

## Quantum-state transfer from light to macroscopic oscillators

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(Received 3 October 2002; revised manuscript received 4 March 2003; published 9 July 2003)

We propose a scheme for transferring quantum states from the propagating light fields to macroscopic, collective vibrational degree of freedom of a massive mirror by exploiting radiation pressure effects. This scheme may prepare an Einstein-Podolsky-Rosen state in position and momentum of a pair of distantly separated movable mirrors by utilizing the entangled light fields produced from a nondegenerate optical parametric amplifier.

DOI: 10.1103/PhysRevA.68.013808

PACS number(s): 42.50.Ct, 03.67.Hk, 42.50.Vk

### I. INTRODUCTION

Quantum state exchange between light and matter is an important ingredient for the future quantum information networks. Photons are the fastest and simplest carriers of quantum information for transmission, but, in general, it is difficult to localize and store photons, so usually one prefers choosing matter as quantum memory elements. The interface between quantum carrier and memory is a key component in the realization of scalable quantum networks [1,2]. Several proposals have addressed the transfer of quantum states from photons to matter. A simplest approach to a quantum memory is to absorb photons with unit probability in an optically thick ensemble of atoms. Nevertheless, it has been shown that such absorption of light leads to a partial mapping of its quantum properties to atomic ensembles [3,4]. The important one of the light storage techniques is the stimulated Raman adiabatic passage [5,6]. The application of this technique in single-atom cavity QED systems has led to a number of important proposals for qubit transfer between atoms and photons as well as for quantum-logic gates [7–9]. Mapping quantum states to collective atomic spin systems [10] and quantized vibrational states of trapped atoms [11,12] by means of the stimulated Raman adiabatic passage technology have been discussed. An alternative approach to mapping a quantum state of light onto an atomic state via electromagnetically induced transparency (EIT) has been also proposed [13,14]. Recent experimental achievements employing cold Na atoms [15] and warm Rb vapor [16,17], which demonstrated the “storage” of a classical light pulse in an atomic ensemble and reversible reduction of the light pulse group velocity to zero, have made important progress in this direction.

Over the last 10 years a lot of theoretical attention has been devoted to optomechanical systems. It has been pointed out that nonclassical quantum states of light can be generated by optomechanical coupling in a cavity with a movable mirror [18–24]. Furthermore, due to recent technological developments in optomechanics, this area is now becoming experimentally accessible [25–27]. Recently, entangling

macroscopic oscillators of movable mirrors in a cavity [28] and teleporting a continuous variable state of a light field onto the vibrational state of a mirror [29] by exploiting the radiation pressure have been presented. In this papers, we propose a scheme of transferring quantum states from the propagating light field to the vibrational state of a movable mirror by exploiting radiation pressure effects. The entanglement of the vibrational state of a movable mirror may be achieved through the transfer of entanglement of quantum-correlated light fields produced from a nondegenerate parametric amplifier (NOPA). Moreover, the proposed scheme might be developed to become important resources used in quantum networks; especially it offer us possibility to obtain the Einstein-Podolsky-Rosen (EPR) states of macroscopic oscillators and teleport the center-of-mass wave function of a massive object [30].

### II. MODEL

In the discussed system, the coherently driven optical cavity consists of a moving mirror ( $M2$ ) and an input and output coupler ( $M1$ ) as shown in Fig. 1. For simplicity, we shall describe the mirror as a single quantum-mechanical harmonic oscillator with mass  $m$  and frequency  $\nu_1$ . In fact, the mirror motion is a result of the excitation of many vibrational modes. The description of the mirror as a single oscillator is, however, a good approximation when frequencies are limited to a bandwidth including a single mechanical resonance. The optomechanical coupling between the mirror and the cavity field is realized by the radiation pressure and one has the following Hamiltonian [28,29]:

$$H = \hbar \omega_a a_1^\dagger a_1 + \hbar \nu_1 (b_1^\dagger b_1 + \frac{1}{2}) + \hbar G a_1^\dagger a_1 x_1 + i \hbar \sqrt{2} \gamma (a_1^{in} e^{-i\omega_{a0} t} a_1^\dagger - a_1^{in*} e^{i\omega_{a0} t} a_1), \quad (1)$$

where  $a_1$  and  $b_1$  are annihilation operators of the cavity field with optical frequency  $\omega_a$  and the quantized vibrational mode of a movable mirror with harmonic-oscillation frequency  $\nu_1$ , respectively. Moreover,  $x_1 = b_1^\dagger + b_1$  is the dimensionless position operator of the movable mirror, and  $G$  is the optomechanical coupling constant (proportional to  $m^{-1/2}$ ,  $m$  is mass of the oscillator). The bright input light field  $a_1^{in}$  can be

