Dissipative Preparation of Spin Squeezed Atomic Ensembles in a Steady State

Citation: Dalla Torre, Emanuele G., Johannes Otterbach, Eugene Demler, Vladan Vuletic, and Mikhail D. Lukin. Dissipative Preparation of Spin Squeezed Atomic Ensembles in a Steady State. Physical Review Letters 110, no. 12 (March 2013).© 2013 American Physical Society.

As Published: http://dx.doi.org/10.1103/PhysRevLett.110.120402

Publisher: American Physical Society

Persistent URL: http://hdl.handle.net/1721.1/81217

Version: Final published version: final published article, as it appeared in a journal, conference proceedings, or other formally published context

Dissipative Preparation of Spin Squeezed Atomic Ensembles in a Steady State

Emanuele G. Dalla Torre,[1,*] Johannes Otterbach,[1,†] Eugene Demler,[1] Vladan Vuletic,[2] and Mikhail D. Lukin[1]

[1]*Physics Department, Harvard University, Cambridge, Massachusetts 02138, USA*

[2]*Physics Department, Massachussetts Institute of Technology, Cambridge, Massachusetts 02139, USA*

(Received 7 September 2012; revised manuscript received 21 December 2012; published 19 March 2013)

We present and analyze a new approach for the generation of atomic spin-squeezed states. Our method involves the collective coupling of an atomic ensemble to a decaying mode of an open optical cavity. We demonstrate the existence of a collective atomic dark state, decoupled from the radiation field. By explicitly constructing this state we find that it can feature spin squeezing bounded only by the Heisenberg limit. We show that such dark states can be deterministically prepared via dissipative means, thus turning dissipation into a resource for entanglement. The scaling of the phase sensitivity taking realistic imperfections into account is discussed.

DOI: 10.1103/PhysRevLett.110.120402 PACS numbers: 42.50.Dv, 03.65.Yz, 42.50.Pq

The realization of spin-squeezed states [1] of atomic ensembles is an important subject in quantum science. Such states play a central role in studies of many-body entanglement [2–4]. In addition, they may lead to practical improvements of state-of-the-art atomic clocks and frequency standards [5–7]. In a spin-squeezed state the atoms are entangled in such a way that the fluctuations of their total spin are smaller than the sum of the fluctuations of the individual atoms. As a rule, such entangled states are extremely fragile and spin squeezing is destroyed due to dissipation or decoherence.

In this Letter we propose a new approach for the realization of spin squeezing using "quantum bath engineering" [8–12] and show how to realize squeezing in the steady state of a dissipative atom-cavity system. In our scheme, the steady state is unique and is reached by the system starting from any initial state, without the need to adiabatically follow a particular path in parameter space. Because the spin squeezing is achieved by optical pumping, our approach can be considerably more robust against noise as compared to existing preparation methods [13–21] that produce short-lived spin-squeezed states, limited by decoherence processes. Moreover, using properly selected atomic transitions allows us to continuously pump the atoms into the desired states and consequently avoid population losses due to scattering into states not participating in the squeezing. A similar approach to continuously entangle two remote atomic ensembles using free space scattering has been experimentally demonstrated in Ref. [22].

The central idea of our work can be understood by considering an ensemble of N atoms interacting with a single radiation mode of an open optical cavity [Fig. 1(a)] and externally driven by a pair of coherent laser fields. The cavity mode and laser fields are tuned to excite a pair of two-photon Raman transitions, each involving one laser field with Rabi frequency $\Omega_\pm$ and a single cavity photon a ($a^\dagger$) [Fig. 1(b)]. Assuming that the coupling of atoms to the cavity mode is uniform, the unitary time evolution of such a cavity-atom system is described by the effective Hamiltonian

$$H_{\rm int} = \frac{g}{\Delta} a^\dagger (\Omega_+ S^+ + \Omega_- S^-) + \text{H.c.}, \tag{1}$$

