Colloquium: Trapping and manipulating photon states in atomic ensembles

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Modern optical techniques allow one to accurately control light using atoms and to manipulate atoms using light. In this Colloquium the author reviews several ideas indicating how such techniques can be used for accurate manipulation of quantum states of atomic ensembles and photons. First a technique is discussed that allows one to transfer quantum states between light fields and metastable states of matter. The technique is based on trapping quantum states of photons in coherently driven atomic media, in which the group velocity is adiabatically reduced to zero. Next, possible mechanisms are outlined for manipulating quantum states of atomic ensembles. Specifically, a "dipole blockade" technique is considered in which optical excitation of mesoscopic samples into Rydberg states can be used to control the state of ensembles at the level of individual quanta. It is also noted that even simple processes involving atom-photon correlations can be used to effectively manipulate the ensemble states. Potentially these techniques can be used for implementation of important concepts from quantum information science.

CONTENTS

I. INTRODUCTION: QUANTUM COHERENCE AND CONTROL

The work described in this Colloquium is part of the broad effort to develop new techniques for manipulating quantum states of matter and light. Here we discuss examples of how the tools from quantum optics and atomic physics can be used to attain this goal. Recently, these efforts have been stimulated by the emerging field of quantum information. Hence, we begin with a brief overview of coherence and quantum control, relating them to ideas of quantum information science.

Quantum theory allows for a system to be placed into a superposition of its distinct states. While undergoing unitary evolution such superposition states can interfere. Quantum superpositions can also involve several parts of a larger system, in which case the different subsystems can be in correlated states. It was recognized early on that these phenomena are perhaps the most intriguing aspects of quantum theory (Bohr, 1935; Einstein *et al.*, 1935; Schrodinger, 1935).

To what extent can these effects be controlled, both in principle and in practice? In the past few decades dramatic conceptual and technological progress in the area of quantum control has been made (Chu, 2002). Quantum-mechanical coherence and interference are now explored in different areas of science and are being used for some noteworthy applications. Specifically, optical and atomic coherence phenomena are central to many areas of quantum optics and atomic physics. Spectacular manifestations ranging from experiments on ultracold atoms and matter waves [involving coherence of center-of-mass motion (Meystre, 2001; Anglin and Ketterle, 2002)] to generation of nonclassical light [involving quantum coherence of photon states (Walls and Milburn, 1994)] and coherent control of atomic and optical properties of resonant media (Scully and Zubairy, 1997) are by now well known. It is not surprising that the best known application of quantum coherence and interference—the atomic clock (Udem *et al.*, 2002)—emerged from many years of atomic and optical physics research.

Recently the ideas of coherence and correlations emerged as the basis for intriguing developments in quantum information science, a new field that combines ideas from quantum mechanics and information theory. Originally stimulated by research into possible limits for information processing, this field by now has a solid theoretical foundation (Preskill, 1998; Steane, 1998; Bouwmeester *et al.*, 2000; Nielsen and Chuang, 2000). For example, a basic concept of quantum computation can be illustrated by noting that a system prepared in a coherent superposition of distinct "registers" and undergoing unitary evolution is capable of effectively evaluating many such "registers" at once, which implies the possibility of the so-called "quantum parallelism."

From the viewpoint of a physicist the implementation of these ideas is essentially an ultimate exercise in controlling quantum systems. In direct analogy to the clas-

sical situation, quantum information should be encoded in a chosen set of physical systems and subsequently stored, manipulated, and communicated (DiVincenzo, 2000). In contrast to the classical case, however, coherent manipulation of quantum superposition states is essential and to ensure unitary evolution during storage and manipulation, quantum states must be well protected from noise and decoherence. In other words, the practical implementation of quantum information requires precise manipulation of many coupled quantum-mechanical systems, which is an extremely challenging task. At present it is not clear if a sufficient degree of such control will ever be attained for the fundamental ideas of quantum information theory to become practical. Nevertheless, these ideas provide an exciting challenge. Furthermore, it is seems likely (as, we hope, this Colloquium will illustrate) that the quantum control techniques now being developed to implement these ideas will allow the study of a wide variety of new problems in physics and beyond. Indeed, it is intriguing to see that the ideas from many areas of physics, ranging from condensed matter to optics and NMR, are now being explored to devise suitable quantum control techniques (Braunstein, 2000). In particular, clear-cut manipulation of entanglement has already been observed in quantum optical systems involving photons, atoms, or trapped ions. There exist excellent reviews of these subjects in various contexts, including quantum cryptography (Gisin *et al.*, 2002), atoms in cavities (Raimond *et al.*, 2001), and ultracold ions and atoms (Monroe, 2002).

In this Colloquium we focus on using techniques involving ensembles of neutral atoms for coherent control of light as well as for storing and manipulating quantum states.

II. "TRAPPING" PHOTON STATES

A. Motivation and overview

Photons are the fastest and very robust carriers of quantum states, but their main strength is also their weakness: they are difficult to localize and store. Ideally one would like to store and manipulate quantum states in matter and to map these states onto photons, when desired. From a quantum information prespective such a system could be viewed as a quantum network in which communication between the nodes occurs via photons (Cirac *et al.*, 1997). Spin states of atoms, for example, represent reliable and long-lived storage units for quantum superpositions. Therefore, the challenge is to develop a technique for coherent transfer of quantum states carried by light into atoms and vice versa. In other words, it is necessary to have a quantum memory that is capable of storing and releasing quantum states on the level of individual quanta. In general, such a device needs to be entirely coherent and reversible and should also allow for suitable manipulation of the stored states.

A conceptually simple approach to quantum memory is to "store" quantum states of single photons in indi-

vidual atoms. This approach involves, in essence, coherent absorption and emission of single photons by single atoms. However, implementing this idea in practice is difficult: the single-atom absorption cross section scales as the square of the optical wavelength and is therefore very small. A very elegant solution is provided by a set of techniques known as cavity QED (Kimble, 1998). Placing an atom in a high-Q resonator effectively enhances its cross section by the number of photon round trips in the cavity and thus makes an effective coupling to single atoms possible. The spectacular experimental progress in this field (Ye *et al.*, 1999; Hennrich *et al.*, 2000; Hood *et al.*, 2000; Foster *et al.*, 2001; Kuhn *et al.*, 2002; McKeever *et al.*, 2002) makes it a viable avenue for studying fundamental physics of atom-photon interaction as well as for quantum networking with possibilities ranging from deterministic single-photon sources to quantum logic operations. However, it is still technically challenging to achieve the required strong-coupling regime and to simultaneously control the motion of individual atoms confined to a small cavity volume.

