Quantum-noise quenching in quantum tweezers

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The efficiency of extracting single atoms or molecules from an ultracold bosonic reservoir is theoretically investigated for a protocol based on lasers, coupling the hyperfine state in which the atoms form a condensate to another stable state, in which the atom experiences a tight potential in the regime of collisional blockade, the quantum tweezers. The transfer efficiency into the single-atom ground state of the tight trap is fundamentally limited by the collective modes of the condensate, which are thermally and dynamically excited and constitute the ultimate noise sources. This quantum noise can be quenched for sufficiently long laser pulses, thereby achieving high efficiencies, and showing that this protocol can be applied for quantum information processing based on tweezer traps for neutral atoms.

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Optical tweezers hold and manipulate microparticles, from molecules to living cells [1–3]. The concept on which they are based find applications down to the level of single atoms [4, 5, 7]. Indeed, the progress in the mechanical manipulation of atoms by means of lasers has allowed, amongst others, to control the position and transport of cold particles loaded in dispersive potentials [5–8]. Such control is at the basis of several protocols for quantum information processing with neutral atoms [9–12]. In this context a relevant issue is the initialization of the quantum register, namely, their preparation in target quantum states of the single atom trap, the quantum tweezers.

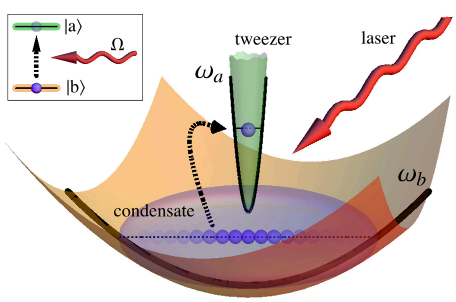

FIG. 1: (Color online) An atom is transferred from a Bose-Einstein condensate (prepared in the electronic state $|b\rangle$) to the ground state of a tweezer trap (realized when the atom is in state $|a\rangle$) by a Raman transition coherently coupling the two states. The protocol is based on spectrally resolving the one-atom ground states of the tweezer trap, which is in the collisional blockade regime [16]. The efficiency of the extraction protocol can be enhanced by quenching the quantum noise due to the condensate excitations.

Proposals for extracting atoms on demand from a quantum reservoir are based either on tunneling and/or dynamical modification of the trapping potential [4, 13–15], or on lasers. In the latter case, transfer of a single atom from a condensate into the ground state of a tweezer trap is achieved by coupling two internal atomic states with different spin-dependent potentials [16–19]. The basic idea is that the ground state of the tweezers trap can be spectrally resolved in the collisional blockade regime, so that the single atom ground state can be coupled on resonance with the ultracold reservoir, while all other states of the quantum tweezers are set significantly out of resonance. This condition however cannot be realized for the collective excitations of the reservoir, the Bose-Einstein condensate. These excitations are present at finite temperature and non-vanishing interactions. They are also created by the dynamics of the extraction process and the collisions between the atom in the tweezers trap and the condensate atoms. They are hence a source of quantum noise, which is inherently due to matter wave fluctuations and which is expected to reduce the efficiency of the protocol.

In this Letter we show that quantum noise due to matter wave fluctuations can be quenched in quantum tweezers for accessible experimental parameters, thereby reaching high fidelities for quantum state preparation of the tweezers. Quenching of noise is achieved by means of a destructive interference between dynamics of different physical origin. Our findings are in agreement with and generalize the predictions in Ref. [18], which have been derived under specific assumptions, and show interesting analogies with quantum noise quenching in quantum optical systems, such as the correlated emission laser [20].

