

Theoretical method for the generation of a dark two-mode squeezed state of a trapped ion

T. Werlang and C. J. Villas-Boas

Departamento de Física, Universidade Federal de São Carlos, Caixa Postal 676, 13595-905 São Carlos, São Paulo, Brazil

(Received 26 February 2008; published 6 June 2008)

Here we show how to generate a dark two-mode squeezed state of a trapped ion, employing a three-level ion in a V configuration with a strong decay of the excited states. The degree of squeezing can be manipulated by choosing the intensity of the driving fields. Our scheme is robust against the usual dissipation mechanism and could be implemented with present-day technology. The validity of the approximations employed in this work was tested by numerical calculations, which agreed completely with the analytical solutions.

DOI: [10.1103/PhysRevA.77.065801](https://doi.org/10.1103/PhysRevA.77.065801)

PACS number(s): 42.50.Dv, 03.67.Bg, 03.65.Yz

The recent experimental advances in quantum optics, especially in the domain of trapped ions, have allowed fundamental features of quantum mechanics, such as geometric phases [1] and Bell inequalities [2], to be investigated, as well as offering potential applications in quantum computation [3,4] and teleportation processes [5]. With the advent of quantum information theory, the generation of entangled states has become essential for the implementation of quantum communication protocols [3] and to improve our understanding of this nonlocal character of quantum theory [6]. In particular, the two-mode vacuum squeezed (TMVS) state, i.e., the original Einstein-Podolsky-Rosen (EPR) state [7], has attracted much attention because it can show a high degree of entanglement [8] and can be useful for teleportation of continuous variable states [9]. Success in generating the TMVS state has been reported in the running wave domain, with a parametric down conversion process [9]. However, the experimental generation of this state in the cavity quantum electrodynamics (QED) or trapped ion domains has not been achieved so far, mainly due to the sensitivity of quantum states to system-environment interaction. In the cavity QED context, several theoretical schemes with three-level atoms [10–12] or even two-level atoms, where the sideband transition is used [13], have been elaborated for the generation of the TMVS state. Also in the trapped ion domain we find some schemes which allow the generation of this entangled state through the manipulation of laser fields [14]. However, none of the schemes cited above take into account the system-environment interaction, which degrades the quantum states so that, in general, the fidelity of the generated states decays quickly. In this scenario, *reservoir engineering* appears to offer a possible way to round this problem and can generate robust nonclassical states of the radiation field or of the ionic motion. For example, using the atomic decay of the internal levels of a single ion, in Ref. [15] the authors showed how to construct a reservoir able to lead the motion of the ion to a squeezed state asymptotically. Similar schemes have been employed to protect various superpositions of coherent states [16–18], and in Ref. [19] the authors showed how to protect any one-dimensional motional state of an ion against decoherence. Also in this context of reservoir engineering, in Ref. [20] we find an effective master equation that, in the stationary state, filters a specific number state of the vibrational motion of a trapped ion. On the other hand, reservoir engineering for multimode states has been addressed only recently, and there are few theoretical

schemes so far. We can cite, for example, in the trapped ion domain, theoretical schemes for the preparation of a pair coherent state [21] and pair “cat” states (i.e., a superposition of two pair coherent states separated in phase by π) [22], SU(1,1) intelligent states [23], and dark SU(2) states of a trapped ion [24]. Recently, Parkins *et al.* [25] have proposed a scheme for the unconditional generation of a two-mode squeezed state of two separated atomic ensembles. A similar scheme was employed for the generation of the TMVS state for the motion of two ions in different traps [26] or even a single ion in a two-dimensional trap [27] inside an optical cavity. In Ref. [28] it was shown theoretically how to generate this entangled state using an atomic reservoir for a two-mode cavity. In Ref. [29], the authors showed how a beam splitter operation may be produced in a single ion in a two-dimensional trap. After generating a robust squeezed state of a single mode of a trapped ion [15], this effective interaction could be directly employed to generate a TMVS state, but in this case the TMVS state would not be the steady state of the system.