At the same time, photons can interact much more strongly with ensembles containing a large number of atoms. For example, photons can be absorbed with unit probability in an optically thick atomic sample. Normally such absorption is accompanied by *dissipative* processes, which result in decoherence (i.e., deteriorating quantum states). In the first part of this Colloquium we describe a technique that allows, in principle, for an ideal transfer of quantum correlations between light fields and metastable states of matter. This approach is based on mapping quantum states of photons into coherently driven atomic media, in which the group velocity of the propagating pulses is adiabatically reduced to zero, resulting in a coherently controlled absorption process. Specifically, a technique known as electromagnetically induced transparency (EIT) (Harri, 1997) is used in which the optical properties of atoms can be manipulated by an external (classical) field. This technique helps to alleviate most of the stringent requirements of single-atom cavity QED, and thus could form the basis for a fast and reliable quantum memory for light.

Before proceeding we note that in early work on this subject (Grom and Kuzmich, 1995) it has been shown that even the usual absorption of light can lead to partial mapping of its quantum properties to the atomic ensemble. These ideas have later been extended theoretically (Kuzmich *et al.*, 1997; Kozhekin *et al.*, 2000) and verified experimentally (Hald *et al.*, 1999). However, as a consequence of dissipation, the absorptive methods do not allow one to store the quantum states of *individual* wave packets (corresponding, e.g., to single photons). Instead, they have to involve a stationary source of light (corresponding in essence to identical copies of photons) such as a stationary source of squeezed vacuum.

B. Electromagnetically induced transparency

The strength of the interaction between light and atoms is a function of the wavelength or frequency of
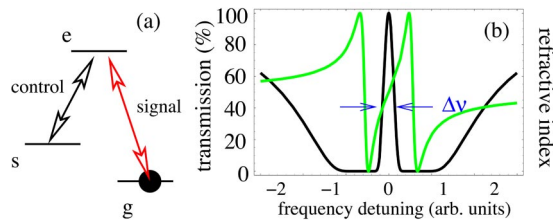
FIG. 1. Electromagnetically induced transparency: (a) Proto-type atomic system for EIT. (b) Spectrum of transmission and refractive index corresponding to EIT. Rapid variation of the refractive index (gray curve) causes a reduction of group velocity (Color in online edition).

light. When the light frequency matches the frequency of a particular atomic transition, a resonance condition occurs and the optical response of the medium is greatly enhanced. Light propagation is then accompanied by strong absorption and dispersion (Scully and Zubairy, 1997), as the atoms are actively promoted into fluorescing excited states. Electromagnetically induced transparency is a technique that can be used to make a resonant, opaque medium transparent by means of quantum interference.

To illustrate this effect, consider the situation in which the atoms have a pair of lower energy states [$|g\rangle$ and $|s\rangle$ in Fig. 1(a)] in each of which the atoms can live for a long time. Such is the case for sublevels of different angular momentum (spin) within the electronic ground state of alkali atoms.[1] In order to modify the propagation through this atomic medium of a light field that couples the ground state $|g\rangle$ to an electronically excited state $|e\rangle$ (the "signal" field), one can apply a second "control" field that is at resonance with the transition $|e\rangle\rightarrow|s\rangle$. The combined effect of these two fields is to stimulate the atoms into a so-called dark superposition of the states $|g\rangle$ and $|s\rangle$. In such a case, the two possible pathways in which light can be absorbed by atoms ($|g\rangle \rightarrow|e\rangle$ and $|s\rangle\rightarrow|e\rangle$) can interfere and cancel each other. The atoms then are said to be in the "dark states." With such destructive quantum interference, none of the atoms are promoted to the excited states, leading to a vanishing light absorption (Arimondo, 1996). This is the essence of the so-called dark resonances and electromagnetically induced transparency (Boller *et al.*, 1991).

In the case when the resonant control field is strong and its intensity is constant in time but the signal field is weak, the response of the atomic ensemble can be described in terms of the linear susceptibility spectrum $\chi(\omega)$:

$$\chi(\omega)=g^2 N \frac{\gamma_{gs}+i\omega}{(\gamma_{ge}+i\omega)(\gamma_{gs}+i\omega)+|\Omega|^2}, \quad (1)$$

where γ_{ij} corresponds to the relaxation rate of the $i\rightarrow j$ coherence, Ω is the Rabi frequency of the control field

[proportional to the electric field amplitude (Scully and Zubairy, 1997)], N is the total number of atoms in the sample, g is the atom-field coupling constant, and ω is the difference between the signal field frequency and the frequency of the atomic transition $|g\rangle\rightarrow|e\rangle$ (with $\omega=0$ corresponding to the exact atom-field resonance). The imaginary part of the susceptibility describes absorptive properties of the medium (thereby modifying the intensity transmission coefficient T), whereas the real part determines the refractive index n:

$$T(\omega)=\exp[-\operatorname{Im}\chi(\omega)kL], \;\; n=1+\frac{1}{2}\operatorname{Re}\chi(\omega), \quad (2)$$

where L is the length of the medium.

Ideal transparency is obtained in the limit when the relaxation of the low-frequency (spin) coherence ($\gamma_{gs}=0$) vanishes, in which case there is no absorption at atomic resonance (see Fig. 1). Many of the important properties of EIT result from the fragile nature of quantum interference in a material that is initially opaque. Indeed, the ideal transparency is attained only at exact resonance, i.e., when the frequency detuning $\omega=0$.[2] Away from this resonance condition the interference is not ideal and the medium becomes absorbing. Hence the transparency spike that appears in the absorption spectrum is typically very narrow [Fig. 1(b)]. At the same time the tolerance to frequency detuning ("transparency window," $\Delta\nu$) can be increased by using stronger coupling fields, since in this case interference becomes more robust.