where S^+ (S^-) is the spin raising (lowering) operator of the collective spin, g is the single atom-field coupling strength, and Δ is the single-photon detuning from the excited states. The coupling $g\Omega_\pm/\Delta$ is therefore an effective Rabi frequency corresponding to off-resonant Raman scattering (cf. Fig. 1). Direct examination of the Hamiltonian (1) shows that it has a dark state $|D\rangle = |\psi_{\rm spin}\rangle|0_{\rm cav}\rangle$, where $|0_{\rm cav}\rangle$ is the cavity vacuum and $|\psi_{\rm spin}\rangle$ satisfies

$$(\Omega_+ S^+ + \Omega_- S^-)|\psi_{\rm spin}\rangle = 0. \tag{2}$$

The interesting features of this state are (i) that it is a pure state not containing any photons and consequently is not affected by the cavity decay, i.e., $|\psi_{\rm spin}\rangle$ is a *dark state* of

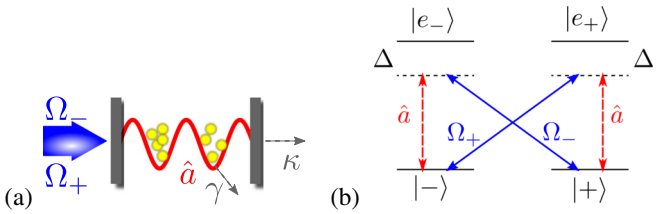

FIG. 1 (color online). (a) Schematic experimental setup of atoms with a spontaneous emission rate γ in a cavity with linewidth κ supporting a single cavity mode a. (b) Effective linkage pattern consisting of two degenerate ground states ($|\pm\rangle$) encoding the effective atomic spin and two excited states $|e_\pm\rangle$. The system is driven via classical control fields $\Omega_\pm$ with large detuning Δ and a single global cavity mode a coupling homogeneously to all atoms. The orientation of the arrows does not necessarily correspond to the actual polarization of the fields in a specific physical realization. We give an example of such an implementation in the Supplemental Material [29].

the cavity decay, and (ii) it describes a highly correlated atomic state containing a high degree of spin squeezing. The latter can be seen by computing the ratio between the spin fluctuations in the x and y direction: $\langle S_x^2\rangle/\langle S_y^2\rangle = \langle (S^+ + S^-)^2\rangle/\langle (S^+ - S^-)^2\rangle = (\Omega_- - \Omega_+)^2/(\Omega_- + \Omega_+)^2$. For any $\Omega_+\Omega_- > 0$ this ratio is smaller than 1, indicating that the dark state is spin squeezed.

To quantify the degree of useful spin squeezing, we use the phase sensitivity introduced by Wineland *et al.* [23],

$$\delta\phi = \frac{\sqrt{\langle S_x^2\rangle}}{|\langle S_z\rangle|}. \tag{3}$$

We first focus on symmetric states with the total spin $S = N/2$, i.e., an initially spin-polarized ensemble, and expand the dark state into eigenstates of S_z as $|\psi_{\text{spin}}\rangle = \sum_m c_m |S = N/2, S_z = -N/2 + m\rangle$. Substituting this into Eq. (2) results in a recursion relation between c_m and c_{m+2}, in close analogy with squeezed light [24],

$$c_{m=2n} = \left(\frac{\Omega_+}{\Omega_-}\right)^n \binom{N/2}{n} \binom{N}{2n}^{-1/2} c_0. \tag{4}$$

Here (:) are the binomial coefficients and c_0 is determined by the normalization condition $\sum_m |c_m|^2 = 1$. The phase sensitivity $\delta\phi$ of $|\psi_{\text{spin}}\rangle$ is shown by the blue solid curve in Fig. 2 as a function of the ratio Ω_+/Ω_- between the control fields. For $\Omega_+/\Omega_- = 0$, the dark state is fully polarized along the z axis, giving a phase sensitivity of $1/\sqrt{N}$ corresponding to the standard quantum limit (SQL). In the opposite limit of $\Omega_+ \to \Omega_-$, the phase sensitivity approaches the Heisenberg limit [16,25] $\delta\phi = 1/\sqrt{N(N/2+1)}$, indicating that the dark state corresponds to an almost maximally squeezed atomic state. A related method to generate spin squeezing in the weak coupling regime, i.e., $\Omega_+/\Omega_- \ll 1$, has been recently proposed in Ref. [26].