In this Brief Report we report a simple feasible scheme for the unconditional generation of the TMVS state in a single trapped ion. For this purpose we have employed a two-dimensional harmonic motion (on the x and y axes) of the center of mass of a single ion in a V configuration (see Fig. 1). The excited states, $|1\rangle$ and $|2\rangle$, are coupled to the ground state $|0\rangle$ through classical fields (propagating along the x and y axes). When the decay of the excited electronic

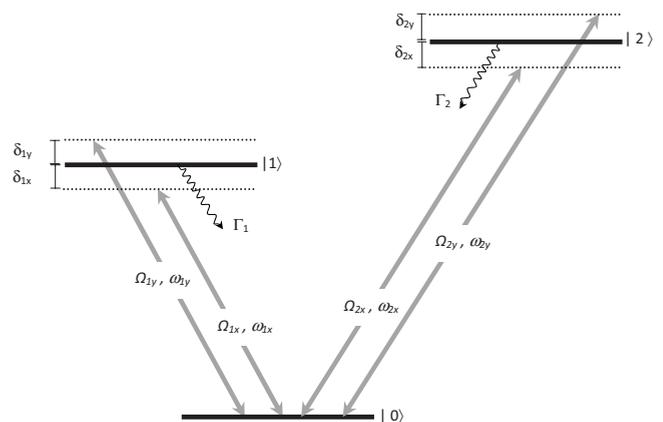


FIG. 1. Atomic levels of the trapped ion. The ground state $|0\rangle$ is coupled to the excited states $|1\rangle$ and $|2\rangle$ through laser fields.

states is stronger than the effective coupling between the vibrational and the internal ionic states, the steady state of the vibrational modes, for convenient choices of the intensity and frequency of the classical fields, turns out to be exactly the TMVS state. Even in the presence of a thermal reservoir for the ionic motion, we show that the generated TMVS state is almost exactly the desired one. Such a scheme, based on reservoir engineering, is robust against dissipative effects of the vibrational modes and could be used to investigate experimentally the entanglement properties of this state. The basic level configuration needed for the implementation of our scheme is sketched in Fig. 1. We consider an ion with mass m in a two-dimensional trap, driven by four classical fields, two of them along the x axis and the other two along the y axis, with complex amplitudes $\Omega_{j\alpha} = |\Omega_{j\alpha}| e^{i\varphi_{j\alpha}}$ ($|\Omega_{j\alpha}|$ being the Rabi frequency and $\varphi_{j\alpha}$ the phase of the classical fields), frequencies $\omega_{j\alpha}$, and wave numbers $k_{j\alpha}$, $j=1,2$ and $\alpha=x,y$. The total Hamiltonian of the system is given by $H = H_0 + V(t)$, with

$$H_0 = \hbar\omega_1\sigma_{11} + \hbar\omega_2\sigma_{22} + \hbar\nu_x a^\dagger a + \hbar\nu_y b^\dagger b, \quad (1a)$$

$$V(t) = \hbar[\Omega_{1x} e^{ik_{1x}x - i\omega_{1x}t} + \Omega_{1y} e^{ik_{1y}y - i\omega_{1y}t}] \sigma_{10} \\ + \hbar[\Omega_{2x} e^{ik_{2x}x - i\omega_{2x}t} + \Omega_{2y} e^{ik_{2y}y - i\omega_{2y}t}] \sigma_{20} + \text{H.c.}, \quad (1b)$$

