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ENTANGLED STATES IN OPTICS

Entangled States in Spin Subsystems of Polyatomic Ensembles

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Abstract—The feasibility of formation entangled states in spin subsystems of two spatially separated atomic ensembles is discussed. It is shown that the process of cooperative Raman scattering of correlated photon pairs emitted by a parametric source in a below-threshold regime may serve as a physical mechanism for the entanglement. A scheme of teleportation of a coherent spin state in systems of spatially separated ensembles is proposed. © *2003 MAIK "Nauka/Interperiodica".*

INTRODUCTION

The growing interest in problems of quantum information, which combines the fields of quantum cryptography, quantum teleportation, and quantum computing, is related both to indisputable successes in model experiments and to a deeper understanding of the practical importance of the problem itself [1]. In particular, the experiment on teleportation of an optical coherent state described in [2, 3] has given an impetus to the use of quantum states in sets of continuous variables, including, first of all, the states of a harmonic oscillator or of a system of harmonic oscillators.

The main physical concept in design of any quantum-information scheme is based on entanglement of quantum states or on inseparability of the density matrix of spatially separated physical objects. Using, for this purpose, systems described by continuous variables is a promising trend characterized by a large number of physical examples and experimental implementations [3–9]. This is natural, because any quantum state of such a fundamental physical object as harmonic oscillator can be controlled, as a rule, with an appropriate experimental arrangement. It is of basic importance that systems of discrete and continuous variables show certain analogy, which allows one to count on wider practical use of the latter systems. Indeed, the basis of discrete variables, for which one can find well-elaborated sets of quantum protocols and algorithms, is the so-called quantum bit of information. A quantum bit is a superposition of possible states of a two-level system and is described by the SU(2) transformations. For continuous variables, the analogue of the quantum bit is a squeezed state, described by the SU(1,1) algebra. This state can be considered as a quantum superposition of a continuous set of coherent states. On the basis of this analogy, a number of different quantum-information protocols and basis operations have been proposed for a quantum computer operating in a system of continuous variables [10–13].

In this paper, we discuss one possible physical implementation of a system of continuous (more exactly, quasi-continuous) variables in the spin subsystems of two spatially separated macroscopic atomic ensembles. The model is based on the fact that, for macroscopic atomic ensembles containing a large number of atoms and characterized by an initial spin orientation along a certain direction, the macroscopic fluctuations of transverse components of the angular momenta are similar to fluctuations of the quadrature components of a harmonic oscillator. We will show how the cooperative Raman scattering of weak spontaneous parametric emission can introduce quantum correlations into this system and, thus, form an entangled quantum state. We will also show how such correlations can be used in the problem of quantum information transfer for teleportation of an unknown coherent quantum state.

COOPERATIVE RAMAN SCATTERING OF CORRELATED PHOTONS IN THE SYSTEM OF A SPATIALLY SEPARATED ATOMS

As was shown in [7], one possible mechanism for creation of entanglement in a spin subsystem of two atoms can be realized by cooperative Raman scattering. If the scattered photons are twins born due to a spontaneous parametric process (down-conversion or four-wave mixing), the correlation between these photons will be partially transferred to the atomic subsystem. In this case, the atoms may be separated in space by a large distance and the photons can be delivered to the point of location of the atoms through a fiber.

An example of such scattering is illustrated by Fig. 1, which shows how identically linearly polarized photons of different spectral modes ω_1 and ω_2 are scattered by the $s = 1/2 \longrightarrow j = 1/2$ transition (s and j are the

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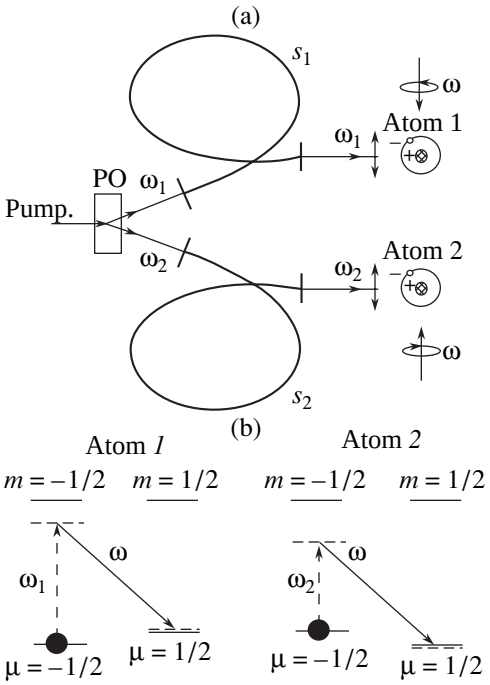