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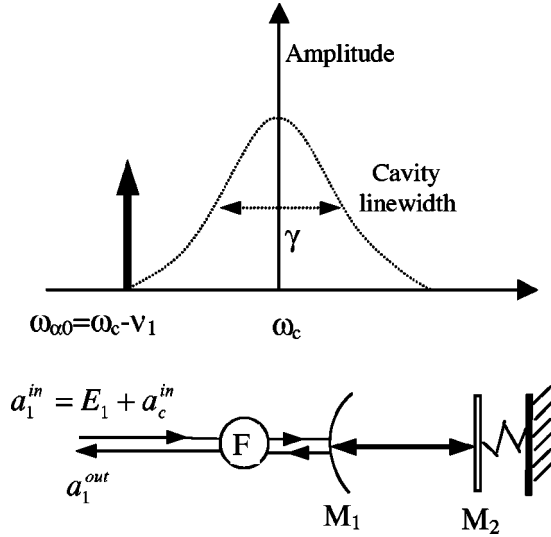


FIG. 1. Schematic description of the discussed system.  $M2$  is the oscillating mirror.  $\Delta = \nu_1$  is the red detuning between the intense driving light and the cavity field. Faraday isolator ( $F$ ) facilitates a unidirectional coupling.

decomposed to an intense carrier  $a_1^{in}$  oscillating at the center frequency  $\omega_{a0}$  with an average amplitude ( $E_1^{in}$ ) which equals the amplitude of its steady-state field, and noise sidebands  $a_{\pm}^{in}(\Lambda)$  surrounding the center frequency and oscillating at frequency  $\omega_{a0} \pm \Lambda$  with zero average amplitude [31]

$$\langle a_1^{in}(\Lambda=0) \rangle = E_1^{in}; \quad \langle a_1^{in}(\Lambda \neq 0) \rangle = 0. \quad (2)$$

Then the Hamiltonian gives rise to nonlinear Langevin equations whose linearization around the steady state leads to

$$\begin{aligned} \dot{a}_1 &= -i\Delta a_1 - iGE_1 x_1 - \gamma a_1 + \sqrt{2\gamma} a_1^{in}, \\ \dot{b}_1 &= -i\nu_1 b_1 - iG(E_1^* a_1 + E_1 a_1^\dagger) - \Gamma b_1 + \xi_1, \end{aligned} \quad (3)$$

and the semiclassical steady values are

$$\begin{aligned} E_1 &= \langle a_1 \rangle = \sqrt{2\gamma} E_1^{in} / (\gamma + i\Delta), \\ \langle b_{1ss} \rangle &= G|E_1|^2 / (\Gamma + i\nu_1), \end{aligned} \quad (4)$$

where  $\Delta = \omega_{a0} - \omega_c + G\langle b_{1ss} \rangle$ . The resonant frequency and the decay rate of the cavity field are  $\omega_c$  and  $\gamma$ , respectively. Furthermore,  $\Gamma$  and  $\xi_1$  are the mechanical damping rate of the mirror Brownian motion and the noise operator of the quantum Brownian motion of the mirror. The general description of quantum Brownian motion of the mirror is presented at any temperatures in Ref. [32]. Assuming that the movable mirror is a miniature high-frequency mechanical oscillator and operates in ultralow temperatures, we consider the noise operator of the mirror as Gaussian white-noise source. Thus the noise correlation of  $\xi_1$  is

$$\begin{aligned} \langle \xi_1(t) \xi_1(t') \rangle &= \langle \xi_1^\dagger(t) \xi_1(t') \rangle = \langle \xi_1^\dagger(t) \xi_1^\dagger(t') \rangle = 0, \\ \langle \xi_1(t) \xi_1^\dagger(t') \rangle &= 2\Gamma(1 + 2n_T) \delta(t - t'), \end{aligned} \quad (5)$$