where ω_1 and ω_2 stand for the atomic transition frequencies between the states $|0\rangle \leftrightarrow |1\rangle$ and $|0\rangle \leftrightarrow |2\rangle$, respectively, $\sigma_{lm} = |l\rangle\langle m|$, $l, m=0,1,2$, are the atomic operators, a (b) and a^\dagger (b^\dagger) are the annihilation and creation operators of the vibrational mode in the x (y) axis, with frequency ν_x (ν_y), and H.c. stands for Hermitian conjugate. In the Lamb-Dicke limit, i.e., $\eta_{j\alpha} \ll 1$, where $\eta_{j\alpha} = k_{j\alpha} \sqrt{\frac{\hbar}{2m\nu_{j\alpha}}}$, $j=1,2$, and $\alpha=x,y$, the above Hamiltonian can be written in the interaction picture as

$$H_I = \hbar\Omega_{1x} e^{i\delta_{1x}t} [1 + i\eta_{1x}(ae^{-i\nu_x t} + a^\dagger e^{i\nu_x t})] \sigma_{10} \\ + \hbar\Omega_{1y} e^{-i\delta_{1y}t} [1 + i\eta_{1y}(be^{-i\nu_y t} + b^\dagger e^{i\nu_y t})] \sigma_{10} \\ + \hbar\Omega_{2x} e^{i\delta_{2x}t} [1 + i\eta_{2x}(ae^{-i\nu_x t} + a^\dagger e^{i\nu_x t})] \sigma_{20} \\ + \hbar\Omega_{2y} e^{-i\delta_{2y}t} [1 + i\eta_{2y}(be^{-i\nu_y t} + b^\dagger e^{i\nu_y t})] \sigma_{20} + \text{H.c.}, \quad (2)$$

with $\delta_{i\alpha} = \omega_{i\alpha} - \omega_i$. Supposing $\delta_{1x} = -\delta_{2x} = \nu_x$, $\delta_{1y} = -\delta_{2y} = \nu_y$, $|\delta_{i\alpha}| \gg |\eta_{i\alpha}\Omega_{i\alpha}|$, and applying a rotating-wave approximation, the effective Hamiltonian becomes

$$H_I = \hbar\lambda_{1x} \left\{ a + \frac{\lambda_{1y}}{\lambda_{1x}} b^\dagger \right\} \sigma_{10} + \hbar\lambda_{2y} \left\{ b + \frac{\lambda_{2x}}{\lambda_{2y}} a^\dagger \right\} \sigma_{20} + \text{H.c.}, \quad (3)$$

where we have defined $\lambda_{j\alpha} \equiv i\eta_{j\alpha}\Omega_{j\alpha}$. As in Ref. [25], we can apply a unitary transformation $\tilde{\rho} = S_{ab}^\dagger(\xi)\rho S_{ab}(\xi)$, with $S_{ab}(\xi) = \exp(\xi^* ab - \xi a^\dagger b^\dagger)$ and $\xi = e^{i\phi}r$, this last being the two-mode squeezing operator (r stands for the squeezing factor and ϕ the angle of squeezing), to obtain the transformed Hamiltonian

$$\tilde{H}_I = S_{ab}^\dagger(\xi) H_I S_{ab}(\xi) = \hbar\tilde{\lambda}_a a \sigma_{10} + \hbar\tilde{\lambda}_b b \sigma_{20} + \text{H.c.}, \quad (4)$$

where $\tilde{\lambda}_a = \lambda_{1x} \cosh(r) - e^{-i\phi} \lambda_{1y} \sinh(r)$, $\tilde{\lambda}_b = \lambda_{2y} \cosh(r) - e^{-i\phi} \lambda_{2x} \sinh(r)$, and we have assumed $\lambda_{1y} \cosh(r) - e^{i\phi} \lambda_{1x} \sinh(r) = \lambda_{2x} \cosh(r) - e^{i\phi} \lambda_{2y} \sinh(r) = 0$, which implies that

$$r = \text{arctanh} \left| \frac{\lambda_{1y}}{\lambda_{1x}} \right| = \text{arctanh} \left| \frac{\lambda_{2x}}{\lambda_{2y}} \right| \quad (5)$$

and

$$\phi = \varphi_{1x} - \varphi_{1y} = -(\varphi_{2x} - \varphi_{2y}). \quad (6)$$