Fig. 1. Cooperative Raman scattering in a system of two spatially separated atoms. (a) The linearly polarized spectral modes ω_1 and ω_2, which are delivered through the waveguides s_1 and s_2 to the points of residence of the atoms and correspond to correlated photon pairs emitted by an optical parametric oscillator (PO) in the below-threshold regime, excite Raman-type optical transitions. The Raman scattering process is additionally stimulated by circularly polarized coherent modes with the frequency ω. (b) Energy-level diagram for the optical transitions under consideration. The stimulated Raman scattering is assumed to be nonresonant for each of the atoms, but resonant for the whole system.

angular momenta of the ground and excited states). Such an optical transition couples the $^2S_{1/2}$ ground spin state of the one-electron atom to the $^2P_{1/2}$ excited state and approximately corresponds to the D_1 transition of an alkali atom in the absence of hyperfine interaction. As shown in Fig. 1, the atoms originally populate a Zeeman state with $\mu = -1/2$, while Raman scattering to the $\mu = +1/2$ state, initiated by correlated quantum modes, is additionally stimulated by the counterclockwise-polarized classical coherent mode at the frequency ω. The atoms are assumed to be in magnetic fields of the same magnitude but opposite direction, and the energy sublevels are shifted in such a way that the degeneracy condition for the states $|-1/2, -1/2\rangle \equiv |1, 1\rangle$ and $|+1/2, +1/2\rangle \equiv |2, 2\rangle$ in the cooperative space of the atoms remains valid.[1] In the general case, the system is supposed to be detuned for stimulated Raman scattering in each of the atoms, but tuned in resonance for the whole system:

$$\omega_1 + \omega_2 = 2\omega. \quad (1)$$

Note that the quantum and classical modes irradiate the atoms in different directions, and, as a result, these channels of excitation do not interfere. It is the appearance of correlated photons of quantum modes that opens an excitation channel.

The state of the system of the two atoms is described by components of the two-particle density matrix $\rho_{ij, kl}$ where the subscript pairs refer, respectively, to the first and second atoms. At the first step of the excitation, the light is scattered from quantum to classical modes. This process is accompanied by transition of the atomic subsystem to a cooperative coherent state, described by the two-particle density matrix $\rho_{12, 12}$. Upon subsequent evolution, however, the cooperative coherence is destroyed due to spontaneous Raman and Rayleigh scattering of the classical coherent modes, occurring in the two atoms independently. In addition, even when the system is excited by photon pairs, there is a small, but not negligible, probability of independent Raman scattering of quantum-mode photons of different pairs. To compare the effect of different excitation channels on the dynamics of a two-particle density matrix, in [7], quantum-electrodynamic calculations were performed for a special type of correlation functions of the quantum field. The results of these calculations are shown in Fig. 2.

Figure 2 shows three dependences: The dotted curve takes into account the cooperative Raman scattering of correlated photons of a pair of atoms into classical modes; the dashed curve additionally takes into account the spontaneous Rayleigh and Raman scattering of classical modes, characterized by the rate constant Γ; and the solid curve represents an exact account of all the processes. As seen from these dependences, at the initial stage of evolution, when the spontaneous processes related to the classical mode have not had time to develop (so that $\Gamma t \ll 1$), the main effect is the appearance of cooperative coherence $\rho_{12, 12}$, with the atoms themselves remaining, with a high probability, in the initial states $\rho_{11, 11} \longrightarrow 1$. This is accompanied, as shown in the last plot, by squeezing of the collective atomic spin with an uncertainty below the standard quantum limit ($\delta S_y < 1/\sqrt{2}$), which is described, with good accuracy, by a dynamic dependence (the dotted curve) and indicates entanglement of the atomic spin states. The main result of the calculations is the possibility of representing the density matrix in the form

$$\hat{\rho}(t)\big|_{t \ll \Gamma^{-1}} = p(t)|\Psi\rangle\langle\Psi| + \hat{\rho}', \quad (2)$$

where

$$|\Psi\rangle = |\Psi(t)\rangle = c_1(t)|1, 1\rangle + c_2(t)|2, 2\rangle \quad (3)$$