where  $n_T = \coth(\hbar\nu_1/2k_B T)$  is the mean thermal photon number with  $k_B$  being the Boltzmann constant and  $T$  being the equilibrium temperature. We choose the negatively detuning  $\Delta = \nu_1$  by adjusting the frequency of driving beam  $E_1$ . Assuming that  $\nu_1 \gg \gamma$ ,  $|GE_1| \gg \Gamma$ , Eq. (3) is simplified in rotating frame by neglecting fast rotating terms

$$\begin{aligned} \dot{\tilde{a}}_1 &= -iGEe^{-i\theta_1} \tilde{b}_1 - \gamma \tilde{a}_1 + \sqrt{2\gamma} \tilde{a}_1^{in}, \\ \dot{\tilde{b}}_1 &= -iGEe^{i\theta_1} \tilde{a}_1 - \Gamma \tilde{b}_1 + \tilde{\xi}_1, \end{aligned} \quad (6)$$

where  $E_1 = Ee^{-i\theta_1}$ . The equations simply describe a pair of coupled harmonic oscillators, which means that the effective interaction between the cavity field and the vibrational mode of the movable mirror is a simply linear coupling of the form

$$H_{eff} = -\hbar\Omega(t)(\tilde{a}_1^\dagger \tilde{b}_1 e^{-i\theta_1} + \tilde{a}_1 \tilde{b}_1^\dagger e^{i\theta_1}), \quad (7)$$

where  $\Omega(t) = GE$  (the controlled driving field  $E$  is time dependent). Equation (7) is the main result of this paper. This kind of interaction configuration has been used in the design for the light storage through either, the Raman process [10,11] or the EIT scheme [13,14].

A further simplification of the dynamics is possible if the decay rate  $\gamma$  of the cavity field is sufficiently large so that the cavity mode dynamics can also be adiabatically eliminated. In particular, if  $\gamma \gg \Omega(t)$  (but still with  $\nu_1 \gg \gamma$ ), the cavity mode is set  $\tilde{a}_1 = 0$  and

$$\tilde{a}_1 \approx \frac{\Omega(t)}{\gamma} \tilde{b}_1 - \sqrt{2/\gamma} \tilde{a}_1^{in} \quad (8)$$

is substituted into the equation for  $\tilde{b}_1$  to give [11]

$$\begin{aligned} \dot{\tilde{b}}_1 &\approx -\frac{[\Omega(t)]^2}{\gamma} \tilde{b}_1 + \Omega(t) \sqrt{2/\gamma} \tilde{a}_1^{in} - \Gamma \tilde{b}_1 + \tilde{\xi}_1 \\ &= -\eta(t) \tilde{b}_1 - \Gamma \tilde{b}_1 + \sqrt{2\eta(t)} \tilde{a}_1^{in} + \tilde{\xi}_1, \end{aligned} \quad (9)$$

where we have set  $\theta_1 = -\pi/2$  for simplicity and defined  $\eta(t) = [\Omega(t)]^2/\gamma$ . The motional dynamics of the mirror thus is reduced to that of a simple damped harmonic oscillator with the time-dependent damping rate  $\eta(t)$  and motional decoherence rate  $\Gamma$ . If  $\eta(t) \gg \Gamma$ , the statistics of the input light field can be “written onto” the vibrational state of a movable mirror.

### III. LIGHT-TO-MOTION QUANTUM-STATE TRANSFER

Recently, entangling two macroscopic oscillators of movable mirrors in a cavity [28] has been proposed. Here, we present a different protocol that transfers entanglement of a pair of quantum-correlated light fields into a pair of macroscopic oscillators of movable mirrors which are the cavity mirrors in two independent optical cavities. The source of quantum-correlated light fields is a NOPA operating below threshold [33,34]. The output light fields are frequency degenerate, but polarization nondegenerate. We denote the annihilation operators for the two intracavity field modes with

