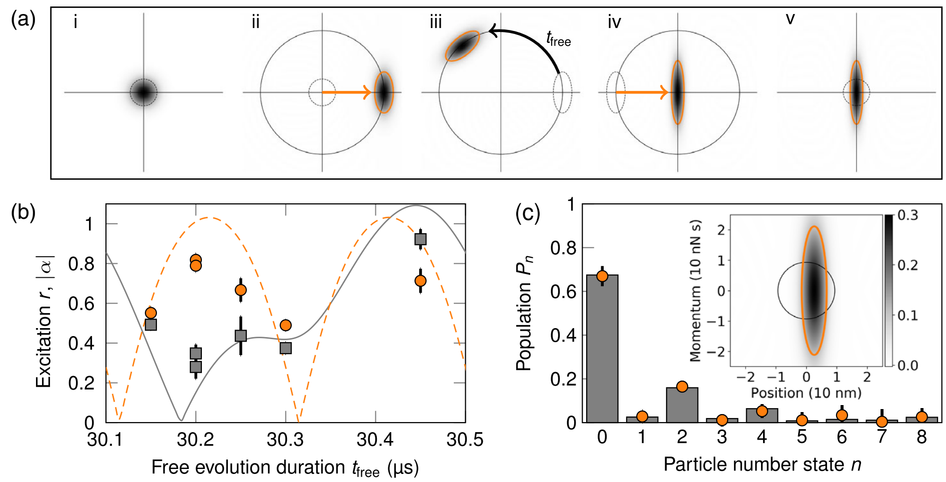

FIG. 3. Experimental results for particle creation and purification of squeezing with an echo sequence consisting of two pulses separated by a free evolution duration $t_{\rm free}$. (a) Phase space illustration of the anticipated echo effect: (i) The initial vacuum state, (ii) squeezed and displaced by the first pulse, (iii) rotates on a circular path in phase space. (iv) After an optimal duration $t_\pi = \pi/\omega_2$ (or odd multiples) the residual coherent displacement is reversed by the second pulse. (v) The final state shows enhanced squeezing excitation. (b) Experimental results for squeezed (circles) and coherent (squares) excitation as a function of $t_{\rm free}$, error bars indicate the standard error of the mean. Numerical simulations (lines) indicate the different oscillation rates of $r(t_{\rm free})$ and $\alpha(t_{\rm free})$; see Ref. [10]. (c) Phonon number distribution (data points, error bars indicate three standard error of the mean) for $t_{\rm free} = 30.2\ \mu$s. From a model fit (bars) we extract $r = 0.83(8)$ and $|\alpha| = 0.29(15)$ and the corresponding Wigner function (inset, phases omitted for clarity). The squeezed state evidences entanglement in the spatial degrees of freedom (d.o.f.) of the two ions, represented by the superposition $|^\bullet{}_\bullet\rangle + |_\bullet{}^\bullet\rangle$; see main text.

$\bar{n}_{\rm th} = 0.03(6)$; see Ref. [10]. Subsequently, we induce dynamics analog to Eq. (3); however, in order to ensure stable preparation and detection conditions, we apply a sequence of two ramps to form one pulse: After ramping down ω_2 and an appropriate duration $t_{\rm hold}$, ω_2 is ramped back up. Finally, we read out the motional state by mapping it onto the electronic state of one ion [10,39], and reconstruct the individual phonon number distribution P_n. To evaluate our results, we decompose the excitation into contributions by r, a coherent displacement α, and a thermal share $\bar{n}_{\rm th}$. We fit the data with a parametrized phonon number distribution $P_n^{\rm par}(r, \alpha, \bar{n}_{\rm th})$ of a Gaussian state, with $r = 0.54(8)$ and $|\alpha| = 0.88(6)$; see Ref. [10]. Since all heating effects remain negligible within experimental sequences, we keep $\bar{n}_{\rm th}$ at the initial value 0.03. We attribute $|\alpha| \neq 0$ to residual static differential stray fields in our setup [10]. Both excitation amplitudes, r and $|\alpha|$, depend on the phase φ_2 accumulated during the pulse; i.e., they are determined by $t_{\rm ramp}$, $t_{\rm hold}$, and $\omega_2(t)$; see Ref. [10].