The basic physics of light interacting with multilevel atoms has been explored over several decades. The first schemes involving atomic coherence generated by strong coupling lasers in Raman-like systems were introduced and studied extensively in the late 1960s (Scully and Zubairy, 1997). More recent work on quantum nondemolition measurements in Raman systems should be especially noted (Roch *et al.*, 1997). The effect that lies at the foundation of current developments, coherent population trapping, was experimentally discovered and theoretically explained by Alzetta *et al.* (1976). There exist a number of reviews on the subject (Arimondo, 1996; Harris, 1997; Lukin and Imamoglu, 2001) that also cover closely related areas such as lasers without inversion (Kocharovskaya, 1992) and subrecoil cooling (Cohen-Tannoudji, 1998).

Since atoms are decoupled from the light fields in an ideal EIT medium, at resonance susceptibility vanishes, and the refractive index is equal to unity. This means that the propagation velocity of a phase front (i.e., the phase velocity) is equal to that in vacuum. However, the narrow transparency resonance is accompanied by a very steep variation of the refractive index with frequency. As a result, the envelope of a wave packet propagating in the medium moves with a group velocity v_g (Harris *et al.*, 1992), where

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[1]In the following we will consider two angular momentum states $|g\rangle$ and $|s\rangle$ as corresponding to generalized spin and we call them spin states accordingly.

[2]More generally, transparency occurs when the frequency difference between the signal and control fields matches the resonance frequency of the two-photon transition $|g\rangle\rightarrow|s\rangle$, ω_{gs}.
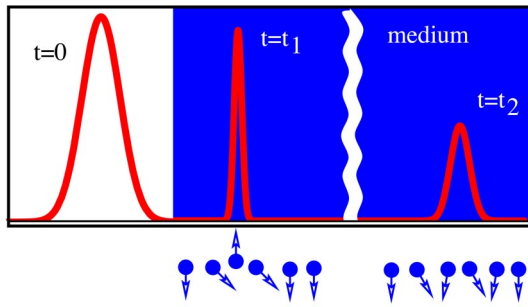
FIG. 2. When a light pulse enters the medium, it exhibits a spatial compression, while photons are converted into atomic (spin) excitation. The slow photonic and spin waves then propagate together. The lossless propagation is limited by the spreading of the pulses due to the narrow bandwidth of the transparency window (Color in online edition).

$$v_g = \frac{c}{1 + g^2 N/|\Omega|^2}, \tag{3}$$

which can be much smaller than the speed of light in vacuum c. Note that v_g depends on the control field intensity and the atomic density: decreasing the control power or increasing the atom density makes v_g slower.

The possibility of manipulating the group velocity and refractive properties in EIT media was pointed out by Harris and co-workers (Harris *et al.*, 1992). Already early experimental work demonstrated these striking properties of EIT (Kasapi *et al.*, 1995; Xiao *et al.*, 1995; Schmidt *et al.*, 1996; Lukin *et al.*, 1997). For example, group velocities of $c/165$ were measured in lead vapor (Kasapi *et al.*, 1995).

The subject was truly brought to the focus of attention by the remarkable experiments of Hau *et al.* (1999), in which an ultracold gas of Na atoms was used to slow light pulses to 17 m/s. Furthermore, this work made use of a large optical density to localize the pulse entirely inside the medium. Experimental work on slow group velocities in hot atomic vapor (Budker *et al.*, 1999; Kash *et al.*, 1999) rapidly followed. Other mechanisms to obtain similarly narrow resonances and correspondingly small group velocities also exist (Inouye *et al.*, 2000).

C. Propagation in an EIT medium and dark-state polaritons

Figure 2 illustrates the dynamics of light propagation in an EIT medium. Initially the pulse is outside the medium in which all atoms are in their ground states ($|g\rangle$). The front edge of the pulse then enters the medium and is rapidly decelerated. Being outside of the medium the back edge still propagates with vacuum speed c. As a result, upon entrance into the cell, the spatial extent of the pulse is compressed by the ratio c/v_g, while its peak amplitude remains unchanged. Clearly the energy of the light pulse is much smaller when it is inside the medium. Photons are being expended to establish the coherence between the states $|g\rangle$ and $|s\rangle$, or, in other words, to

change atomic state, with any excess energy carried away by the control field. As the pulse exits the medium its spatial extent increases again and the atoms return to their original ground state; the pulse however, is delayed as a whole by

$$\tau = (1/v_g - 1/c)L = L/c \times \frac{g^2 N}{|\Omega|^2}. \tag{4}$$

Inside the medium the wave of flipped spins propagates together with the light pulse. The photons in the pulse are therefore strongly coupled to atoms. It turns out that it is possible to associate a quasiparticle with such a slow propagation. This quasiparticle, called a dark-state polariton (Fleischhauer and Lukin, 2000), is a combined excitation of photons and spins.

To see how it emerges, let us consider the quantum evolution of a propagating signal field and atoms in the Heisenberg picture for the case when the decay rate of coherence between states $|g\rangle$ and $|s\rangle$ is negligible. We can describe a propagating signal by the electric field operator $\hat{E}(z,t) = \Sigma_k a_k(t) e^{ikz}$, where the sum is over the free-space modes with wave vectors k and corresponding bosonic operators a_k. To describe the properties of the medium, we use collective atomic operators, averaged over small but macroscopic volumes containing $N_z \gg 1$ particles at position z. In particular, the operator $\hat{e}(z,t) = (N_z)^{-1} \Sigma_{j=1}^{N_z} |g_j\rangle\langle e_j| e^{-i\omega_{ge}t}$ describes the atomic polarization oscillating at an optical frequency, whereas the operator $\hat{s}(z,t) = (N_z)^{-1} \Sigma_{j=1}^{N_z} |g_j\rangle\langle s_j| e^{-i\omega_{gs}t}$ corresponds to a low-frequency spin wave. Here ω_{ge} and ω_{gs} are resonance frequencies of optical and spin-flip transitions, respectively. The control field is assumed to be strong and is treated classically. The atomic evolution is governed by a set of Heisenberg equations: $\dot{\hat{A}} = i/\hbar[\hat{H}, \hat{A}]$, where $\hat{H}$ is the atom-field interaction Hamiltonian (Scully and Zubairy, 1997) and $\hat{A} = \{\hat{e}, \hat{s}\}$. These equations can be simplified assuming that the signal field is weak and that Ω and $\hat{E}$ change in time sufficiently slowly, i.e., adiabatically. To leading order in the signal field $\hat{E}$ we find (see, e.g., Fleischhauer and Lukin, 2000)

$$\hat{e}(z,t) = -\frac{i}{\Omega}\left[\frac{\partial}{\partial t}\hat{s}(z,t)\right], \tag{5}$$

$$\hat{s}(z,t) = -\frac{g\hat{E}(z,t)}{\Omega}. \tag{6}$$

The evolution of the signal field is described by the Heisenberg equation

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{E}(z,t) = igN\hat{e}(z,t), \tag{7}$$

which is essentially analogous to a classical propagation equation.