In this way, the squeezing factor r and the squeezing angle ϕ can be manipulated, respectively, by the intensities and the phases $\varphi_{j\alpha}$ of the classical fields. When we take into account the atomic decay of levels $|1\rangle$ and $|2\rangle$, decay rates Γ_1 and Γ_2 , respectively, the dynamics of the system, in the transformed picture, is determined by the master equation

$$\dot{\tilde{\rho}} = -\frac{i}{\hbar} [\tilde{H}_I, \tilde{\rho}] + \mathcal{L}_1 \tilde{\rho} + \mathcal{L}_2 \tilde{\rho}, \quad (7)$$

where $\mathcal{L}_1 \tilde{\rho} = \frac{\Gamma_1}{2} (2\sigma_{01} \tilde{\rho} \sigma_{10} - \sigma_{11} \tilde{\rho} - \tilde{\rho} \sigma_{11})$ and $\mathcal{L}_2 \tilde{\rho} = \frac{\Gamma_2}{2} (2\sigma_{02} \tilde{\rho} \sigma_{20} - \sigma_{22} \tilde{\rho} - \tilde{\rho} \sigma_{22})$. The steady state of Eq. (7) is the vacuum for both modes and $|0\rangle$ for the electronic state. We have assumed a strong decay of both excited electronic states once, as pointed out in Ref. [15], we need two distinct dissipation channels ($\mathcal{L}_1 \tilde{\rho}$ and $\mathcal{L}_2 \tilde{\rho}$) to protect a two-mode quantum state of a trapped ion. (Without this assumption we cannot ensure that the steady state of both modes, in the transformed picture, is the vacuum state.) By applying the reverse unitary transformation, it is readily shown that the steady state of this system is

$$\rho(t \rightarrow \infty) = S_{ab}(\xi) \tilde{\rho} S_{ab}^\dagger(\xi) = S_{ab}(\xi) |0,0\rangle\langle 0,0| S_{ab}^\dagger(\xi) \otimes |0\rangle\langle 0|, \quad (8)$$

which is a pure state for the vibrational modes a and b , i.e., exactly the two-mode squeezed vacuum state $|\Psi\rangle = \sum_n \tanh^n(r) / \cosh(r) |n,n\rangle_{ab}$. The degree of squeezing r is determined by the amplitudes of the classical fields $\Omega_{j\alpha}$ since $\tanh(r) = \left| \frac{\lambda_{1y}}{\lambda_{1x}} \right| = \left| \frac{\lambda_{2x}}{\lambda_{2y}} \right|$ and $\lambda_{j\alpha} \equiv i\eta_{j\alpha}\Omega_{j\alpha}$. This steady state does not depend on the initial electronic or motional state of the ion. Thus the ion does not have to be cooled to the fundamental state in order to prepare such a state. Also, as the entangled state is generated through the engineered reservoir, there is no requirement for a precisely timed interaction between the ion and the laser fields and the degree of entanglement (r) is determined only by the ratio of the amplitudes of the classical fields [see Eq. (5)]. In this scheme, the TMVS state is generated when the system reaches the steady state. As pointed out in Refs. [25,26], the time needed for the system to reach the steady state is defined by the atomic decay rate Γ . For $|\tilde{\lambda}_a| \sim |\tilde{\lambda}_b| > \Gamma$, this time will be of the order of a few times $1/\Gamma$. In Ref. [27] Hua-Tang *et al.* showed how to generate the TMVS state in a single ion in a two-dimensional trap inside a nonideal optical cavity. Different from our scheme, where the required dissipation channels are played by the decay of the excited electronic levels, in Ref.