The setup we consider is sketched in Fig. 1. Here, a coherent Raman transition couples two internal, stable states of the atoms, $|b\rangle$ and $|a\rangle$, in which the atoms experience a shallow and a steep confining potential, respectively. Transitions between the two states hence allow for switching between the two confinements. The atoms are identical bosons of mass M and form a Bose-Einstein condensate in state $|b\rangle$, i.e., the quantum reservoir. In absence of perturbations their dynamics is described by Hamiltonian $\mathcal{H}_b = \int \mathrm{d}\mathbf{r}\, \psi_b^\dagger(\mathbf{r}) \left[-\frac{\hbar^2}{2M}\nabla^2 + V_b(\mathbf{r}) + (g_b/2)\psi_b^\dagger(\mathbf{r})\psi_b(\mathbf{r})\right] \psi_b(\mathbf{r})$, with $V_b(\mathbf{r})$ the potential, ψ_b and $\psi_b^\dagger$ the bosonic field operators annihilating and creating an atom in state $|b\rangle$ at position $\mathbf{r}$, and $g_b = 4\pi\hbar^2 a_b/M$ the interaction strength of two-body, s-wave collisions with scattering length a_b. A radiation pulse couples state $|b\rangle$ to $|a\rangle$, in which the atomic center of mass is confined by the steep potential $V_a(\mathbf{r})$ in the regime of collisional blockade, i.e., the quantum tweezers. We denote by $\psi_a(\mathbf{r})$ and $\phi_a(\mathbf{r})$ the field operator and wave function of the atom in the ground state of the tweezer trap, here assumed to be harmonic. The pulse is a standing-wave with wave vector $\mathbf{k}$ and is homogeneous over the volume of the tweezers. It has duration τ, characteristic frequency ω_L and maximum value of the Rabi frequency Ω_0. The Hamiltonian, describing the dynamics due to atom-light coupling, reads $\mathcal{H}_{\rm int}(t) = \mathcal{H}_{\rm r}(t) + \mathcal{H}_{\rm off}(t) + \mathcal{H}_{\rm c}$, where

$$\mathcal{H}_{\rm r}(t) = \frac{\hbar\Omega_0}{2} f(t) \int \mathrm{d}\mathbf{r}\, \cos(\mathbf{k}\cdot\mathbf{r})\psi_a^\dagger(\mathbf{r})\psi_b(\mathbf{r})\mathrm{e}^{-\mathrm{i}\omega_L t} + \mathrm{H.c.} \tag{1}$$

is the Hamiltonian for the resonant coupling between the condensate and the single-atom ground state in the tweezers, with $f(t)$ the temporal shape of the pulse, here assumed to be a step function, while $\mathcal{H}_{\rm off}(t)$ includes the coupling to all other bound states of the tweezers, and $\mathcal{H}_{\rm c} = (g_{ab}/2)\int \mathrm{d}\mathbf{r}\, \psi_b^\dagger(\mathbf{r})\psi_a^\dagger(\mathbf{r})\psi_a(\mathbf{r})\psi_b(\mathbf{r})$ describes s-wave collisions with strength g_{ab} between atoms in $|b\rangle$ and $|a\rangle$. In the regime of collisional blockade the frequency $\omega_{\rm gap} \sim (g_a/2\hbar)\int \mathrm{d}\mathbf{r}\, |\phi_a(\mathbf{r})|^4$ gives the gap between the single- and the two-atom ground state in the tweezers, with g_a the strength of interparticle collisions in $|a\rangle$. For a harmonic potential $V_a(\mathbf{r})$ with frequencies of the order of hundreds of KHz till MHz, $\omega_{\rm gap}$ can reach the order of several to hundreds KHz, and the gap can be spectrally resolved [4, 16–18]. In this regime the laser resonantly couples the condensate with the single-atom ground state of the tweezers with strength $\Omega_{\rm eff} = \Omega_0 \int \mathrm{d}\mathbf{r}\, \cos(\mathbf{k}\cdot\mathbf{r})\phi_a(\mathbf{r})\phi_b(\mathbf{r})$, while the dynamics due to $\mathcal{H}_{\rm off}$ can be neglected. This requires $\omega_{\rm gap}\tau \gg 1$ and $\Omega_{\rm eff} \ll \omega_{\rm gap}$. Correspondingly, the Hilbert space of the tweezers is reduced to the states $|0\rangle_a$ and $|1\rangle_a$, i.e., no atoms and one atom in the tweezers ground state, respectively. It is convenient to define the operators $\sigma = |0\rangle_a\langle 1|$, such that $\psi_a(\mathbf{r}) = \phi_a(\mathbf{r})\sigma$.