In order to reverse the coherent displacement, we perform purifying echo sequences consisting of two pulses, separated by a duration of free evolution $t_{\rm free}$. The anticipated echo effect is illustrated by the corresponding phase space dynamics; see Fig. 3(a). The initial thermal state (i) is squeezed and coherently displaced by the first pulse (ii). During $t_{\rm free}$ the state oscillates, i.e., rotates on a circular path in phase space (iii). ψ_2 accumulates two distinct phases corresponding to excitations r and α with frequencies $2\omega_2(t_{\rm free})$ and $\omega_2(t_{\rm free})$, respectively. Ideally, the second pulse is applied after a duration $\pi/\omega_2(t_{\rm free})$ (or odd multiples), where the coherent displacement has evolved from α to $-\alpha$ (iv), and the final state after the second pulse (v) is characterized by $\alpha \approx 0$ and enhanced squeezing [10]. We experimentally realize the two-pulse sequence for variable $t_{\rm free}$, perform the motional state analysis, and depict resulting squeezing and coherent excitations in Fig. 3(b). As indicated by the data, for $t_{\rm free} \approx 30.2\ \mu$s, the squeezing is significantly increased, while the coherent excitation is reduced with respect to the single pulse sequence. Accordingly, in Fig. 3(c) we depict the P_n for $t_{\rm free} = 30.2\ \mu$s. Here the squeezing is directly evidenced by increased populations P_n for even states only, while odd states remain nearly unpopulated, cf. Fig. 1(d). This corresponds to the creation of pairs of phonons. For example, we detect single phonon pairs ($n = 2$) in $P_2 \approx 20\%$ ($\approx$8000 in total) and double phonon pairs ($n = 4$) in $P_4 \approx 5\%$ of our realizations, and the Wigner function indicates a suppression of the variance $(\Delta x)^2$ by $\approx$7.2 dB.

Moreover, features of entanglement of the particle pairs can be witnessed in the spatial degree of freedom (d.o.f.) of our two-ion crystal [49]. To first order, the wave function's squeezing excitation can be visualized in a simplifying way with a term $\propto r(|^\bullet{}_\bullet\rangle + |_\bullet{}^\bullet\rangle)$, resembling a Schrödinger cat state. Here, $|^\bullet{}_\bullet\rangle$ and $|_\bullet{}^\bullet\rangle$ indicate the ions' nonclassical anticorrelation that can be identified with particles with momenta $+\hbar k_2$ and $-\hbar k_2$, respectively. We note that the creation of multiple particle pairs per state ($n = 4, 6, 8$) that we observe evidences the bosonic character of the underlying quantum statistics of phonons [10]. By employing the criterion in Ref. [27] for Gaussian states we quantify the entanglement of formation $E_F \approx 0.41$, significantly increased with respect to the intrinsic vacuum state entanglement $E_F \approx 10^{-5}$; see Ref. [10]. In future studies, we may be able to effectively decouple both ions, while preserving their nonclassical correlation, and individually map the entanglement onto the ions' electronic d.o.f. [25,62].

Conclusion and outlook.—In general, many fundamental predictions rely on the premise that the laws of quantum theory remain valid, not only across very different length and energy scales, but also during drastic changes of the external conditions. Here, we tested this premise in our trapped ion analog. We can continue these studies for generalized analogies involving scalar or vector fields in different dimensions. Depending on the interpretation of the time coordinate (e.g., transformation from proper to conformal time), we may simulate different dynamics of the related scale parameter [10] and investigate its consequences. For this, we can further tune durations and shapes of the ramps of our potential or add parametric driving. Our platform may allow us to study the causal connections of squeezing, pair creation, and entanglement, e.g., in the context of Hawking radiation due to surface gravity, the crossing of cosmic horizons during inflation, and the Sauter-Schwinger effect relating to high-field lasers [10]. Considering the controlled coupling of our closed quantum system to environments and noise fields, we might simulate realistically extended analogs. Further, our method represents a novel tool to create squeezing that can assist in gaining sensitivity for quantum metrology applications, cf. Ref. [63]. In addition, squeezing has also been proposed to substantially enhance effective spin-spin interactions, required for experimental simulations of quantum spin models [64,65], and has recently been used to implement qubits for quantum information processing in the motional states of trapped ions [66]. Finally, our results regarding squeezing and purification from coherent excitations should be considered when implementing (multi-ion) entangling gates on multiplex trap architectures, where rapid changes of potential landscapes are required for scaling towards a universal quantum computer [67,68].

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