The solution of Eqs. (5)–(7) can be obtained by introducing a new quantum field $\hat{\Psi}(z,t)$ that is a superposition of photonic and spin-wave components:

$$\hat{\Psi}(z,t) = \cos\theta\, \hat{E}(z,t) - \sin\theta\, \sqrt{N}\,\hat{s}(z,t), \qquad (8)$$

$$\cos\theta = \frac{\Omega}{\sqrt{\Omega^2 + g^2 N}}, \qquad \sin\theta = \frac{g\sqrt{N}}{\sqrt{\Omega^2 + g^2 N}}. \qquad (9)$$

The field $\hat{\Psi}$ obeys the following equation of motion:

$$\left[\frac{\partial}{\partial t} + c\cos^2\theta\,\frac{\partial}{\partial z}\right]\hat{\Psi}(z,t) = 0, \qquad (10)$$

which describes a shape-preserving propagation with velocity $v_g = c\cos^2\theta$ that is proportional to the magnitude of its photonic component. Before proceeding, several important properties of the dark-state polaritons should be noted. First of all, its Fourier components $\hat{\Psi}_k$ and $\hat{\Psi}_k^\dagger$ obey the bosonic commutation relations since $[\hat{\Psi}_k, \hat{\Psi}_{k'}^\dagger] \approx 1$ is the weak signal limit considered here. Second, operators defined by Eq. (8) do not contain the electronic excited state components and hence are immune to spontaneous emission. Finally, we note that the properties of dark-state polaritons (such as propagation velocity and mixing angle) can be easily manipulated, e.g., by changing the intensity of the control field.

The above description is an ideal scenario. In practice two main limitations have to be taken into account. First, there always exists a finite rate of decoherence between different spin states (γ_{gs}). This effect results in damping of the polariton amplitude [see Eq. (13) below] with a rate proportional to its spin component $\gamma_{gs}\sin^2\theta$. This implies that after a characteristic time $\tau_{coh} = \gamma_{gs}^{-1}$ the atoms will end up either in state $|g\rangle$ or $|s\rangle$; the transparency will then be lost and the polariton will disappear. Second, even for infinite τ_{coh}, the pulse delay is limited by the bandwidth of the transparency window, which decreases with propagation distance. This is due to the fact that at higher densities or propagation distances the medium becomes increasingly opaque at frequencies other than the line center; as a result the available transparency window becomes smaller. The EIT medium behaves like a nonabsorbing, linear dispersive medium only within a certain frequency window around the two-photon resonance (Lukin *et al.*, 1997). The adiabatic approximation described above essentially assumes that all relevant dynamics happens within this frequency window.

The transparency window is defined by the intensity transmission of the medium, Eq. (2). Close to EIT resonance, $\omega \to 0$, we can expand $\mathrm{Im}\,\chi(\omega)$ in a power series to find

$$T(\omega) \approx \exp\{-\omega^2/\Delta\nu^2\} \quad \text{with} \quad \Delta\nu = \left[\frac{c}{\gamma_{ge}L}\frac{|\Omega|^4}{g^2 N}\right]^{1/2}. \qquad (11)$$

One recognizes that the transparency bandwidth decreases with increasing length or decreasing control intensity. After a sufficiently long propagation this results in spreading of the signal pulse as indicated in Fig. 2. In order to preserve the pulse, its bandwidth $1/T$ should be smaller than the transparency bandwidth: $1/T < \Delta\nu$. This puts a limit on the ratio of delay and pulse duration.

Using $\Delta\nu$ from Eq. (11) together with the expression for τ from Eq. (4), we find that τ/T should be smaller than $\sqrt{g^2 NL/(c\gamma_{ge})}$, which is the square root of the optical depth of the medium. Hence, the delay can exceed the pulse duration, or in other words, the entire pulse can be localized inside the medium only if an optically dense medium is used, i.e.,

$$\frac{g^2 NL}{c\gamma_{ge}} \gg 1. \qquad (12)$$

Note that this condition contains a total number of atoms N, which is a signature of so-called "collective enhancement."

The concept of "polariton" is well known in condensed matter physics. For instance, interaction of excitons with light in semiconductor materials is often described in terms of exciton polaritons (Yamamoto and Imamoglu, 1999). Already early experiments demonstrated that the dispersion relation for exciton polaritons can dramatically alter group velocities (Ulbrich and Fehrenbach, 1979). The polariton picture of Raman adiabatic passage was first discussed by Mazets and Matisov (1996). The finite bandwidth associated with EIT is the basis of the pulse matching effect studied by Harris (1993), and by Fleischhauer and Manka (1996).

D. From EIT to quantum memory: the idea

The idea for quantum memory is closely related to the dark-state polariton concept. When a polariton propagates in an EIT medium, it preserves its amplitude and shape,

$$\hat{\Psi}(z,t) = \hat{\Psi}\left[z - c\int_0^t d\tau\cos^2\theta(\tau), t=0\right], \qquad (13)$$

but its properties can be modified by simply changing the intensity of the control beam. As the control intensity is decreased, $\cos^2\theta \sim |\Omega|^2$ becomes very small, and the group velocity ($v_g = c\cos^2\theta$) is slowed. At the same time the contribution of photons in the polariton state is reduced, see Eq. (8). In particular, if the control beam is turned off [$\Omega(t) \to 0$] the polariton's group velocity is reduced to zero [$\theta(t) \to \pi/2$] and it becomes purely atomic:

$$\hat{\Psi}(z,t) \to -\sqrt{N}\,\hat{s}(z). \qquad (14)$$

At this point, the photonic quantum state is mapped onto long-lived spin states of atoms. As long as the trapping process is sufficiently smooth [i.e., adiabatic (Oreg *et al.*, 1984)], the entire procedure has no loss and is completely coherent. The stored quantum state can be easily retrieved by simply reaccelerating the stopped polariton. This is illustrated in Fig. 3, which shows the evolution of a "signal" light pulse, spin coherence, and polariton when the control beam is turned off and on. The amplitude of the signal pulse decreases as it is being decelerated whereas the spin coherence grows; the procedure is reversed when the control beam is turned back on.
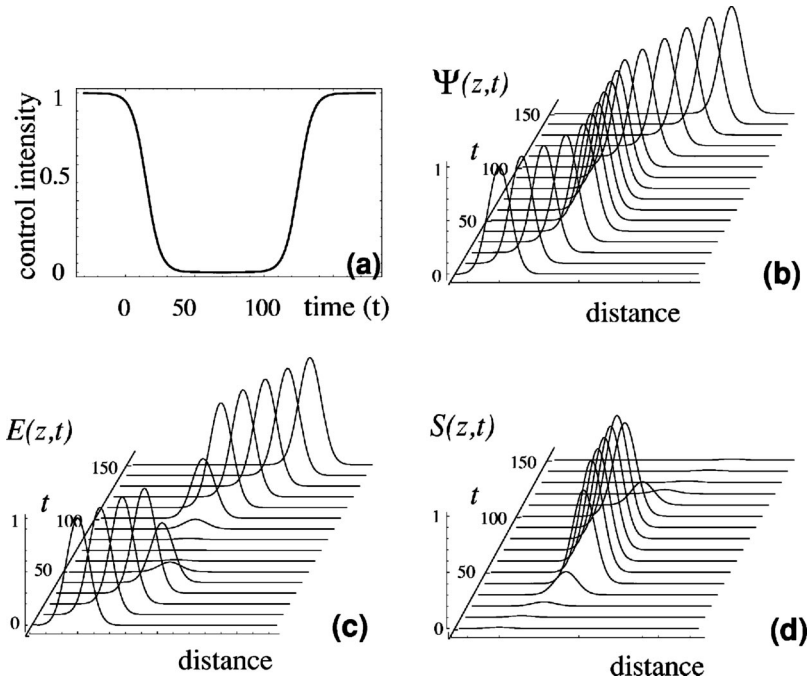
FIG. 3. A dark-state polariton can be stopped and reaccelerated by ramping the control field intensity, as shown in (a). The coherent amplitudes of the polariton Ψ, the electric field E, and the spin components s are plotted in (b)–(d).

We have already argued that for EIT to be effective in eliminating dissipation, the light pulse spectrum should be contained within a relatively narrow transparency window [Fig. 1(b)]. A vanishing control beam intensity implies that the transparency window would become infinitely narrow and eventually disappear. How can one avoid loss in such a case? The essence of adiabatic following in polaritons is that a dynamic reduction in group velocity is accompanied by narrowing of the polariton frequency spectrum, such that it is not destroyed even if $v_g = 0$. To see why it happens, we note that during the process of adiabatic slowing the *spatial* profile and, in particular, the spatial width of the wave packet remains unaffected (see Fig. 3), as long as the group velocity $v_g(t)$ is only a function of time (Fleischhauer and Lukin, 2002). At the same time, the amplitude of the electric field gets reduced and its temporal profile is stretched due to the reduction of the group velocity.

The spectrum of the signal field is reduced in proportion to $v_g/c \sim |\Omega|^2$, i.e., by exactly the same factor as the transparency bandwidth $\Delta\nu$. Therefore, the conditions for adiabatic following are very simple: the entire pulse should be within the medium at the beginning of the trapping procedure, and its spectrum should be contained within the original transparency window. Once again, these conditions are satisfied only if an optically dense medium (12) is used. It is also worth noting that the rate at which the group velocity is turned to zero can be quite fast, especially if the initial group velocity of the light pulse is much smaller than c. The adiabaticity conditions have been analyzed in detail by Matsko *et al.* (2001a), and by Fleischhauer and Lukin (2002).

The concept of adiabatic passage in multilevel systems was first introduced by Oreg *et al.* (1984) and was experimentally rediscovered by Gaubatz *et al.* (1990). Its application for quantum state transfer was first pointed out by Parkins *et al.* (1993). Extensions and detailed analysis of such techniques were considered by Parkins and Kimble (1999). Recent experimental progress toward implementation of these ideas (Kuhn *et al.*, 2002) should be especially noted. Csesznegi and Grobe (1997) pointed out that the spatial profile of an atomic Raman coherence can be mirrored into the electromagnetic field by coherent scattering, whereas time-varying fields can be used to create spatially nonhomogeneous matter excitations. These techniques were reviewed by Bergmann *et al.* (1998). There is by now a considerable literature investigating various aspects of storage in atomic ensembles (Juzelinas and Carmichael, 2002; Mewes and Fleischauer, 2002) as well as nonclassical light generation (Poulsen and Molmer, 2001) using these techniques. See also the review by Fleischhauer and Mewes (2001).

Finally, it should be remarked here that the essential point of this technique is not to store the energy or momentum carried by photons but their quantum states. In fact, in practice almost no energy or momentum is actually stored in the EIT medium. Instead, both are being transferred into (or borrowed from) the control beam in such a way that an *entire* optical pulse is coherently converted into a low-energy spin wave. This is the key feature that distinguishes the present approach from earlier studies in optics [involving, e.g., traditional photon echo techniques (Boyd, 1992) or nuclear physics (Shvydko *et al.*, 1996)], and that enables potential applications in quantum information science. A different proposal to "freeze" light pulses in a moving medium was suggested by Kocharovskaya *et al.* (2001) and the possibility to observe phenomena resembling black holes was considered by Leonhardt (2001).

E. Collective enhancement and stored states

The above considerations indicate that, in principle, complete storage and retrieval of the input state is pos-

sible. The key feature of the technique is that it involves atomic excitations in an optically dense, many-atom system, which allows one to make an efficient and robust quantum state transfer. Specifically, the requirement of adiabatic following can be fulfilled as long as the medium is optically dense, i.e., a large number of atoms interact with the light. This is a specific manifestation of the collective enhancement: the rate of coherent coupling between photons and matter increases with the total number of atoms N. It is natural to wonder what the corresponding scaling for decoherence processes is.