We now focus on the effect of the condensate excitations over the efficiency of the extraction dynamics. In the following we assume that the probability of populating non-condensed states during the extraction process is small. For sufficiently low temperature the field operator for the atoms in the condensate can be decomposed into the sum

$$\psi_b(\mathbf{r}) = \phi_b(\mathbf{r}) + \delta\psi_b(\mathbf{r})\,, \tag{2}$$

where $\phi_b(\mathbf{r})$ is the macroscopic wave function of the condensate, satisfying the Gross-Pitaevskii equation $\left(-\frac{\hbar^2\nabla^2}{2M} + V_b(\mathbf{r}) + g_b|\phi_b|^2\right)\phi_b(\mathbf{r}) = \mu\phi_b(\mathbf{r})$ with μ the chemical potential. Operator $\delta\psi_b(\mathbf{r})$ represents the quantum fluctuations about the mean value, and in the Bogoliubov expansion reads $\delta\psi_b(\mathbf{r}) = \sum_{\mathbf{q}}\left[u_{\mathbf{q}}(\mathbf{r})b_{\mathbf{q}} - v_{\mathbf{q}}^*(\mathbf{r})b_{\mathbf{q}}^\dagger\right]$, with $b_{\mathbf{q}}$ and $b_{\mathbf{q}}^\dagger$ the annihilation and creation operators, respectively, of a quasiparticle with frequency $\omega_{\mathbf{q}}$, and $u_{\mathbf{q}}(\mathbf{r})$ and $v_{\mathbf{q}}(\mathbf{r})$ the corresponding wave functions, such that the Hamiltonian for the atoms in state $|b\rangle$ reads $\mathcal{H}_b \simeq \hbar\sum_{\mathbf{q}}\omega_{\mathbf{q}}\, b_{\mathbf{q}}^\dagger b_{\mathbf{q}}$. The total system dynamics is thus mapped to a spin-boson model [18, 21], where the bosonic bath are the Bogoliubov excitations of the condensate and the spin is composed by the tweezers states $|0\rangle_a$ and $|1\rangle_a$, eigenstates of the Pauli matrix σ_z. Using Eq. (2) into Eq. (1), the coupling between condensate and ground state of the tweezers is given by Hamiltonian $\mathcal{H}_{\rm s} = \hbar\Omega_{\rm eff} f(t)\sigma_x/2$, while the coupling involving the Bogoliubov excitations, which emerges from the corresponding terms in $\mathcal{H}_{\rm r} + \mathcal{H}_{\rm c}$, reads

$$\mathcal{H}_{\rm sb} = \frac{\hbar}{2}\sum_{\mathbf{q}}\left(\alpha_{x,\mathbf{q}}(t)\sigma_x + \mathrm{i}\alpha_{y\mathbf{q}}(t)\sigma_y + 2\alpha_{z\mathbf{q}}\sigma_z\right) b_{\mathbf{q}} + \mathrm{H.c.} \tag{3}$$

Here, the first two terms on the right hand side originate from the coupling of the pseudospin to the bosonic bath via the laser, with coupling strengths

$$\alpha_{x\mathbf{q}}(t) = \frac{\Omega_0}{2} f(t)\int \mathrm{d}\mathbf{x}\, \cos(\mathbf{k}\cdot\mathbf{x})\phi_a(\mathbf{x})\left[u_{\mathbf{q}}(\mathbf{x}) - v_{\mathbf{q}}(\mathbf{x})\right] \tag{4}$$

$$\alpha_{y\mathbf{q}}(t) = \frac{\Omega_0}{2} f(t)\int \mathrm{d}\mathbf{x}\, \cos(\mathbf{k}\cdot\mathbf{x})\phi_a(\mathbf{x})\left[u_{\mathbf{q}}(\mathbf{x}) + v_{\mathbf{q}}(\mathbf{x})\right] \tag{5}$$

while the third term is due to collisions between the condensate trap and the tweezers,

$$\alpha_{z\mathbf{q}} = \frac{g_{ab}}{2\hbar}\int \mathrm{d}\mathbf{x}\, |\phi_a(\mathbf{x})|^2\,\phi_b(\mathbf{x})\left[u_{\mathbf{q}}(\mathbf{x}) - v_{\mathbf{q}}(\mathbf{x})\right]. \tag{6}$$

These two kinds of perturbation couple the effective spin with the condensate excitations and may interfere [18].