We now turn to the preparation of the dark state (2) and study the corresponding quantum dynamics, focusing first on the effect of the cavity decay. To this end, consider an ensemble of N four-state atoms, consisting of two (meta) stable states $|+\rangle$, $|-\rangle$ serving as our effective spin states and two excited states $|e_\pm\rangle$. The atoms are coupled to a

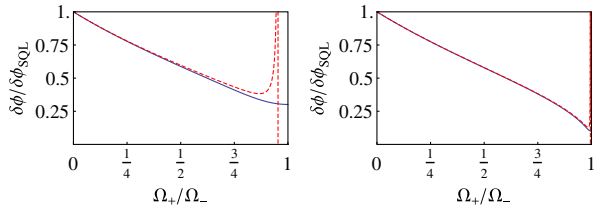

FIG. 2 (color online). Phase sensitivity of the dark state $|\psi_{\text{spin}}\rangle$. Exact numerical solution (solid blue curve) and mean-field approximation (dashed red curve) for $N = 10$ (left) and $N = 100$ (right) atoms. The optimal squeezing is obtained for $\Omega_+ \to \Omega_-$ and approaches the Heisenberg limit $\delta\phi = 1/\sqrt{N(N/2+1)} \sim 1/N$.

single mode of an open cavity (cf. Fig. 1) with resonance frequency ω_a, volume V, and photon escape rate κ driving the transitions $|\pm\rangle \to |e_\pm\rangle$ and strong control fields $\Omega_\pm$ with frequency ω_c driving the transitions $|\mp\rangle \to |e_\pm\rangle$, respectively. The atom-cavity field coupling strength is given by $g = \wp/\hbar\sqrt{\hbar\omega_a/2\varepsilon_0 V}$, where $\wp$ is the dipole matrix element of the respective transition. We assume the control fields and the cavity mode to be far detuned from the resonant transition frequency $\omega_{\text{res}} = \omega_\pm - \omega_{e_\pm}$ between ground state and excited state manifold. This allows us to adiabatically eliminate the excited states, and we obtain the effective Hamiltonian (1) with $S^+ = \sum_j^N \sigma_j^+ = \sum_j^N |+\rangle_{jj}\langle -|$, $S^- = (S^+)^\dagger$, and $|\Delta| = |\omega_{\text{res}} - \omega_{a,c}|$. For simplicity of the calculation we assume here that we can realize two independent Λ schemes according to Fig. 1 and neglect any possible hyperfine structure in the excited states. While this approximation is only true if the detunings from the respective excited state is smaller than the hyperfine splitting in the excited state manifold, Hamiltonian (1) is valid for more general conditions, when Raman processes mediated by multiple virtual states can be added up [27,28]. A more detailed derivation of the Hamiltonian, also including different detunings, can be found in the Supplemental Material [29]. Assuming $\kappa \gg |\Delta|$, we can neglect the occupation of the cavity mode, and using standard techniques [30] we obtain the master equation in Lindblad form after tracing over the cavity modes:

$$\frac{\partial}{\partial t}\rho = -\gamma_{\text{cav}}[\{\mathcal{I}^\dagger \mathcal{I}, \rho\}_+ - 2\mathcal{I}\rho\mathcal{I}^\dagger]. \tag{5}$$

Here $\{\cdot,\cdot\}_+$ denotes the anticommutator, $\gamma_{\text{cav}} = g^2\Omega^2/\Delta^2\kappa$ with $\Omega^2 = \Omega_+^2 + \Omega_-^2$, $\mathcal{I} = \sin\theta S^+ + \cos\theta S^-$ is the Lindblad operator accompanying the emission of a cavity photon, and we defined $\tan\theta = \Omega_-/\Omega_+$.