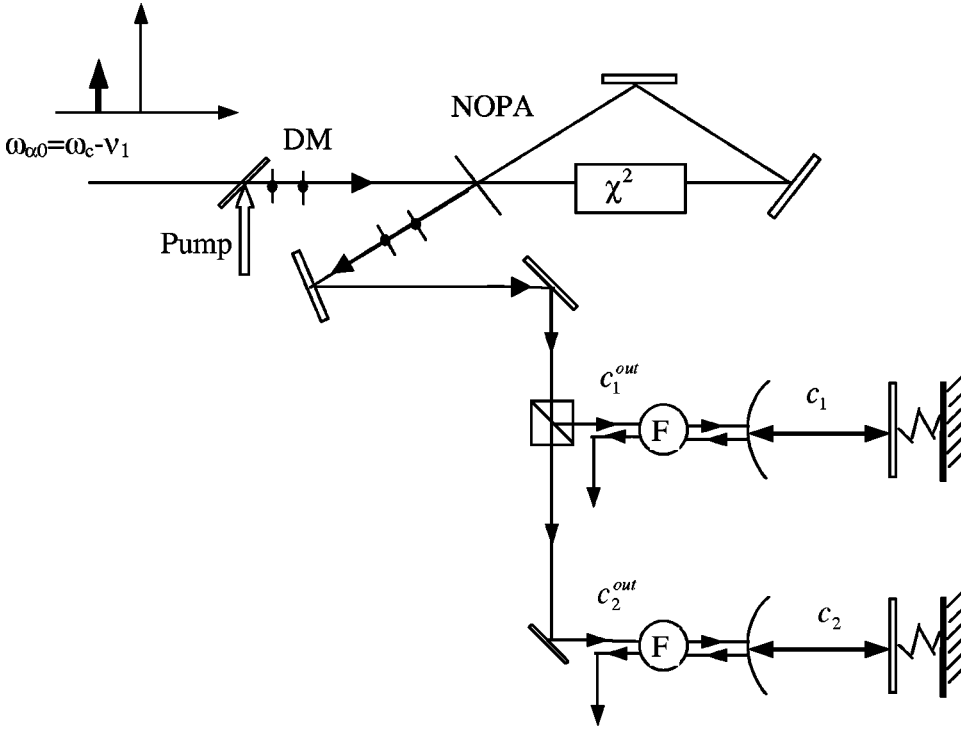


FIG. 2. Preparation of an EPR state of the motion of two movable mirrors. The output modes from a NOPA are incident on two separated movable mirrors.  $F$  is used for facilitating a unidirectional coupling between the entangled light source and the cavity systems with a movable mirror. The driving field is injected into NOPA with the red detuning  $\nu_1$ . The polarizations of two output modes from NOPA are orthogonal, which are separated by a polarized beam splitter (PBS).

orthogonal polarization by  $c_1$  and  $c_2$ , and the Hamiltonian describing the coupling between these modes takes the forms (in a rotating frame)

$$H_1 = ih\chi(c_1c_2 - c_1^\dagger c_2^\dagger), \quad (10)$$

where  $\chi$  is the coupling strength that is proportional to the nonlinear susceptibility of the intracavity medium and the intensity of the coherent pump field. Assuming that the cavity modes are damped with the same rate  $k_c$ , the equations of motion for the mode operators can be derived as

$$\begin{aligned} \dot{c}_1 &= -k_c c_1 - \chi c_2^\dagger + \sqrt{2k_c} c_1^{\text{in}}, \\ \dot{c}_2 &= -k_c c_2 - \chi c_1^\dagger + \sqrt{2k_c} c_2^{\text{in}}, \end{aligned} \quad (11)$$

where  $c_1^{\text{in}}$  and  $c_2^{\text{in}}$  are the vacuum input fields of two cavity modes of NOPA. The output fields from the NOPA then follow from the boundary conditions

$$\begin{aligned} c_1^{\text{out}} &= \sqrt{2k_c} c_1 - c_1^{\text{in}}, \\ c_2^{\text{out}} &= \sqrt{2k_c} c_2 - c_2^{\text{in}}. \end{aligned} \quad (12)$$

Quadrature phase amplitudes for the output fields are defined by

$$\begin{aligned} X_i^{\text{out}} &= c_i^{\text{out}} + c_i^{\text{out}\dagger}, \\ Y_i^{\text{out}} &= -i(c_i^{\text{out}} - c_i^{\text{out}\dagger}), \end{aligned} \quad (13)$$

where  $i$  is 1 or 2. Equation (11) are readily solved in a Fourier-transformed space. The sum of the quadrature amplitudes and the difference of the quadrature phases are derived as