[1] We assume that the light shifts of the right-hand sublevels, related to classical modes, can be compensated by introducing additional magnetic fields into the atomic Hamiltonians, so that the degeneracy in the system of cooperative states $|1, 1\rangle$ and $|2, 2\rangle$ holds.

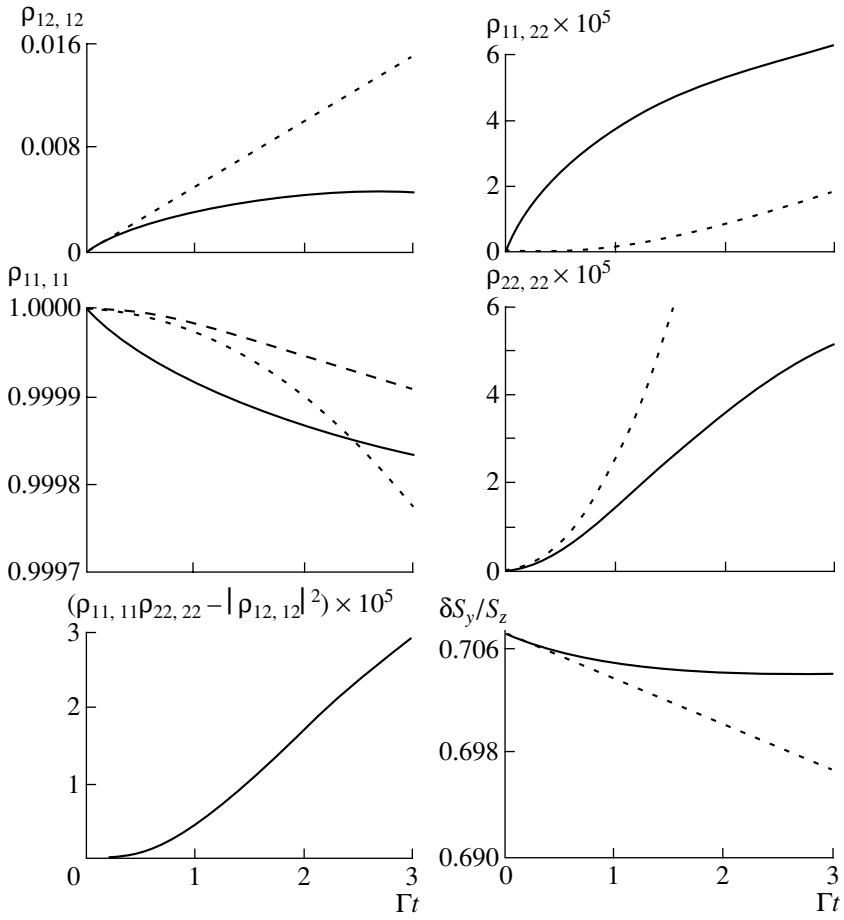

Fig. 2. Time evolution of components of a two-particle density matrix and squeezing effect in the spin system of two spatially separated atoms according to results in [7]. The solid curve corresponds to the exact solution; the dashed curve, to solution of the rate equation with allowance for only cooperative scattering and spontaneous Rayleigh and Raman scattering of the classical mode; and the dotted curve, to the effective Hamiltonian approximation. The dependences are considered as functions of time in units of Γ^{-1} (see the text).

corresponds to transition of the system to an entangled state. It is important that now $|c_2| \sim |\rho_{12,12}| \gg \mathrm{Tr}\hat{\rho}'$ and $p(t) \longrightarrow 1$.