We evaluate now the fidelity of preparing the tweezers in a target state, which we denote by $|\theta\rangle_a = \cos\theta|0\rangle_a - \mathrm{i}\sin\theta|1\rangle_a$, assuming that initially the tweezers trap is empty and all atoms are in state $|b\rangle$ at finite temperature T. The density matrix of tweezers and condensate excitations at $t = 0$ reads $\rho(0) = |0\rangle_a\langle 0| \otimes \rho_B$, with $\rho_B = \mathrm{e}^{-H_b/K_BT}/Z$ and $Z = \mathrm{Tr}\left\{\mathrm{e}^{-H_b/K_BT}\right\}$. The fidelity can be cast in the form

$$\begin{aligned} P(\theta,\tau) &= \mathrm{Tr}\{|\theta\rangle\langle\theta|\mathrm{e}^{-\mathrm{i}\mathcal{H}_{\rm eff}\tau/\hbar}\rho(0)\mathrm{e}^{\mathrm{i}\mathcal{H}_{\rm eff}\tau/\hbar}\} \\ &= \cos^2(\theta - \Omega_{\rm eff}\tau/2) - g(\theta,\tau) \end{aligned} \tag{7}$$

where $\mathcal{H}_{\rm eff} = \mathcal{H}_s + \mathcal{H}_b + \mathcal{H}_{sb}$ and $g(\theta,\tau)$ is due to the Bogoliubov modes. In second order in the coupling $\mathcal{H}_{\rm sb}$, maximum transfer efficiency is achieved setting $\tau = \tau_0 \equiv 2\theta/\Omega_{\rm eff}$, namely, to the value at which perfect transfer is observed in the ideal dynamics. With this choice, the transfer efficiency reads $P(\theta,\tau_0) = 1 - g(\theta,\tau_0)$, where

$$\begin{aligned} g(\theta,\tau_0) &= \pi^2\sum_{\mathbf{q}}\{A_{1\mathbf{q}}\cos\theta \\ &\quad +(2n_{\mathbf{q}}+1)\left[A_{2\mathbf{q}}\cos 2\theta + A_{3\mathbf{q}} + A_{4\mathbf{q}}\right]\} > 0\,, \end{aligned} \tag{8}$$

with $\bar{n}_{\mathbf{q}} = \mathrm{Tr}\left\{b^\dagger_{\mathbf{q}} b_{\mathbf{q}} \rho_B\right\}$ the mean thermal phonon number of the mode $\omega_{\mathbf{q}}$. The other coefficients take the form

$$\begin{aligned}
A_{1\mathbf{q}} &= -\alpha_{x\mathbf{q}} \delta^{(\tau_0)}\left(\omega_{\mathbf{q}}\right)\left[\alpha_{-\mathbf{q}}\delta^{(\tau_0)}\left(\omega^-_{\mathbf{q}}\right) + \alpha_{+\mathbf{q}}\delta^{(\tau_0)}\left(\omega^+_{\mathbf{q}}\right)\right] , \\
A_{2\mathbf{q}} &= \left(\frac{\alpha_{+\mathbf{q}}\alpha_{-\mathbf{q}}}{2}\right)\delta^{(\tau_0)}\left(\omega^-_{\mathbf{q}}\right)\delta^{(\tau_0)}\left(\omega^+_{\mathbf{q}}\right) , \\
A_{3\mathbf{q}} &= \frac{1}{4}\left[\left(\alpha_{-\mathbf{q}}\delta^{(\tau_0)}\left(\omega^-_{\mathbf{q}}\right)\right)^2 + \left(\alpha_{+\mathbf{q}}\delta^{(\tau_0)}\left(\omega^+_{\mathbf{q}}\right)\right)^2\right] , \\
A_{4\mathbf{q}} &= \alpha^2_{x\mathbf{q}}\delta^{(\tau_0)}\left(\omega_{\mathbf{q}}\right)^2 ,
\end{aligned}$$

with $\delta^{(\tau)}(x) = \sin(x\tau/2)/(\pi x)$, $\alpha_{\pm\mathbf{q}} = \alpha_{y\mathbf{q}} \pm 2\alpha_{z\mathbf{q}}$, and $\omega^\pm_{\mathbf{q}} = \omega_{\mathbf{q}} \pm \Omega_{\mathrm{eff}}$. Rewriting Eq. (8) as a function of these latter coefficients explicitly shows that the quantum noise due to the laser can interfere with the quantum noise due to interspecies collisions at sufficiently long times, for which one can spectrally resolve the frequencies $\omega^\pm_{\mathbf{q}}$.