$$\begin{aligned} X_1^{\text{out}}(\omega) + X_2^{\text{out}}(\omega) &= \frac{k_c - \chi + i\omega}{k_c + \chi - i\omega} [X_1^{\text{in}}(\omega) \\ &\quad + X_2^{\text{in}}(\omega)] \xrightarrow{(\chi \rightarrow k_c, \omega \rightarrow 0)} 0, \end{aligned}$$

$$\begin{aligned} Y_1^{\text{out}}(\omega) - Y_2^{\text{out}}(\omega) &= \frac{k_c - \chi + i\omega}{k_c + \chi - i\omega} [Y_1^{\text{in}}(\omega) \\ &\quad - Y_2^{\text{in}}(\omega)] \xrightarrow{(\chi \rightarrow k_c, \omega \rightarrow 0)} 0. \end{aligned} \quad (14)$$

Obviously, the two output fields are highly correlated if  $\chi \rightarrow k_c$  and  $\omega \rightarrow 0$ , it means that the quadrature amplitudes exhibit exactly the property of the EPR correlation.

As depicted in Fig. 2, the two output fields of the NOPA are incident on two space-separated optical cavities, with a movable mirror as described in Sec. II. We require that the output fields from the NOPA are resonant with the cavity mode frequency and the driving fields are red detuned  $\Delta = \nu_1$ . We assume that  $\eta(t) = \eta$ , a constant, and, for simplicity,  $\eta$  is the same for both configurations. Denoting the motional mode operators of the two mirrors with  $\tilde{b}_1$  and  $\tilde{b}_2$ , the two systems are described by

$$\begin{aligned} \dot{\tilde{b}}_1 &= -\eta \tilde{b}_1 - \Gamma \tilde{b}_1 + \sqrt{2\eta} c_1^{\text{out}} + \tilde{\xi}_1, \\ \dot{\tilde{b}}_2 &= -\eta \tilde{b}_2 - \Gamma \tilde{b}_2 + \sqrt{2\eta} c_2^{\text{out}} + \tilde{\xi}_2, \end{aligned} \quad (15)$$

where we assume that the coupling between the NOPA and the cavities is unidirectional. If the bandwidths of the input light fields from the NOPA are sufficiently broad, in particular, if  $k_c \gg \eta$ ,  $c_1^{\text{out}}$  and  $c_2^{\text{out}}$  can be regarded as quantum

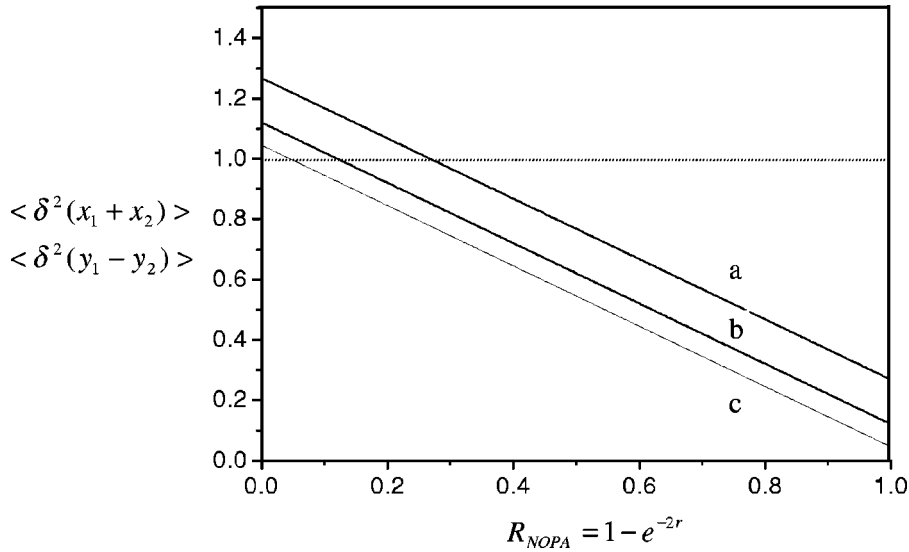


FIG. 3. Evolution of the variance  $\langle \delta^2(x_1 + x_2) \rangle$  and  $\langle \delta^2(y_1 - y_2) \rangle$  of the positions and momenta of the two movable mirrors versus the squeezing rate  $R_{NOPA}$  from the NOPA output field for given values of  $\Omega/(2\pi)$ , (a)  $2 \times 10^5$  Hz; (b)  $3 \times 10^5$  Hz; (c)  $5 \times 10^5$  Hz; with  $V_b = 5306$  at a temperatures of 4 K.