Note that the above mechanism of entanglement or squeezing in the spin subsystem of two atoms is rather weak and the mean collective spin momentum is conserved with a high accuracy, $S_z \approx -1$. The main impediment is the presence of concomitant uncorrelated random processes that destroy the quantum correlations. However, as will be shown below, the cooperative Raman scattering process may cause substantial entanglement of spin fluctuations in a macroscopic system, when, instead of two atoms, we deal with two spatially separated polyatomic ensembles.

ENTANGLEMENT OF SPIN FLUCTUATIONS IN POLYATOMIC ENSEMBLES

Consider the above cooperative Raman scattering process in a system of two spatially separated macroscopic atomic ensembles comprised of N_1 and N_2 atoms, respectively. Let us assume that the time of irradiation of the system (t) is limited by the following two conditions. First, we assume that $\Gamma t \ll 1$, which allows us to neglect the effects of spontaneous scattering of classical modes. Second, we assume that the effect of independent Raman-type transitions due to scattering of photons of different pairs does not substantially redistribute the atoms. By denoting the rate of independent Raman transition in one atom by R, we obtain that $R(N_1 + N_2)t \ll 1$. This inequality means, in fact, that the atoms of the two ensembles stay in states $|1\rangle$, during the whole interaction cycle, and the cooperative scattering effect is reduced to the appearance of coherent (nondiagonal) components of the many-particle density matrix.

Since the excitation geometry proposed corresponds to irradiation of the atoms by quantum and classical fields along different directions, we will assume, in the polyatomic case, that the atoms of each ensemble are

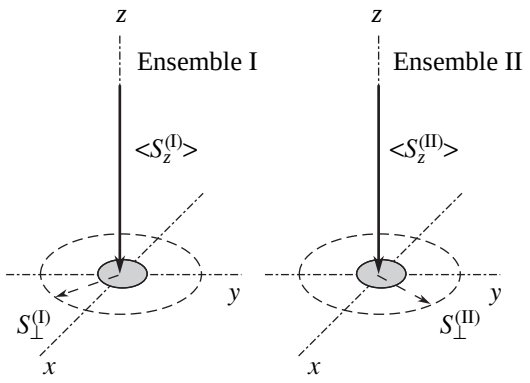

Fig. 3. The regions of quantum uncertainties for transverse fluctuations of the collective spins of atomic ensembles (initially characterized by Poissonian scale and shaded in the figure) significantly increase due to the optical interaction. The collective spin angular momentum fluctuations, in this case, become strongly correlated along the *x* axis and anticorrelated along the *y* axis.

confined within a spatial scale comparable with the irradiating wavelength or are arranged along the line of intersection of the wavefronts of the quantum and classical fields. This condition is a technical limitation, related to a particular way of realization of the Λ scheme in the case of the 1/2 ⟶ 1/2 transition. Actually, we need this condition to provide equivalent conditions for excitation of all the atoms inside the ensemble. We may also assume that, after excitation, the atoms can be redistributed due to thermal motion inside any macroscopic volume. It is exactly this situation that will be considered in the next section.

Under these conditions, the evolution of the combined system is dynamic and is described by the following effective Hamiltonian:

$$\begin{aligned}\mathcal{H}_{\text{eff}} &= \sum_{a,b} i\hbar\chi s_+^{(a)} s_+^{(b)} - i\hbar\chi^* s_-^{(a)} s_-^{(b)} \\ &= i\hbar\chi S_+^{(\text{I})} S_+^{(\text{II})} - i\hbar\chi^* S_-^{(\text{I})} S_-^{(\text{II})},\end{aligned} \tag{4}$$

where χ is the constant of cooperative scattering of the correlated photons. This constant, being proportional to the anomalous amplitude correlations of the quantum field, is responsible for the appearance of the cooperative coherence in the two-atom problem and satisfies the important inequality $|\chi| \gg R$. The spin operators of creation and annihilation for excitations of individual atoms pertaining to the first ($s_\pm^{(a)}$, $a = 1$–N_1) and second ($s_\pm^{(b)}$, $b = 1$–N_2) ensembles are defined in a standard way [14]. Due to the equivalence of all the atoms inside each ensemble, their spin variables can be represented in the Hamiltonian by the appropriate collective spin operators $S_\pm^{(\text{I})}$ and $S_\pm^{(\text{II})}$.