We now focus on the parameter regime in which quantum noise can be quenched, and study the dependence of the interference condition on the temperature T and on the "Bloch" angle θ of the target state. At zero temperature ($n_{\mathbf{q}} = 0$) the condition on the parameters, for which function (8) is minimal, depends on the angle θ. Maximal quantum noise quenching for the target state $|\theta\rangle_a = \pm|1\rangle_a$ requires that condition

$$\omega_{\mathbf{q}}\alpha_{y\mathbf{q}} - 2\Omega_{\mathrm{eff}}\alpha_{z\mathbf{q}} \sim 0 \tag{9}$$

is fulfilled, for which $g_{\min}(\pi/2, \tau_0) = \pi^2 \sum_{\mathbf{q}} A_{4\mathbf{q}}$. Using Eqs. (4)-(6) in Eq. (9), one finds that the parameter Ω_0 simplifies, such that the interference condition just depends on the trap, density, and interparticle interaction strength. Condition (9) can be further simplified in the limit in which only long-wavelength modes of the condensate are involved, and reduces to the expression $g_{ab} = g_b$, which agrees with the result obtained in Ref. [18] and which was derived for a low energy model. For other target states one finds different conditions on the parameters, and also lower efficiencies: It results, in fact, that for $\theta \neq m\pi/2$, it is not possible to disentangle the condensate excitations from the tweezer state. In particular, it turns out that the target state at $\theta = (2m+1)\pi/4$, with $m = 0, 1, 2, 3$, corresponding to the most non-classical state, equal superposition of one and zero atom in the tweezers, is most sensitive to quantum noise. The dependence of the fidelity on the target state θ is even more enhanced at finite temperatures: In this case the most efficient procedure is the preparation of state $|\theta\rangle_a = \pm|1\rangle_a$, for which the quenching condition is still given by Eq. (9), and $g_{\min}(\pi/2, \tau_0) = \pi^2 \sum_{\mathbf{q}} (2n_{\mathbf{q}}+1) A_{4\mathbf{q}}$. We remark that one general physical consequence is that maximal coherence can be achieved provided that the noise due to collision between condensate and tweezers is significantly different from zero: This noise source can be tuned by means of a Feshbach resonance so to interfere destructively with the laser excitation. We also note that the function $g_{\min}(\pi/2, \tau_0)$ scales with the coupling strength Ω_0^2, as it originates from the excitations due to the laser

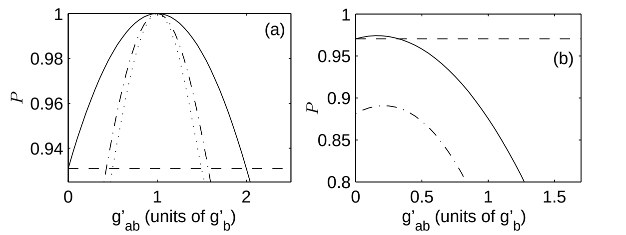

FIG. 2: Efficiency of atomic extraction P from a quasi one-dimensional Bose-Einstein condensate as a function of the interspecies collision strength in 1D, g'_{ab}, in units of g'_b (in 1D) when the target state is (a) $|1\rangle_a$ ($\theta = \pi/2$) and (b) $(|0\rangle - \mathrm{i}\,|1\rangle)_a/\sqrt{2}$ ($\theta = \pi/4$). The curves are found by summing over 500 Bogoliubov modes in Eq. (7) for a condensate of ^{87}Rb atoms in a harmonic trap $V_b(\mathbf{r})$ with axial and radial frequencies $\omega_b = 2\pi \times 200$ Hz and $\omega_\perp = 2\pi \times 0.3$ MHz, density at the center $n_L = 10^8\mathrm{m}^{-1}$, and temperature $T = 0$ (solid line), $T = 50$ nK (dash-dotted line), $T = 100$nK (dotted line in (a)). Here, $\omega_a = 2\pi \times 1$ MHz, $\omega_{\mathrm{gap}} \sim 2\pi \times 0.2$ MHz and $\Omega_{\mathrm{eff}} \sim 2\pi \times 0.45$ kHz. The dashed line gives the efficiency when $g'_{ab} = 0$ and $T = 0$.