To compute the phase sensitivity $\delta\phi$, we use (5) to determine the time evolution of the expectation values of the spin operators S_z and S_x^2:

$$\frac{d\langle S_z\rangle}{dt} = -\gamma_{\text{cav}}[\cos^2(\theta)\langle S^+S^-\rangle - \sin^2(\theta)\langle S^-S^+\rangle], \tag{6}$$

$$\frac{d\langle S_x^2\rangle}{dt} = -\gamma_{\text{cav}}(\sin(\theta) - \cos(\theta)) \times \langle(\sin(\theta)S^- + \cos(\theta)S^+)(S_xS_z + S_zS_x) + \text{H.c.}\rangle. \tag{7}$$

Note that the dark state $|\psi_{\text{spin}}\rangle$, defined by Eq. (2), is indeed a steady state of Eqs. (6) and (7). Equations (6) and (7) involve higher moments of the spin, whose time evolution should be computed independently. To avoid this complication and obtain an analytic solution, we linearize the equations of motion around $\langle S_z\rangle \approx -N/2$. Defining the small fluctuations as $\delta S_z = S_z + N/2$ and using the approximation $\langle S_x^2 + S_y^2\rangle = \langle \mathbf{S}^2\rangle - \langle S_z^2\rangle \approx N\langle \delta S_z\rangle + N/2$, we find

$$\frac{d\langle \delta S_z\rangle}{dt} = -\gamma_{\rm cav} N[\cos(2\theta)\langle \delta S_z\rangle - \sin^2(\theta)], \quad (8)$$

$$\frac{d\langle S_x^2\rangle}{dt} = -\gamma_{\rm cav} N\left[\cos(2\theta)\langle S_x^2\rangle + [1-\sin(2\theta)]\frac{1}{4}\right]. \quad (9)$$

The solutions to these equations decay exponentially in time with an effective rate $\gamma_{\rm eff} = \gamma_{\rm cav}\cos(2\theta)N$. From Eqs. (8) and (9) we obtain the phase sensitivity in the steady state as

$$\delta\phi^2 = \frac{N[1-\sin(2\theta)]\cos(2\theta)}{[N\cos(2\theta) - 2\sin^2(\theta)]^2}. \quad (10)$$

In Fig. 2 we compare this approximate expression (dashed red curve) with the exact result obtained from Eq. (4) (solid blue curve). The two curves significantly deviate only for $\Omega_+/\Omega_- \approx 1 - \mathcal{O}(1/N)$, or equivalently $\theta \approx \pi/4 - \mathcal{O}(1/N)$. In this regime the present linearized approximation fails because $\langle \delta S_z\rangle = \sin^2(\theta)/\cos(2\theta) \approx \mathcal{O}(N)$. Note that, approaching the limit $\Omega_+ = \Omega_-$, the steady-state spin squeezing reaches its maximal Heisenberg-limited value. But, at the same time, the effective dark-state pumping rate $\gamma_{\rm eff}$ tends to zero, making the spin squeezing process extremely slow. This renders competing processes, such as e.g., spontaneous Raman scattering of individual atoms, very important, as will be discussed below.

Up to this point, we have restricted ourselves to the manifold of total spin $S = N/2$. This approach is valid if the system is initially prepared in the maximally polarized state $|S = N/2, S_z = N/2\rangle$ because Hamiltonian (1) commutes with the total spin $\mathbf{S}^2$. Consequently, if the system is initially prepared in a state with a different total spin, the final steady state will be different and will in general contain a smaller amount of spin squeezing. This is shown in Fig. 3(a), where we compare the linearized solution with the exact numerical solution of the Lindblad equation (5) for $N = 10$ spins with different initial conditions.