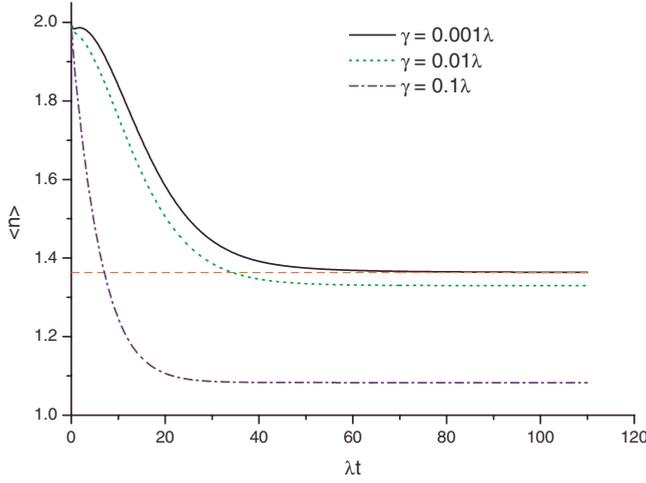


FIG. 2. (Color online) The time evolution of the mean number of quanta, $\langle a^\dagger a \rangle$, of the vibrational mode x for $\Gamma=10\lambda$, $\bar{n}_{th}=0.5$, $r=1$, and three values of the decay rate of the vibrational modes: $\gamma=0.001\lambda$ (solid line), $\gamma=0.01\lambda$ (dotted line), and $\gamma=0.1\lambda$ (dashed-dotted line). The dashed line (straight line) represents the expected value.

[27] the required dissipation channel is played by the decay of the cavity mode.

To check our result we solve numerically the master equation for our system in the interaction picture,

$$\dot{\rho} = -\frac{i}{\hbar}[H_I, \rho] + \mathcal{L}_1\rho + \mathcal{L}_2\rho + \mathcal{L}_{ab}\rho, \quad (9)$$

where H_I is given by Eq. (3). In this equation we have introduced the Liouvillian $\mathcal{L}_{ab}\rho$ which describes the action of the thermal reservoir on both atomic motions, i.e.,

$$\mathcal{L}_{ab}\rho = \sum_{\alpha=a,b} \left\{ \frac{(\bar{n}_{th}+1)\gamma_\alpha}{2} (2\alpha\rho\alpha^\dagger - \alpha^\dagger\alpha\rho - \rho\alpha^\dagger\alpha) + \frac{\bar{n}_{th}\gamma_\alpha}{2} (2\alpha^\dagger\rho\alpha - \alpha\alpha^\dagger\rho - \rho\alpha\alpha^\dagger) \right\}, \quad (10)$$

\bar{n}_{th} being the mean number of quanta of the thermal reservoir and γ_a (γ_b) the decay rate of the vibrational mode x (y). We start with the ion in the internal ground state $|0\rangle$ and both modes in the thermal state, $\rho_{ab}(0) = \rho_a \otimes \rho_b$, with $\rho_a = \rho_b = \sum_{n=0}^{\infty} \frac{\bar{n}^n}{(1+\bar{n})^{n+1}} |n\rangle\langle n|$, \bar{n} being the initial mean number of quanta for each mode. To solve numerically the master equation (9) we adopt the same coupling, $\lambda_{1x} = \lambda_{2y} = \lambda$, $\lambda_{1y} = \lambda_{2x} = \lambda \tanh(r)$, and the same decay rate for both excited electronic states, $\Gamma_1 = \Gamma_2 = \Gamma$, and the same decay rate for the vibrational modes, $\gamma_a = \gamma_b = \gamma$. In Fig. 2, we have plotted the mean number of quanta of mode a against time (the evolution of mode b is identical) for different values of γ and for $\bar{n}=2$, $\bar{n}_{th}=0.5$, $\Gamma=10$, and $r=1$ [which implies $\tanh(r) = \left| \frac{\lambda_{1y}}{\lambda_{1x}} \right| = \left| \frac{\lambda_{2x}}{\lambda_{2y}} \right| \approx 0.76$]. For an ideal two-mode vacuum squeezed state, the mean number of quanta for each mode is $\langle n_{a,b} \rangle = \frac{\tanh(r)^2}{1 - \tanh(r)^2}$, which for $r=1$ gives $\langle n_{a,b} \rangle \approx 1.4$. We can see in