A particular feature associated with an ensemble approach is that here it is impossible to determine which particular atom absorbs a photon. This implies that the signal pulse couples to collective spin states. As before, we assume that initially all atoms are in the ground state, $|g.s.\rangle \equiv |g_1, \ldots, g_N\rangle$. If the field was initially in the vacuum state, then no change will take place for the atomic states. However, if the field was initially in a Fock state with a single photon, $|1\rangle$, then at the end of the trapping procedure one spin will be flipped. However, it is impossible to determine which particular atom will be spin flipped. To illustrate the concept, consider the situation in which the excitation is delocalized over a certain volume in which each of N atoms has an equal probability of absorption. The state that will result from the trapping procedure is a symmetric collective state

$$|1\rangle|g.s.\rangle \rightarrow |0\rangle \sum_{i=1}^{N} \frac{1}{\sqrt{N}} |g_1, \ldots, s_i, \ldots, g_N\rangle. \qquad (15)$$

More generally, the ideal storage procedure will transform any superposition of photon states into corresponding superpositions of atomic spin states:

$$\sum_i \alpha_i |i\rangle|g.s.\rangle \rightarrow \sum_i \alpha_i |0\rangle|\mathbf{s}^i\rangle, \qquad (16)$$

where $|i\rangle$ are number states of the photon field and $|\mathbf{s}^i\rangle$ are collective states containing i flipped spins. The state of the type given by Eq. (15) is a complex state of N particles, and one might wonder if those will be fragile with respect to decoherence and losses. This is not so. For example, if one atom is lost, the resulting state coincides almost exactly (with an error that scales as $1/N$) with a symmetrized $N-1$ atom state. This implies that the characteristic decoherence rate of the single photon storage state is equal to the single particle decay rate γ_{gs}. Note that the effect of loss can be very different for other kinds of many-particle states, such as Greenberger-Horne-Zeilinger-type entangled states ($\sim |g_1, \ldots, g_N\rangle + |s_1, \ldots, s_N\rangle$), in which the decoherence rate is proportional to N (Bouwmeester *et al.*, 2000). Hence a remarkable feature of the special type of collective states [sometimes also called W states (Dur

et al., 2000)] used for storage is that it is very robust with respect to decoherence and particle loss.[3]

We have already noted that the idea of using atomic ensembles with light to prepare nontrivial states of atoms was first originated and explored in the context of absorptive interactions (Grom and Kuzmich, 1995). Off-resonant light can also interact with atoms by means of the refractive index (i.e., dispersively) in which case, for instance, its polarization state can be modified (Duan *et al.*, 2000; Kuzmich and Polzik, 2000). In such a dispersive interaction photons can become correlated with atoms. Although such interactions are weaker than those at resonance, the collective enhancement still allows one to perform nontrivial quantum state manipulations, such as quantum nondemolition measurements of collective atomic spins (Kuzmich *et al.*, 2000). Recently these ideas have led to the striking experimental demonstration of entanglement of two macroscopic samples of atoms (Julsgaard *et al.*, 2001). Likewise, such dispersive interactions could be used to facilitate quantum state exchange between atoms and continous wave (cw) fields in nonclassical states. Recent progress towards implementing these ideas is noteworthy (Schori *et al.*, 2002). Finally we point to the recent theoretical studies on decoherence properties of collective states (Dur *et al.*, 2000).

F. Experimental demonstrations

Recent experiments have already demonstrated some of the light manipulation effects described above by showing that the weak pulses can be "trapped" and released after some storage interval. It must be emphasized that all of the experiments carried out so far were in the classical domain, since they involved weak laser pulses in coherent states. Hence, these experiments did not explore what is probably the most interesting feature of the quantum memory technique (i.e., storage of nonclassical correlations of light). Nevertheless, they are able to probe some important features of the dynamic trapping in polaritons, showing in particular that the weak pulses are not destroyed during the trapping and release processes and demonstrating a phase coherence of the entire procedure.

The experiment of Liu and co-workers (Liu *et al.*, 2001) used a setup similar to that of the earlier slow light demonstration. A light pulse has been slowed and then "trapped" in an ultracold atomic sample slightly above the point of quantum degeneracy for up to 1.5 ms. Another experiment (Phillips *et al.*, 2001) involved Zeeman sublevels of the electronic ground state of warm Rb atoms and used polarization control techniques to manipulate them. In order to ensure long coherence times a buffer gas was used to effectively slow the atomic mo-

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[3]In the case when the incident field is in a coherent photon state, a coherent spin state is created. The latter can be factorized into a product of individual atom states, in agreement with the usual semiclassical theory.
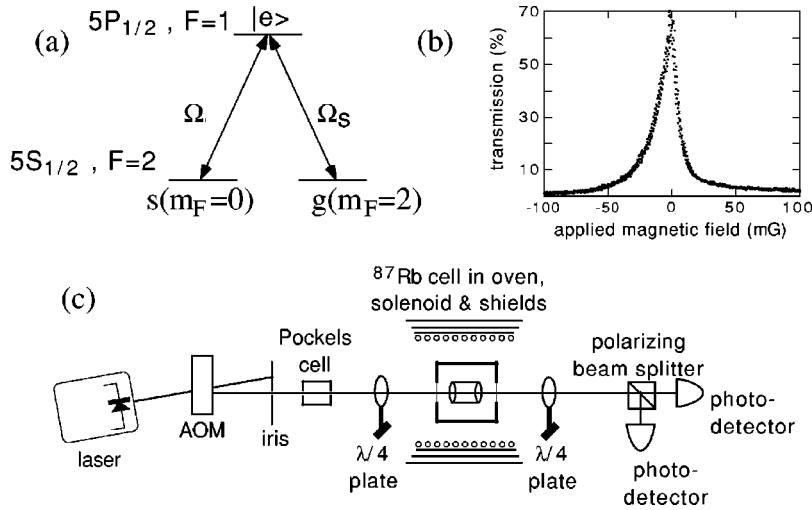
FIG. 4. Experiment by Phillips *et al.* (2001): (a) Λ-type configuration of ^{87}Rb atomic states resonantly coupled to a control field (Ω) and a signal field (Ω_s). (b) A typical observed Rb Faraday resonance in which the transmission intensity for a cw signal field is shown as a function of magnetic field. The full width of this resonance is 20 mG, which corresponds to a 15-kHz shift in the Zeeman levels. (c) Schematic of the experimental setup.

tion, similar to the earlier studies (Kash *et al.*, 1999). This realization requires a very simple experimental arrangement, possibly adaptable to an advanced undergraduate laboratory, that we describe next.