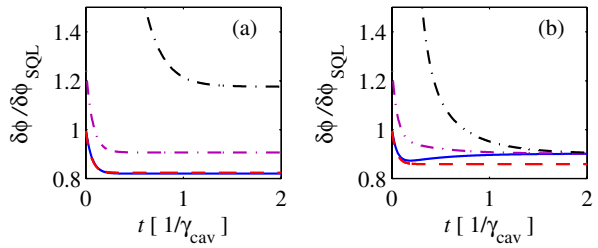

FIG. 3 (color online). Time evolution of the phase sensitivity for $N = 10$ spins and $\Omega_-/\Omega_+ = 0.2$. Numerical solution of the master equation, with fully polarized (solid curves) or randomly generated (dash-dotted curves) initial state, and mean-field approximation (dashed curves). (a) The atoms are coupled only to the cavity, i.e., $\gamma = 0$. (b) Finite scattering rate $\gamma \neq 0$ with single-atom cooperativity $\chi = 1$. Observe that the steady state does not depend on the initial state.

The linearization works well only when the spins are initially polarized, as would be expected.

So far our analysis completely disregarded the effect of spontaneous Raman scattering of the individual atoms into free space. Being a single atom process, it breaks the conservation of the total spin and thus competes with the squeezing process. To investigate the effect of spontaneous Raman scattering we introduce a spontaneous decay rate $\gamma \ll |\Delta|$ to the states $|e_\pm\rangle$. The corresponding master equation is easily obtained using standard techniques and reads

$$\frac{\partial}{\partial t}\rho = -\gamma_{\rm cav}[\{\mathcal{I}^\dagger\mathcal{I},\rho\}_+ - 2\mathcal{I}\rho\mathcal{I}^\dagger] - \gamma_{\rm spont}\sum_{\alpha=\sigma^\pm,\pi}\sum_{j=1}^{N}[\{\mathcal{L}_{j,\alpha}{}^\dagger\mathcal{L}_{j,\alpha},\rho\}_+ - 2\mathcal{L}_{j,\alpha}\rho\mathcal{L}_{j,\alpha}{}^\dagger]. \quad (11)$$

The first term on the right-hand side of Eq. (11) is equivalent to (5), whereas the second term describes the spontaneous Raman scattering of atom j into free space with rates $\gamma_{\rm spont} = \gamma\Omega^2/\Delta^2$. The Lindblad operators are given by $\mathcal{L}_{j,\sigma^\pm} = (\cos\theta\sigma_j^+ + \sin\theta\sigma^-)\sigma^\mp$ and $\mathcal{L}_{j,\pi} = \cos\theta\sigma_j^+ + \sin\theta\sigma^-$. The numerical solution of Eq. (11) is shown in Fig. 3(b) for different initial conditions and $N = 10$ spins. The spontaneous Raman scattering couples different total spin-S manifolds and, as expected, lifts the degeneracy between them. Consequently, it leads to a mixed but nevertheless unique steady state, characterized by a unique density matrix, independent of the initial state, and hence presents a feature rather than a detrimental problem.

Still, Raman scattering leads to a reduction of the achievable phase sensitivity. To quantify this effect we obtain the contribution of the single-spin decay to the time evolution of the collective variables, by using the second row of Eq. (11), leading to

$$\frac{d\langle \delta S_z\rangle}{dt} = -\gamma_{\rm spont}[\langle \delta S_z\rangle - \sin^2(\theta)N], \quad (12)$$

$$\frac{d\langle S_x^2\rangle}{dt} = -2\gamma_{\rm spont}\left(1 - \frac{1}{2}\sin(2\theta)\right)\left[\langle S_x^2\rangle - \frac{1}{4}\right]. \quad (13)$$

Adding Eqs. (12) and (13) to the contribution of the collective coupling to the cavity, Eqs. (8) and (9), we find that in the steady state the competition between the two processes leads to

$$\langle \delta S_z\rangle_\infty = \frac{\sin^2(\theta)(\frac{Ng^2}{\kappa} + N\gamma)}{\frac{Ng^2}{\kappa}\cos(2\theta) + \gamma}, \quad (14)$$

$$\langle S_x^2\rangle_\infty = \frac{1}{4}\frac{\frac{Ng^2}{\kappa}[1-\sin(2\theta)] + 2\gamma[1-\frac{1}{2}\sin(2\theta)]}{\frac{Ng^2}{\kappa}\cos(2\theta) + 2\gamma[1-\frac{1}{2}\sin(2\theta)]}. \quad (15)$$