Let us introduce the normalized operators

$$\begin{aligned}a_+^{(\text{I})} &= \frac{1}{\sqrt{N_1}} S_+^{(\text{I})}, \quad a_-^{(\text{I})} = \frac{1}{\sqrt{N_1}} S_-^{(\text{I})}, \\ a_+^{(\text{II})} &= \frac{1}{\sqrt{N_2}} S_+^{(\text{II})}, \quad a_-^{(\text{II})} = \frac{1}{\sqrt{N_2}} S_-^{(\text{II})},\end{aligned} \tag{5}$$

which obey the commutation relationships

$$\begin{aligned}[a_+^{(\text{I})}, a_-^{(\text{I})}] &= \frac{2}{N_1} S_z^{(\text{I})} \longrightarrow -1, \\ [a_+^{(\text{II})}, a_-^{(\text{II})}] &= \frac{2}{N_2} S_z^{(\text{II})} \longrightarrow -1.\end{aligned} \tag{6}$$

In conformity with the above assumptions, in the process of excitation, the atoms of both ensembles remain for the most part on the sublevels $|1\rangle$. Thus, when this relationship is averaged over a time-evolving density matrix, its right-hand side can be considered as time-independent. In this form, commutation relationships (6) are similar to those for two harmonic oscillators, and Hamiltonian (4) describes their parametric excitation. In this case, as follows from the above relations, the parametric excitation process is described by the constant $\tilde{\chi} = \sqrt{N_1 N_2}\chi \gg \chi$, which substantially exceeds the seed constant χ of the two-atom problem. In what follows, these constants are assumed to be real and positive.

This enhancement results from the fact that, in macroscopic systems, the number of pairs participating in the cooperative scattering increases in proportion to the product $N_1 N_2$, whereas the number of uncorrelated transitions increases in proportion to the number of atoms $N_1 + N_2$. Thus, the entanglement of the spin variables of the first and second ensembles can be effective implemented even after a small number of events of cooperative scattering of the correlated photon pairs. In this case, because of weakness of the irradiating flow of quantum modes, concomitant random processes of scattering of photons of different pairs do not have enough time to appear.

Let us assume that the number of atoms in the two ensembles is the same ($N_1 = N_2 \equiv N$). Then, the solution of the equations of motion yields the following result for variations of the mutual transverse spin fluctuations in the first and second ensembles:

$$\begin{aligned}\langle [S_x^{(\text{I})} - S_x^{(\text{II})}]^2 \rangle_t &= \frac{N}{2} e^{-2\tilde{\chi}t} \longrightarrow 0, \\ \langle [S_y^{(\text{I})} + S_y^{(\text{II})}]^2 \rangle_t &= \frac{N}{2} e^{-2\tilde{\chi}t} \longrightarrow 0.\end{aligned} \tag{7}$$

This result shows that, in the process of the cooperative Raman scattering of correlated photons, the spin subsystems of both ensembles pass to the entangled state. The arising correlation of the spin fluctuations is vividly illustrated by Fig. 3. After the optical coupling, the magnitudes of fluctuations of the transverse spin com-

ponents significantly increase, but become strongly correlated in accordance with relations (7). We will note, without going into detail, that, due to statistically independent scattering of photons of different pairs, the correlations are partially destroyed, and the ultimate suppression of the standard quantum limit on the left-hand side of Eq. (7) is restricted by the parameter $R/\chi \ll 1$. In other words, the proposed mechanism of entanglement of the spin states works more efficiently for a weaker irradiating quantum beam comprised of a rarefied flow of correlated photon pairs.