In the experiment of Phillips *et al.* (2001) the control field and signal field were represented, respectively, by the two helicities of circularly polarized light (σ_+ and σ_-) derived from a single laser beam by carefully controlling the light polarization, as shown in the experimental schematic in Fig. 4(c). These light fields couple pairs of Zeeman hyperfine sublevels of electronic ground-state ($5^2S_{1/2}$) Rb atoms ($|s\rangle, |g\rangle$), with magnetic quantum numbers differing by two, via the excited ($5^2P_{1/2}$) state [cf. Fig. 4(a)]. In this configuration the effective two-photon detuning can be controlled by applying an external magnetic field, which causes Zeeman level shifts between $|s\rangle$ and $|g\rangle$. For the Rb data presented here the $5^2S_{1/2}, F=2 \rightarrow 5^2P_{1/2}, F=1$ transition in ^{87}Rb was employed. The control field was always much stronger than the signal field; hence most of the relevant atoms were in the $5^2S_{1/2}, F=2, M_F=+2$ magnetic sublevel. In this case the states $|s\rangle, |g\rangle$ of the simplified three-level model correspond, respectively, to $|F=2, M_F=0\rangle$ and $|F=2, M_F=+2\rangle$. By using a fast Pockels cell we slightly rotated the polarization of the input light to create a weak pulse of σ_- light, which served as the signal field. Figure 4(b) displays a typical transmission spectrum for the signal (σ_-) field obtained by scanning the magnetic field and thereby changing the effective two-photon detuning. Outside of the transparency window (magnetic fields >20 mG) the Rb vapor is completely opaque to σ_- light. Rb vapor experiments were carried out at temperatures $\sim$70–90 °C, which correspond to atomic densities $\sim 10^{11}$–10^{12} cm^{-3}. For comparison, in experiments involving ultracold Na atoms, hyperfine sublevels were used with control and signal beams separated in frequency by about 1.7 GHz.

We turn next to a demonstration of the storage process. Typical input σ_- signal pulses had a temporal length of $\sim$10–30 μs. Upon entrance into the Rb cell the signal pulse was spatially compressed to a length of a few centimeters due to the reduction in group velocity.

In order to trap, store, and release the signal pulse, we used an acousto-optic modulator to turn off the control field smoothly over about 3 μs while much of the signal pulse was contained in the Rb cell. After some time interval, we turned the control field on again, thereby releasing the stored portion of the signal pulse.

An example of the observed storage process is shown in Fig. 5. Typically, two time-resolved σ_- signal pulses were registered by the photodetector. First, a fraction of the signal pulse left the cell before the control field was turned off, which resulted in an observed signal that was not affected by the storage operation (peak I in each plot of Fig. 5). This untrapped light was delayed by about 30 μs as compared to free-space propagation due to the slow group velocity ($v_g \sim 1$ km/s). The second observed signal pulse was light that was stored in atomic excitations for a time interval τ. Note that the released signal pulse was detected only after the control field was turned back on (peak II in each plot of Fig. 5). We observed that the amplitude of the released signal pulse decreased as the τ increased. We could resolve released light pulses without signal averaging for storage intervals up to $\tau \simeq 0.5$ ms.

An important feature of the storage technique that distinguishes it, e.g., from the usual photodetection, is its coherence properties. The coherence properties were probed in the experiments of Mair *et al.* (2002). There we applied a pulsed magnetic field during the light storage interval, to vary controllably the phase of the Zeeman coherence in the Rb vapor. We then converted the spin excitations back into light and detected the resultant phase shift in an optical interferometric measurement. To form an interferometer for the two fields, we adjusted the $\lambda/2$ plate such that a small fraction (<10%) of the control field was mixed into the signal detection channel.

By adjusting the applied magnetic field pulse during the storage, we could easily modify the phase Φ of the coherent excitation. If a pulsed magnetic field $B_z(t)$ is applied in a direction parallel to that of the light propagation, then the Zeeman sublevels are differentially shifted in energy, producing a phase shift in the Zeeman coherence given by
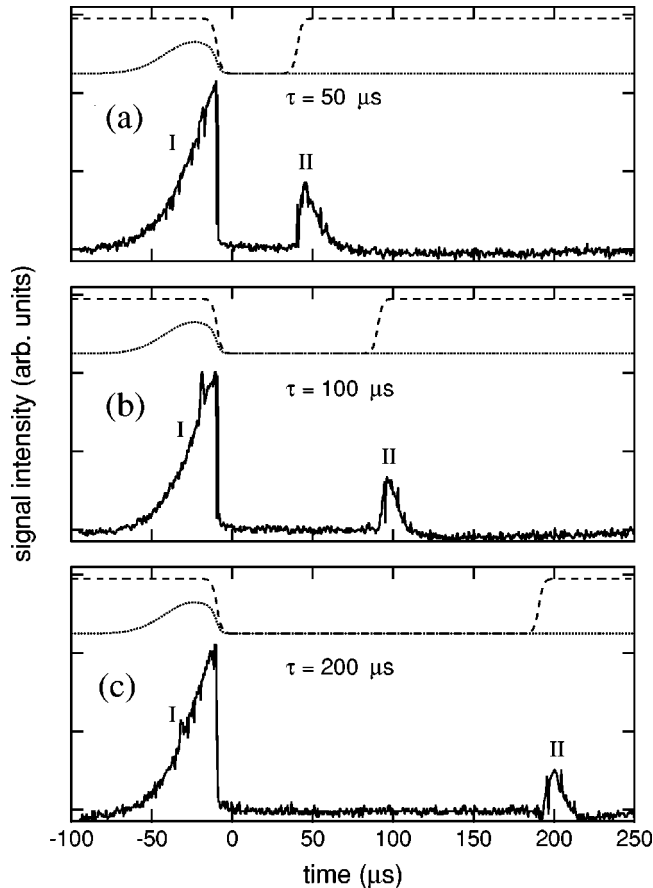
FIG. 5. Observed light pulse storage in a ^{87}Rb vapor cell. Examples are shown for storage times of (a) 50 μs, (b) 100 μs, and (c) 200 μs. (Background transmission from the control field, which leaks into the signal field detection optics, has been subtracted from these plots.) Shown above the data in each graph are calculated representations of the applied control field (dashed line) and input signal pulse (dotted line).