For a given single-atom cooperativity $\chi = \gamma_{\rm cav}/\gamma_{\rm spont} = g^2/\kappa\gamma$, the optimal value of the phase sensitivity $\delta\phi$ can be obtained by numerically minimizing the corresponding expression with respect to the mixing angle θ while keeping χ fixed. The resulting curve is shown in Fig. 4 for $N = 10^6$. To gain insight into this result, we will now separately investigate the limits of small and large single-atom cooperativity.

If the single-atom cooperativity is small, but $\chi N \sim 1$, the optimal spin-squeezing is obtained for $\theta \ll \pi/4$. Expanding Eqs. (14) and (15) around $\theta \approx 0$ we obtain a phase sensitivity $\delta\phi^2 \approx \frac{1}{N}(1 - N\chi\theta + 4\theta^2)$ with the optimum value obtained for $\theta = N\chi/8$:

$$\delta\phi_{\rm opt} = \frac{1}{\sqrt{N}}\left(1 - \frac{\chi^2 N^2}{32}\right). \tag{16}$$

It should be noted that, due to the collective enhancement by a factor of N, one gains a quadratic improvement in the phase sensitivity even for small cooperativities.

If the cooperativity is larger, i.e., $\chi N \gg 1$, the optimal squeezing is obtained for $\Omega_+ \approx \Omega_-$ and one may naively expect the linearization to fail. However, to render the problem tractable we nevertheless perform the large-N expansion and explicitly check the validity of the obtained results below. To obtain the scaling of the phase sensitivity, we expand Eqs. (14) and (15) around $\theta = \pi/4 - \epsilon$. For small $\epsilon \sim 1/\sqrt{N} \ll 1$ and neglecting terms of order $1/N$, we find $\delta\phi^2 = \frac{1}{N}\frac{1+2N\chi\epsilon^2}{2N\chi\epsilon}$, which has a minimum at $\epsilon = \sqrt{1/2N\chi}$. This leads to the optimal sensitivity

$$\delta\phi = \left(\frac{2}{\chi N^3}\right)^{1/4}. \tag{17}$$

Note that, for the optimal value of θ, the ensemble is still mostly polarized, as can be checked by explicitly computing $\langle \delta S_z \rangle \approx \frac{\chi+1}{4\chi\epsilon} \sim \sqrt{N}$. This observation provides a self-consistency check for our linearization procedure in the limit of a large ensemble, where $\sqrt{N} \ll N$. Equation (17) shows that the achievable phase sensitivity scales as $\delta\phi \sim N^{-3/4}$ and thus offers a significant improvement over the SQL. With respect to other proposals with a similar scaling [16,19,31], our approach deterministically generates spin squeezing in a steady state by optical pumping and can thus potentially be more robust against external perturbations.

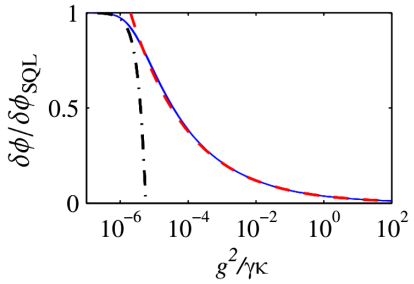

FIG. 4 (color online). Optimal phase sensitivity in the steady state, in the presence of single spin decay for $N = 10^6$. The solid curve is given by the numerical minimization of Eqs. (14) and (15). The dash-dotted and dashed curves are Eqs. (16) and (17), respectively.