TELEPORTATION OF QUANTUM STATES OF SPIN SUBSYSTEMS

As an example of use of entangled states of the Einstein–Podolsky–Rosen (EPR) type, described above, we consider the mechanism of teleportation of a coherent spin state. The scheme presented below was briefly described in [7] and is based on the idea of quantum information exchange between a field and atoms via coherent forward scattering, proposed in [15]. We will assume that, along with ensembles I and II described above, there exists a third similar ensemble, III. All the ensembles are classically equivalent and contain the same number of atoms N. To simplify further reasoning, we will assume that all the ensembles are in a zero magnetic field and that the atoms inside each ensemble occupy a macroscopic volume exceeding the scale of the wavelength for the relevant optical transitions. These restrictions are not fundamental and are used mainly for convenience. In this case, the spin angular momenta do not precess and the light scattering by the atomic ensembles occurs in a coherent way in the forward direction only.

Ensemble III, whose state is being copied, is in a coherent state; i.e., the macroscopic angular momentum of ensemble III is oriented at a small angle to the xy plane, so that both components $S_x^{(III)}$ and $S_y^{(III)}$ are nonzero, but $S_z^{(III)} \approx -N/2$. The transverse components of the angular momentum are small and are comparable with the scale of the fluctuations. The state of this ensemble is unknown and should be teleported to another macroscopic ensemble. Ensemble II is the teleportation target and originally is entangled with ensemble I.

The scheme for copying the quantum state, the teleportation protocol, is shown in Fig. 4. At the first stage, one measures the Faraday rotation angle, induced by the transverse spin orientation of ensembles III and I, for two beams of probe radiation linearly polarized along the z axis that propagate along two different directions.[2] The first beam is directed along the y axis and intersects both ensembles in the same direction.

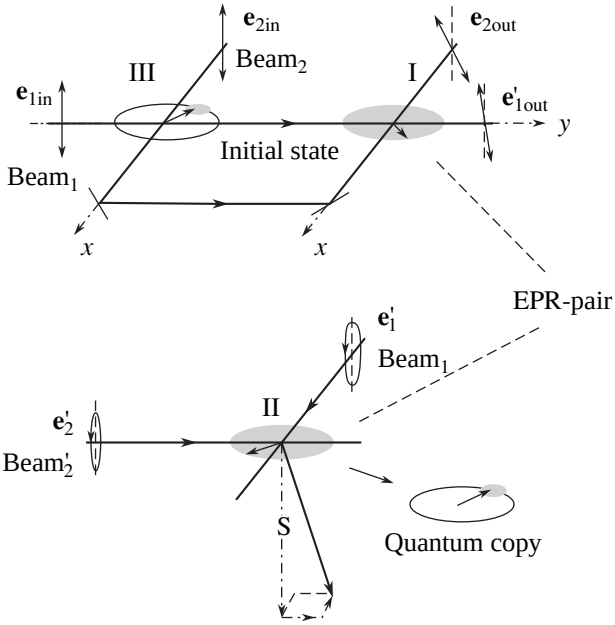

Fig. 4. Teleportation of the coherent state of ensemble III to the spin state of ensemble II. Ensemble III, which is being copied, and one of the ensembles of the EPR pair (in the figure, this is ensemble I) are probed with two nonresonant and nonsaturating beams with the linear polarizations $\mathbf{e}_{1\,in}$ and $\mathbf{e}_{2in}$. The measured rotation angles $\mathbf{e}_{1\,out}$ and $\mathbf{e}_{2\,out}$ of the polarization planes are used for the subsequent unitary transformation of the polarizations $\mathbf{e}'_1$ and $\mathbf{e}'_2$ from linear to elliptical for the two nonresonant beams controlling the state of ensemble II.

The second beam intersects ensemble III in the positive direction of the x axis and ensemble I in the negative direction. The rotation angles measured in this configuration are related to the values of the spin variables as follows:

$$\delta\varphi_1 = \frac{2\pi\bar{\omega}}{cS_0}\alpha(\bar{\omega})[S_y^{(III)} + S_y^{(I)}],$$
$$\delta\varphi_2 = \frac{2\pi\bar{\omega}}{cS_0}\alpha(\bar{\omega})[S_x^{(III)} - S_x^{(I)}], \quad (8)$$

where $\bar{\omega}$ is the frequency of the probe beam, S_0 is the beam cross-section area, and $\alpha(\bar{\omega})$ is the dynamic polarizability for the $1/2 \longrightarrow 1/2$ transition for a circularly polarized beam [16]. The square brackets on the right-hand sides of Eqs. (8) contain combinations of spin operators, which, after the measurement, turn into known c-number quantities. It should be noted that the possibility of simultaneously measuring these combinations is a consequence of the commutativity of the relevant operators. Note also that the spin observables, treated separately for each of the ensembles, conserve their quantum uncertainty after such measurement.