white-noise operators in Eq. (15). Then the variances of positions and momenta of the movable mirrors are given by

$$\begin{aligned} \langle \delta^2(x_1 + x_2)/2 \rangle = \langle \delta^2(y_1 - y_2)/2 \rangle = & \frac{\eta'}{(1 + \Gamma')(\eta' + \Gamma')} e^{-2r} \\ & + \frac{\Gamma'(\eta' + \Gamma') + \Gamma'}{(1 + \Gamma')(\eta' + \Gamma')} V_b, \end{aligned} \quad (16)$$

where  $\eta' = \eta/\gamma$  and  $\Gamma' = \Gamma/\gamma$ ;  $r = -\ln[(k_c - \chi)/(k_c + \chi)]$  is the squeezing parameter of the NOPA and  $V_b = 1 + 2n_T$  is noise introduced from the mechanical damping rate of the movable mirror. So in the limit  $r \rightarrow \infty$  and  $\gamma \gg \Omega \gg \Gamma$ , an EPR state in the positions and momenta of the two movable mirrors is established, i.e.,  $\langle \delta^2(x_1 + x_2)/2 \rangle = \langle \delta^2(y_1 - y_2)/2 \rangle = e^{-2r} \rightarrow 0$ . The coupling to the external fields from the NOPA and driving beam can then be turned off by setting to zero. The result is a stored EPR state for the motion of two mirrors that would persist for a duration set by the time scale  $(\Gamma n_T)^{-1}$  for motional decoherence [18–22].

#### IV. DISCUSSION AND CONCLUSION

We now consider briefly the conditions under which the most significant assumptions required by our model may be satisfied. We assume that the mechanical oscillators with mass  $m = 10^{-6}$  Kg and resonance frequency  $\nu_1/(2\pi) = 10$  MHz are operated at a reasonable temperatures, e.g., 4 K and the mechanical damping rate for the mirror  $\Gamma/(2\pi)$  is  $\sim 1$  Hz. It means there are purely quantum effects at macroscopic scale notwithstanding  $k_B T \gg \hbar \nu_1$ , where  $k_B$  is the Boltzmann constant and  $T$  is the equilibrium temperature of the movable mirror. First, the neglect of the negative sideband  $\tilde{a}_-$  in the effective interaction Hamiltonians requires that the red-detuned frequency of the driving field is equal to

the mechanical resonance frequency  $\nu_1$  and is larger than the cavity decay rate  $\gamma$ . Second, in the overdamped regime [i.e., for Eq. (9)] the cavity decay rate  $\gamma$  is larger than the magnitude of the effective coupling rate  $\Omega(t)$ . If we assume, for example, that the mirrors forming the cavity are separated by a distance  $l = 0.5$  mm, one obtains  $\gamma/(2\pi) = 2$  MHz for a cavity finesse of 75 000. In Fig. 3 we show the evolution of the variance  $\langle \delta^2(x_1 + x_2) \rangle$  and  $\langle \delta^2(y_1 - y_2) \rangle$  of the positions and momenta of the two movable mirrors versus the squeezing rate  $R_{NOPA}$  of the NOPA output field ( $R_{NOPA} = 1 - e^{-2r}$ ) for given values of  $\Omega$ . Figure 3 shows that the entanglement transfer will be perfect when  $\Omega/(2\pi) > 5 \times 10^5$  Hz.

In conclusion, we have proposed a scheme to transfer a quantum state of a light field into a macroscopic, collective vibrational degree of freedom of a massive mirror. Our proposal may realize an EPR state in position and momentum for a pair of massive mirrors at distinct locations by exploiting NOPA. Moreover, our model may be extended to read out the quantum information written in the mirror motion [11] and teleport the quantum state of the mirrors [30]. The possibility of preparing entangled state at a macroscopic level may prove to be useful for high-precision metrology and for implementing a variety of continuous variable quantum computation and communication protocols [35–41], as well as opening the door to further investigations of fundamental aspects of entangled quantum systems.

#### ACKNOWLEDGMENTS

This research was supported by the National Fundamental Research Program (Grant No. 2001CB309304), the National Natural Science Foundation of China (Approval Nos. 60178012, 60238010), and the Shanxi Province Young Science Foundation (Grant No. 20021014).

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