A possible experimental realization could be set up in ^{87}Rb using the clock states $|F = 1, m_F = 0\rangle$ and $|F = 2, m_F = 0\rangle$ in the $5S_{1/2}$ ground-state manifold. Circularly σ_+-polarized cavity and control fields couple these states to the states $|F = 1, m_F = +1\rangle$ and $|F = 2, m_F = +1\rangle$ of the $5P_{1/2}$ manifold. In this case the inclusion of the hyperfine structure leads to an effectively larger Raman rate, as the scattering pathways via the two hyperfine states interfere constructively for suitably chosen detunings (see the Supplemental Material [29]). For current experiments [18], we have $N \approx 10^6$ and $\chi \approx 0.1$, predicting a phase sensitivity $\delta\phi/\delta\phi_{\rm SQL} \approx 0.07$, i.e., an improvement of more than 1 order of magnitude with respect to the SQL.

We now discuss two additional effects related to the experimental realization of the proposed model. First, in a multilevel atom (such as the suggested ^{87}Rb), the spontaneous Raman scattering can lead to states outside of the effective four-state system. Because, as shown above, the squeezed steady state is (mixed but) unique and is achieved from any initial state, we can use additional optical pumping fields to repump the atoms into the correct configuration. In the case of ^{87}Rb this can be achieved through linearly polarized fields resonantly driving the $F = 1 \rightarrow F' = 1$, $F = 2 \rightarrow F' = 2$, and $\Delta m_F = 0$ transition. These fields do not couple directly to the clock states $|F = 1, m_F = 0\rangle$ and $|F = 2, m_F = 0\rangle$ and therefore act as a repumping field for the proposed four-state scheme. The repump processes can be easily incorporated in our formalism in the form of additional single-atom scattering channels [29]. These processes effectively increase the single-atom scattering rate and reduce the cooperativity χ. In addition, the use of a large detuning Δ makes it possible to obtain repumping rates that are much larger than the decay rate into the external states and to minimize the loss of atoms into these states. We conclude that our results, and in particular Eq. (17), are valid even in the presence of additional repumped decay channels, provided that the cooperativity rate χ is computed with respect to the total linewidth of the excited states.

Second, in the ideal model (1), we assumed a spatially homogeneous coupling of the atoms to the cavity mode, which can be approximately achieved only in a ring cavity. In a standing-wave cavity the off-resonant Raman coupling $g\Omega_\pm/\Delta$ depends on the position of the atoms and can be either positive or negative, leading to a cancellation of spin squeezing. A solution is to localize the atoms in space and choose an appropriate geometry for the control fields and

the cavity mode. This ensures that all coupling constants have the same sign, as experimentally shown in Ref. [18]. In this case, the symmetry of the collective state is preserved, leading only to a small reduction of the effective spin cooperativity.

To summarize, we presented a method for the deterministic generation of spin-squeezed states in a driven ensemble of effective two-state atoms in a strongly dissipative cavity. The generating process can be understood as optical pumping into a nonequilibrium steady state of the atom that at the same time is a dark state with respect to the cavity decay. Introducing spontaneous Raman scattering we showed that the squeezed steady state is unique and does not depend on the initial state of the system. We discussed the effect of the single atom spontaneous Raman scattering on the achievable phase sensitivity and found that it scales favorably with the single-atom cooperativity, indicating that the present method can be of direct importance for, e.g., optical atomic clocks.

We thank S. Bennett, M. Fleischhauer, M. Hafezi, N. Yao, and P. Zoller for useful discussions and NSF, CUA, DARPA, ARO MURI, and Packard for financial support. J. O. acknowledges support by the Harvard Quantum Optics Center. The numerical solution of the master equations was performed on the Odyssey cluster supported by the FAS Science Division Research Computing Group at Harvard University. Emanuele G. Dalla Torre and Johannes Otterbach contributed equally to this work.