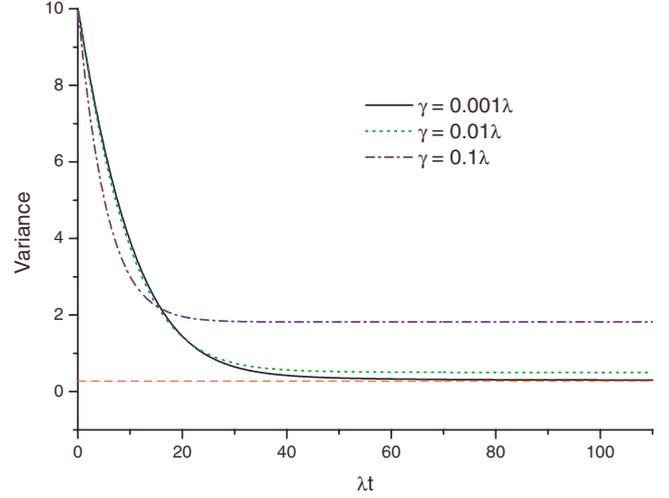


FIG. 3. (Color online) The time evolution of the total variance $\langle (\Delta\hat{u})^2 + (\Delta\hat{v})^2 \rangle$ for the same parameters used in Fig. 2.

Fig. 2 that, for $\gamma=0.001\lambda$ and $\gamma=0.01\lambda$, a mean number of quanta close to this value is reached asymptotically, but this is not so for $\gamma=0.1\lambda$, because of the competition between the engineered and thermal reservoirs. Another parameter we have used to analyze the fidelity of the generated state is the total variance $\langle (\Delta\hat{u})^2 + (\Delta\hat{v})^2 \rangle$ of a pair of EPR-like operators $\hat{u} = |\varepsilon| \hat{x}_a + \frac{1}{\varepsilon} \hat{x}_b$ and $\hat{v} = |\varepsilon| \hat{p}_a - \frac{1}{\varepsilon} \hat{p}_b$ [8], with $\hat{x}_\alpha = (\hat{\alpha} + \hat{\alpha}^\dagger) / \sqrt{2}$ and $\hat{p}_\alpha = -i(\hat{\alpha} - \hat{\alpha}^\dagger) / \sqrt{2}$, $\alpha = a, b$. According to Ref. [8], a two-mode Gaussian state is entangled if and only if $\langle (\Delta\hat{u})^2 + (\Delta\hat{v})^2 \rangle < \varepsilon^2 + 1/\varepsilon^2$. For $\varepsilon=1$ and an ideal two-mode vacuum squeezed state, the total variance is $\langle (\Delta\hat{u})^2 + (\Delta\hat{v})^2 \rangle = 2e^{-2r}$, which for $r=1$ gives us $\langle (\Delta\hat{u})^2 + (\Delta\hat{v})^2 \rangle \approx 0.27$. As we can see in Fig. 3, this value is reached and approximately reached for $\gamma=0.001\lambda$ and $\gamma=0.01\lambda$, respectively. Again, for $\gamma=0.1\lambda$, the action of the thermal reservoir does not allow the ideal generation of the two-mode entangled state. Instead of applying unitary transformations to the density matrix ρ , which led to Eq. (4), and thus making it easy to find the steady state, we could have proceeded by looking for an engineered Liouvillian for the engineered reservoir, as in Ref. [19]. For an atomic decay rate Γ much stronger than the effective coupling λ and the decay rate of the vibrational modes γ , the effective decay rate for the engineered reservoir is given by $\Gamma_{eng} = 4\lambda^2/\Gamma$. In our numerical solution of the master equation (9) we have used $\Gamma=10\lambda$, which results in $\Gamma_{eng} = 0.4\lambda$, that is, of the same order of magnitude as $\gamma = 0.1\lambda$. Then, for this value of γ , the influence on the generated TMVS state of the natural reservoir is almost the same as that of the engineered reservoir. Hence to minimize the influence of the natural reservoirs, we must have $\Gamma_{eng} = 4\lambda^2/\Gamma \gg \gamma$. For example, for $\eta=0.1$ (Lamb-Dicke limit), $\Omega_{i\alpha} \sim 1$ MHz and $\Gamma \sim 1$ MHz, which can easily be achieved with current technology, we have $|\lambda| \sim 0.1$ MHz and $\Gamma_{eng} = 40$ KHz, which is much stronger than $\gamma \sim 2$ KHz, found in present-day experiments. (The chosen values above also satisfy the requirements for the approximations employed to obtain the effective Hamiltonian: for $\nu_x \sim \nu_y \sim 30$ MHz $\gg |\eta\Omega| \sim 0.1$ MHz and $\Gamma \sim 1$ MHz.)