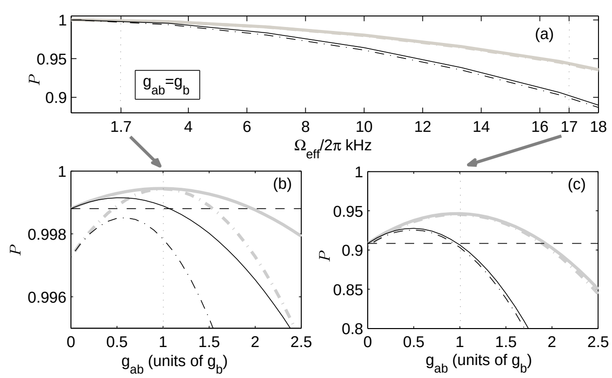

FIG. 3: (a) Efficiency of preparing the tweezers in state $|1\rangle_a$ as a function of Ω_{eff} (for the corresponding value of $\tau = \tau_0$) when the condensate is confined by a spherical harmonic trap with $\omega_b = 2\pi \times 200$ Hz and density at the trap center $n = 2 \times 10^{21} m^{-3}$ (3×10^6 atoms). The curves are found by setting $g_{ab} = g_b$ and $T = 0$ (solid line), and $T = 300$ nK (dash-dotted line), taking $\omega_{j,\ell} = \sqrt{2j^2 + 2j\ell + 3j + \ell}$ for the condensate modes [22, 24]: the gray lines are a sum over $j \in [1, 500]$ and $\ell = 0$, the black lines take into account also the modes with $\ell = 2$ (modes with ℓ odd do not couple to the tweezer ground state). The other parameters are as in Fig. 2. The subplots show P at (b) $\Omega_{\mathrm{eff}} = 2\pi \times 1.7$ kHz and (c) $\Omega_{\mathrm{eff}} = 2\pi \times 17$ kHz as a function of the interparticle collision strength g_{ab} in units of g_b.

coupling which do not interfere with interspecies collisions. This implies that higher efficiency are attained for lower values of Ω_0, and hence for longer transfer pulses.

We now provide some examples with experimental numbers, and consider a condensate of ^{87}Rb atoms in the Thomas-Fermi regime. When evaluating the fidelities we consider only long-wavelength excitations of the condensate, for which the explicit dispersion relations are known for several trap geometries [22]. Figure 2(a) displays the transfer efficiency, Eq. (7), for $\theta = \pi/2$ when the reservoir is a quasi-one dimensional condensate, realized in a highly anisotropic trap. In this case the mean field in Eq. (2) is not strictly valid, but will be used here in order to compare the transfer efficiency for reservoirs of different dimensions. The collisional strength in the condensate is $g_b' = g_b M\omega_\perp/2\pi\hbar$, with $\omega_\perp$ the transverse trap frequency [22, 23]. The interspecies strength g_{ab}' also accounts for the condensate geometry. One clearly observes a maximum at $g_b' \sim g_{ab}'$, where the effects of quantum noise are expected to interfere destructively, and which lies well above the efficiency one would obtain when $g_{ab}' = 0$. By increasing the temperature the condition on the parameters becomes more sensitive to parameter fluctuations. The transfer efficiency is lower when the target state is at $\theta = \pi/4$ (Fig. 2(b)). In this latter case quantum noise quenching improves by a little amount the efficiency of transfer one would find setting $g_{ab}' = 0$. Figure 3(a) displays the efficiency of preparing the state $|1\rangle_a$ by coupling the tweezers with a condensate in a spherical harmonic trap, when the condition $g_{ab} = g_b$ is met, illustrating that quantum noise quenching is optimal when the laser pulse is sufficiently long, so to minimize the effect of the excitations created by the laser coupling which is out of phase with the interspecies collisions. The subplots (b) and (c) highlight the condition on the physical parameters in order to maximize the transfer efficiency, showing that the contribution of excitations at higher angular momentum decreases the transfer efficiency. This latter coupling can be minimized by an accurate design of the setup.