*Corresponding author.
emanuele@physics.harvard.edu
†Corresponding author.
jotterbach@physics.harvard.edu

[1] M. Kitagawa and M. Ueda, Phys. Rev. A **47**, 5138 (1993).
[2] J. K. Korbicz, J. I. Cirac, and M. Lewenstein, Phys. Rev. Lett. **95**, 120502 (2005).
[3] G. Tóth, C. Knapp, O. Guhne, and H. J. Briegel, Phys. Rev. A **79**, 042334 (2009).
[4] O. Gühne and G. Tóth, Phys. Rep. **474**, 1 (2009).
[5] A. M. Childs, J. Preskill, and J. Renes, J. Mod. Opt. **47**, 127 (2000).
[6] K. Hammerer, A. S. Sørensen, and E. S. Polzik, Rev. Mod. Phys. **82**, 1041 (2010).
[7] J. Ma, X. Wang, C. P. Sun, and F. Nori, Phys. Rep. **509**, 89 (2011).
[8] F. Verstraete, M. M. Wolf, and J. I. Cirac, Nat. Phys. **5**, 633 (2009).
[9] S. Diehl, A. Micheli, A. Kantian, B. Kraus, H. P. Büchler, and P. Zoller, Nat. Phys. **4**, 878 (2008).
[10] C. A. Muschik, E. S. Polzik, and J. I. Cirac, Phys. Rev. A **83**, 052312 (2011).
[11] M. J. Kastoryano, F. Reiter, and A. S. Sørensen, Phys. Rev. Lett. **106**, 090502 (2011).
[12] G. Watanabe and H. Mäkelä, Phys. Rev. A **85**, 023604 (2012).
[13] A. Kuzmich, N. P. Bigelow, and L. Mandel, Europhys. Lett. **42**, 481 (1998).
[14] A. André, L.-M. Duan, and M. D. Lukin, Phys. Rev. Lett. **88**, 243602 (2002).
[15] W. Wasilewski, K. Jensen, H. Krauter, J. J. Renema, M. V. Balabas, and E. S. Polzik, Phys. Rev. Lett. **104**, 133601 (2010).
[16] A. André and M. D. Lukin, Phys. Rev. A **65**, 053819 (2002).
[17] A. Louchet-Chauvet, J. Appel, J. J. Renema, D. Oblak, N. Kjaergaard, and E. S. Polzik, New J. Phys. **12**, 065032 (2010).
[18] I. D. Leroux, M. H. Schleier-Smith, and V. Vuletić, Phys. Rev. Lett. **104**, 073602 (2010).
[19] M. H. Schleier-Smith, I. D. Leroux, and V. Vuletić, Phys. Rev. A **81**, 021804 (2010).
[20] Z. Chen, J. G. Bohnet, S. R. Sankar, J. Dai, and J. K. Thompson, Phys. Rev. Lett. **106**, 133601 (2011).
[21] R. J. Sewell, M. Koschorreck, M. Napolitano, B. Dubost, N. Behbood, and M. W. Mitchell, Phys. Rev. Lett. **109**, 253605 (2012).
[22] H. Krauter, C. A. Muschik, K. Jensen, W. Wasilewski, J. M. Petersen, J. I. Cirac, and E. S. Polzik, Phys. Rev. Lett. **107**, 080503 (2011).
[23] D. J. Wineland, J. J. Bollinger, W. M. Itano, F. L. Moore, and D. J. Heinzen, Phys. Rev. A **46**, R6797 (1992).
[24] H. Yuen, Phys. Lett. A **51**, 1 (1975).
[25] Z. Y. Ou, Phys. Rev. A **55**, 2598 (1997).
[26] S. B. Zheng, Phys. Rev. A **86**, 013828 (2012).
[27] H. Xia, S. J. Sharpe, A. J. Merriam, and S. E. Harris, Phys. Rev. A **56**, R3362 (1997).
[28] H. Xia, A. J. Merriam, S. J. Sharpe, G. Y. Yin, and S. E. Harris, Phys. Rev. A **59**, R3190 (1999).
[29] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.110.120402 for a more detailed derivation of the effective Hamiltonian.
[30] C. Gardiner and P. Zoller, *Quantum Noise* (Springer, Heidelberg, 2010).
[31] A. S. Sørensen and K. Mølmer, Phys. Rev. A **66**, 022314 (2002).