Let us now apply to ensemble II nonsaturating nonresonant classical electromagnetic fields propagating in two directions, as shown in Fig. 4. The two probe

[2] To increase the sensitivity of the measurement, the probing is better performed with polarization squeezed light, as described in [8].

beams involved in this process are assumed to be classically similar to the beams used in the Faraday measurements. They are characterized by the same frequency $\bar{\omega}$, but are elliptically polarized. It is crucial that the elliptical polarization is produced from linear polarization by means of some unitary optical transformation. This transformation reduced to a rotation of the polarization plane of the beams with an additional phase delay. As a result, an additional degree of circular polarization arises, characterized by the Stokes parameters A_1 (probing along the x axis) and A_2 (probing along the y axis; see Fig. 4). The magnitude of the introduced circular polarization is specified by the angles $\delta\varphi_1$ and $\delta\varphi_2$ found above. Let us show how the relation between these parameters arises.

We should take into account that the interaction of a coherent nonsaturating beam with an atomic subsystem is described by an effective Hamiltonian [16]. For the polarization-dependent part of this Hamiltonian, we have

$$H_{\text{eff}} = \hbar A_1 \frac{2\pi\bar{\omega}}{cS_0} J\alpha(\bar{\omega}) S_x^{(\text{II})} + \hbar A_2 \frac{2\pi\bar{\omega}}{cS_0} J\alpha(\bar{\omega}) S_y^{(\text{II})}. \quad (9)$$

We introduced here the total (integrated over the beam cross section) photon flux J, which is assumed to be the same for the first and second beams. As a result of a short interaction with a duration of Δt, the spin operators of the transverse components, treated in the Heisenberg representation, are transformed as follows:

$$\begin{aligned} S_x^{(\text{II})}(\Delta t) &= S_x^{(\text{II})} + A_2 \frac{2\pi\bar{\omega}}{cS_0} J\Delta t\alpha(\bar{\omega}) S_z^{(\text{II})}, \\ S_y^{(\text{II})}(\Delta t) &= S_y^{(\text{II})} - A_1 \frac{2\pi\bar{\omega}}{cS_0} J\Delta t\alpha(\bar{\omega}) S_z^{(\text{II})}. \end{aligned} \quad (10)$$

On the right-hand side of these relationships, the macroscopic angular momentum $S_z^{(\text{II})}$ may be replaced by its mean value, i.e., by $-N/2$. Let us choose the Stokes parameters A_1 and A_2 based on the conditions

$$\begin{aligned} A_1 J\Delta t &= \left[\frac{cS_0}{2\pi\bar{\omega}\alpha(\bar{\omega})}\right]^2 \frac{2}{N}\delta\varphi_1, \\ A_2 J\Delta t &= -\left[\frac{cS_0}{2\pi\bar{\omega}\alpha(\bar{\omega})}\right]^2 \frac{2}{N}\delta\varphi_2. \end{aligned} \quad (11)$$

Taking into account the relation between the measured angles of Faraday rotation and the operators of the spin observables (8), transformations (10) yield the following representation for the spin observables of the second ensemble after the interaction:

$$\begin{aligned} S_x^{(\text{II})}(\Delta t) &= S_x^{(\text{II})} - S_x^{(\text{I})} + S_x^{(\text{III})} \longrightarrow S_x^{(\text{III})}, \\ S_y^{(\text{II})}(\Delta t) &= S_y^{(\text{II})} + S_y^{(\text{I})} + S_y^{(\text{III})} \longrightarrow S_y^{(\text{III})}, \end{aligned} \quad (12)$$

where the right-hand sides are transformed with allowance for the strong correlation of the spin fluctuations in the first and second ensembles. Relationships (12) show that, after all the above manipulations, further measurements of any expected values for the spin observables of ensemble II is equivalent, in the statistical sense, to the measurements that were possible for the original ensemble III. However, the initial quantum state of ensemble III was completely destroyed and, therefore, ensemble II can be regarded as a quantum copy of ensemble III.