To conclude, the condensate excitations limit the efficiency of preparing quantum tweezers by loading atoms from a condensate, nevertheless their effect can be quenched by means of an interference process emerging from the dynamics induced by laser and particle-particle collisions. This requires sufficiently long transfer pulses, and tuning the various parameters, so to maximize the interference and achieving high fidelities. These concepts can be extended to protocols for creating entangled atoms in two distant tweezers traps coupled to the same condensate, developing on proposals in Refs. [25, 26]. In a more general framework, the dynamics here reported are another example of quantum reservoir engineering [27], where quantum noise may compete in a counterintuitive way in establishing quantum coherence in a physical system. An interesting analogy is here found with the correlated emission laser, where vacuum fluctuations at two different emission frequencies of an atom interfere locking the phase difference between the electric field amplitudes [20]. These phenomena are robust against parameters fluctuations and set the basis to novel paths to quantum technologies.

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[1] A. Ashkin, Phys. Rev. Lett. **24**, 156 (1970); A. Ashkin, *et al.*, Opt. Lett. **11**, 288 (1986).
[2] C. Bustamante, J. Liphardt and F. Ritort, Physics Today **58**, 43 (2005).
[3] W. Singer, *et al.*, Phys. Rev. E **75**, 011916 (2007).
[4] R. B. Diener, *et al.*, Phys. Rev. Lett. **89**, 070401 (2002).
[5] S. Kuhr, *et al.*, Science **293**, 278 (2001); Y. Miroshnychenko, *et al.*, Nature **442**, 151 (2006).
[6] R. Dumke, *et al.*, Phys. Rev. Lett. **89**, 220402 (2002).
[7] J. Beugnon, *et al.*, Nature Phys. **3**, 696 (2007).
[8] P. Böhi, *et al.*, Nature Phys. **5**, 592 (2009).
[9] T. Calarco, *et al.*, Phys. Rev. A **61**, 022304 (2000).
[10] U Dorner, *et al.*, J. Opt. B **7** 341 (2005).
[11] D. Hayes, P. S. Julienne, and I. H. Deutsch, Phys. Rev. Lett. **98**, 070501 (2007).
[12] M. Saffman, T. G. Walker, and K. Moelmer, Rev. Mod. Phys. **82**, 2313 (2010).
[13] A. M. Dudarev, M. G. Raizen, and Q. Niu, Phys. Rev. Lett. **98**, 063001 (2007); C.-S. Chuu, *et al.*, Phys. Rev. Lett. **95**, 260403 (2005).
[14] M. G. Raizen, *et al.*, Phys. Rev. A **80**, 030302 (2009).
[15] C. Kollath, M. Köhl, and T. Giamarchi, Phys. Rev. A **76**, 063602 (2007).
[16] B. Mohring, *et al.*, Phys. Rev. A **71**, 053601 (2005).
[17] P. Rabl, *et al.*, Phys. Rev. Lett. **91**, 110403 (2003).
[18] A. Recati, *et al.*, Phys. Rev. Lett. **94**, 040404 (2005).
[19] M. Bruderer and D. Jaksch, New J. Phys. **8**, 87 (2006).
[20] W. P. Schleich and M. O. Scully, Phys. Rev. A **37**, 1261 (1988); W. Schleich, M. O. Scully, and H.-G. von Garssen, Phys. Rev. A **37**, 3010 (1988).
[21] A. J. Leggett, *et al.*, Rev. Mod. Phys. **59**, 1 (1987)
[22] L. Pitaevskii and S. Stringari, *Bose-Einstein Condensation* (Clarendon Press, Oxford, 2003).
[23] D.S. Petrov, G.V. Shlyapnikov, and J.T.M Walraven, Phys. Rev. Lett. **85**, 3745 (2000).
[24] P. Öhberg, *et al.*, Phys. Rev. A **56**, R3346 (1997).
[25] A. Klein and M. Fleischhauer, Phys. Rev. A **71**, 033605 (2005); see also M. A. Cirone, *et al.*, New J. Phys. **11**, 103055 (2009).
[26] A. Recati, *et al.*, Phys. Rev. A **72**, 023616 (2005).
[27] S. Diehl, *et al.*, Nature Phys. **4**, 878 (2008); F. Verstraete, M. M. Wolf, and J. I. Cirac, Nature Phys. **5**, 633 (2009).