Summarizing, we have presented a simple scheme to prepare a two-mode vacuum squeezed state for the two-dimensional (2D) motion of a trapped ion via the generation of an artificial reservoir. Our scheme is robust against the usual mechanism of dissipation and could be implemented with the present-day technology and we hope it could be employed to test experimentally the entanglement properties of Gaussian states. The approximations employed in this work were validated by numerical calculations, which

showed complete agreement with the analytical solutions. To prove the engineered two-mode state a tomographic method could be employed that enables the Wigner function of the entangled state to be reconstructed [30].

We wish to acknowledge the support of the Brazilian agencies CNPq, FAPESP (under Contract No. 2005/04105-5), and Brazilian Millennium Institute for Quantum Information.

-
- [1] M. V. Berry, Proc. R. Soc. London, Ser. A **392**, 45 (1984); D. Leibfried, B. DeMarco, V. Meyer, D. Lucas, M. Barrett, J. Britton, W. M. Itano, B. Jelenkovic, C. Langer, T. Rosenband, and D. J. Wineland, Nature (London) **422**, 412 (2003).
- [2] J. S. Bell, Physics (Long Island City, N.Y.) **1**, 195 (1965); M. A. Rowe, D. Kielpinski, V. Meyer, C. A. Sackett, W. M. Itano, C. Monroe, and D. J. Wineland, Nature (London) **409**, 791 (2001); B. B. Blinov, D. L. Moehring, L. M. Duan, and C. Monroe, *ibid.* **428**, 153 (2004).
- [3] M. A. Nielsen and I. L. Chuang, *Quantum Computation and Quantum Information* (Cambridge University Press, Cambridge, England, 2000).
- [4] J. I. Cirac and P. Zoller, Phys. Rev. Lett. **74**, 4091 (1995); Nature (London) **404**, 579 (2000); D. Kielpinski, C. Monroe, and D. J. Wineland, *ibid.* **417**, 709 (2002); S. Gulde, M. Riebe, G. P. T. Lancaster, C. Becher, J. Eschner, H. Häffner, F. Schmidt-Kaler, I. L. Chuang, and R. Blatt, *ibid.* **421**, 48 (2003).
- [5] M. Riebe, H. Häffner, C. F. Roos, W. Hansel, J. Benhelm, G. P. T. Lancaster, T. W. Korber, C. Becher, F. Schmidt-Kaler, D. F. V. James, and R. Blatt, Nature (London) **429**, 734 (2004); M. D. Barrett, J. Chiaverini, T. Schaetz, J. Britton, W. M. Itano, J. D. Jost, E. Knill, C. Langer, D. Leibfried, R. Ozeri, and D. J. Wineland, *ibid.* **429**, 737 (2004).
- [6] Q. A. Turchette, C. S. Wood, B. E. King, C. J. Myatt, D. Leibfried, W. M. Itano, C. Monroe, and D. J. Wineland, Phys. Rev. Lett. **81**, 3631 (1998).
- [7] A. Einstein, B. Podolsky, and N. Rosen, Phys. Rev. **47**, 777 (1935).