Let us make a few remarks. Note that the left-hand sides of Eqs. (11) contain, instead of the Stokes parameter, the difference between the number of right- and left-handed photons passed through the ensemble during the time of the interaction. It is exactly this quantity that may be controlled by a unitary transformation. Indeed, the light linearly polarized along the z axis does not affect the state of the system to the extent that the incoherent light scattering can be neglected compared with the coherent forward scattering. The unitary transformer, which renders the beam elliptically polarized, controls the whole photon flux transmitted through the ensemble during the interaction time. The right-hand sides of relationships (11) should be large, so that the fluctuations of the flux can be neglected and, in Hamiltonian (9), the quantum characteristics of the probe fields can be replaced by their mean classical values. Note also that the optimum value of the parameter S_0 corresponds to the size of the atomic cloud. In this case, all the atoms of the ensembles are efficiently coupled to all the photons of the light beams, which was implied in our reasoning.

CONCLUSIONS

In this paper, we consider an optical scheme for generating entangled states in spin subsystems of two spatially separated atomic ensembles. The mechanism of cooperative Raman scattering of correlated photon pairs emitted by a parametric source in the below-threshold regime is proposed as a physical process leading to the entanglement of spin fluctuations. An analogy is shown between the description of macroscopic fluctuations of transverse components of collective spin angular momenta and the description of quadrature components of a harmonic oscillator. A scheme of teleportation of a coherent spin state is proposed on the basis of the Faraday effect and the possibility (under conditions of coherent forward scattering) of dynamic interaction of a probe field with the spin subsystem. The processes considered in the paper are primarily interest for problems of quantum information, when a system with continuous variables is used as an information carrier.

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REFERENCES

1. D. Bouwmeester, A. Ekert, and A. Zeilinger, *The Physics of Quantum Information* (Springer, Heidelberg, 2000).
2. S. L. Braunstein and H. J. Kimble, Phys. Rev. Lett. **80**, 869 (1998).
3. A. Furusawa, J. L. Sørensen, S. L. Braunstein, *et al.*, Science **282**, 706 (1998).
4. J. Hald, J. L. Sørensen, C. Shori, and E. S. Polzik, Phys. Rev. Lett. **83**, 1319 (1999).
5. B. Jualgaard, A. Kozhekin, and E. S. Polzik, Nature **413** (6854), 400 (2001).
6. Ch. Silberhorn, P. K. Lam, O. Weiss, *et al.*, Phys. Rev. Lett. **86**, 4267 (2001).
7. D. V. Kupriyanov, I. M. Sokolov, and A. V. Slavgorodskii, Phys. Rev. A **63**, 063811 (2001).
8. C. Shori, B. Julsgaard, J. L. Sørensen, and E. S. Polzik, quant-ph/0203023, Phys. Rev. Lett. (in press).
9. S. Mancini, V. Giovanetti, D. Vitali, and P. Tombesi, Phys. Rev. Lett. **88**, 120401 (2002).
10. S. Lloyd and S. L. Braunstein, Phys. Rev. Lett. **82**, 1784 (1999).
11. T. C. Ralph, Phys. Rev. A **61**, 010303 (1999).
12. Ch. Silberhorn, N. Korolkova, and G. Leuchs, Phys. Rev. Lett. **88**, 167902 (2002).
13. S. D. Bartlett, B. C. Sanders, B. T. H. Varcoe, and H. Guise, quant-ph/0011080 v3 (2002).
14. L. D. Landau and E. M. Lifshitz, *Course of Theoretical Physics*, Vol. 3: *Quantum Mechanics: Non-Relativistic Theory*, 4th ed. (Nauka, Moscow, 1989; Pergamon, New York, 1977).
15. A. Kuzmich and E. S. Polzik, Phys. Rev. Lett. **85**, 5639 (2000).
16. W. Happer, Rev. Mod. Phys. **44**, 169 (1972).

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