- [8] L.-M. Duan, G. Giedke, J. I. Cirac, and P. Zoller, Phys. Rev. Lett. **84**, 2722 (2000).
- [9] A. Furusawa, J. L. Sørensen, S. L. Braunstein, C. A. Fuchs, H. J. Kimble, and E. S. Polzik, Science **282**, 706 (1998).
- [10] C. J. Villas-Bôas, N. G. de Almeida, R. M. Serra, and M. H. Y. Moussa, Phys. Rev. A **68**, 061801(R) (2003); R. M. Serra, C. J. Villas-Bôas, N. G. de Almeida, and M. H. Y. Moussa, *ibid.* **71**, 045802 (2005).
- [11] C. J. Villas-Boas and M. H. Y. Moussa, Eur. Phys. J. D **32**, 147 (2005).
- [12] R. Guzmán, J. C. Retamal, E. Solano, and N. Zagury, Phys. Rev. Lett. **96**, 010502 (2006).
- [13] F. O. Prado, N. G. de Almeida, M. H. Y. Moussa, and C. J. Villas-Boas, Phys. Rev. A **73**, 043803 (2006).
- [14] S.-B. Zheng and G.-C. Guo, Quantum Semiclass. Opt. **10**, 441 (1998); X. B. Zou, K. Pahlke, and W. Mathis, Phys. Rev. A **65**, 064303 (2002); H. S. Zeng, L. M. Kuang, and K. L. Gao, Phys. Lett. A **300**, 427 (2002).
- [15] J. I. Cirac, A. S. Parkins, R. Blatt, and P. Zoller, Phys. Rev. Lett. **70**, 556 (1993); J. F. Poyatos, J. I. Cirac, and P. Zoller, *ibid.* **77**, 4728 (1996).
- [16] R. L. de Matos Filho and W. Vogel, Phys. Rev. Lett. **76**, 608 (1996).
- [17] R. L. de Matos Filho and W. Vogel, Phys. Rev. A **54**, 4560 (1996).
- [18] S.-C. Gou, J. Steinbach, and P. L. Knight, Phys. Rev. A **55**, 3719 (1997).
- [19] A. R. R. Carvalho, P. Milman, R. L. de Matos Filho, and L. Davidovich, Phys. Rev. Lett. **86**, 4988 (2001).
- [20] H. Moya-Cessa and P. Tombesi, Phys. Rev. A **61**, 025401 (2000).
- [21] S.-C. Gou, J. Steinbach, and P. L. Knight, Phys. Rev. A **54**, R1014 (1996).
- [22] S.-C. Gou, J. Steinbach, and P. L. Knight, Phys. Rev. A **54**, 4315 (1996).
- [23] C. C. Gerry, S.-C. Gou, and J. Steinbach, Phys. Rev. A **55**, 630 (1997).
- [24] Z. Kis, W. Vogel, L. Davidovich, and N. Zagury, Phys. Rev. A **63**, 053410 (2001).
- [25] A. S. Parkins, E. Solano, and J. I. Cirac, Phys. Rev. Lett. **96**, 053602 (2006).
- [26] G. X. Li, H. T. Tan, S. P. Wu, and G. M. Huang, Phys. Rev. A **74**, 025801 (2006).
- [27] T. Hua-Tang, R. Tang, and G.-X. Li, Chin. Phys. Lett. **25**, 945 (2008).
- [28] S. Pielawa, G. Morigi, D. Vitali, and L. Davidovich, Phys. Rev. Lett. **98**, 240401 (2007).
- [29] S.-C. Gou and P. L. Knight, Phys. Rev. A **54**, 1682 (1996).
- [30] S. Wallentowitz, R. L. de Matos, S. C. Gou, and W. Vogel, Eur. Phys. J. D **6**, 397 (1999).