$$\Phi = (g_g - g_s)\frac{\mu_B}{\hbar}\int_0^T B(t')dt', \tag{17}$$

where g_g and g_s are Landé factors corresponding to different Zeeman states and T is the time during which the magnetic field is applied. Figure 6 shows 20 stored light experiments for which we increased the Zeeman phase shift by approximately 0.2π for each successive run. Trace A in Fig. 6 shows the result for $\Phi \simeq 0$ and hence maximum constructive interference between the output signal light and the control field. As we increased the pulsed magnetic field to change the phase by π, we observed destructive interference (e.g., trace B). As we increased the pulsed magnetic field still further, we alternatively observed constructive and destructive interference (as expected) at $\Phi \simeq 2\pi, 3\pi, 4\pi$, etc. (traces C–E in Fig. 6). We observed up to 10 periods of phase accumulation (i.e., $\Phi \simeq 20\pi$) without loss of coherence.

Similar effects have been observed in several other related experiments demonstrating releasing the pulse in the direction opposite to the original (Zibrov *et al.*, 2002) and investigating the effects of detuning (Kuzuma *et al.*, 2002). These effects have also been observed in
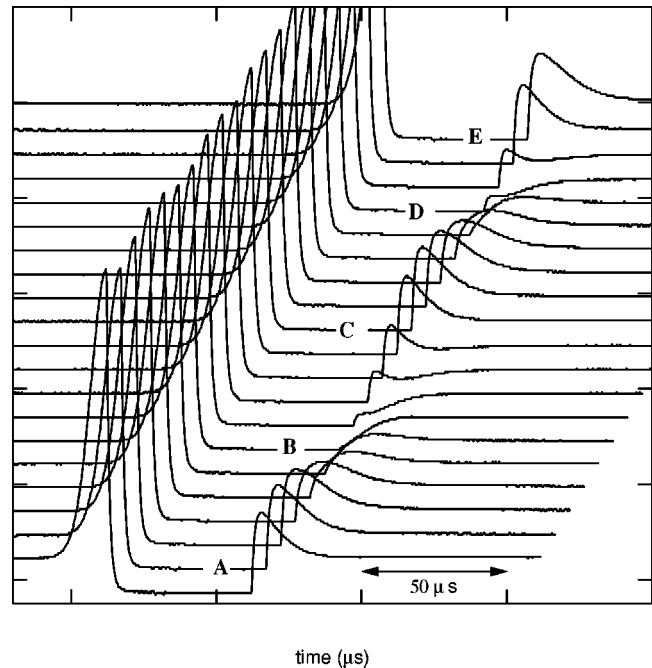

FIG. 6. Results of interferometric measurements of released photonic excitations. Curves represent 20 light storage experiments. In each experiment the magnetic field was pulsed during the storage interval with increasing strength from trace A to E such that the accumulated phase difference [Φ from Eq. (17)] varied from approximately 0 to 4π. (See text for details.) Note that there is a small phase offset at zero pulsed magnetic field, caused by the Pockels cell.

solid-state media (Turukhin *et al.*, 2002). For a recent review, see Matsko *et al.* (2001b). Work on probing the quantum aspects of the storage technique is currently under way.

III. QUANTUM STATE MANIPULATION IN ATOMIC ENSEMBLES

A. Motivation and overview

The above ideas for quantum storage are based on a notion of long-lived coherences involving spin states of atomic ensembles. This also implies that the interactions between atoms in such states should be sufficiently weak. Furthermore, the storage procedure described above is, in essence, a linear optical technique. At the same time, one is naturally interested in developing techniques for nontrivial manipulation of stored quantum states. As a rule, such techniques always involve strong interactions, or, in other words, large nonlinearities. In this section, we outline several approaches for the coherent manipulation of collective quantum states.

The idea of using atoms or ions for controlled quantum state manipulation and entanglement is very attractive, in view of very long coherence times and the well-developed techniques for cooling and trapping. It is worth emphasizing that perhaps the most impressive progress in the entire area of experimental quantum information has been achieved in systems of trapped ions

(Cirac and Zoller, 1995; Monroe *et al.*, 1995; Sackett *et al.*, 2001). The most recent work, in particular, indicates good prospects for scalability (Row *et al.*, 2002). At the same time, a number of spectacular experiments on controlled nonclassical state generation and entanglement has been carried out using neutral atoms in Rydberg states and high-Q microwave cavities (Varcoe *et al.*, 2000; Raimond *et al.*, 2001). The main obstacle for scaling of this system to many atoms and operation is due to the absence of deterministic source of single Rydberg atoms. (The "dipole blockade" technique described shortly can in fact be used for such a source.[4]) Several ideas for using confined neutral atoms for quantum computation have also been put forward (Brennen *et al.*, 1999; Jaksch *et al.*, 1999, 2000; Lukin and Hemmer, 2000). Some very recent fascinating experiments have made definitive steps toward achieving a required degree of quantum control (Schlosser *et al.*, 2001; Orzel *et al.*, 2001; Greiner *et al.*, 2002). At the same time active searches are now under way for new techniques that could alleviate stringent requirements on quantum manipulation schemes.

B. Mesoscopic phenomena in small atomic samples: dipole blockade

One of the most straightforward ways to induce strong interactions between atoms is to promote one of the spin states of atoms into highly excited electronic states (such as, e.g., Rydberg states) with a resonant laser field for a well-defined period of time. Atoms in such Rydberg states have a large size and can therefore have large dipole moments resulting in strong atom-atom interaction (Gallagher, 1994). The resulting energy shift scales as $U \sim \wp^2/R^3$, where $\wp$ is the dipole moment and R is the separation between atoms, and this interaction can be used to entangle atoms (Jaksch *et al.*, 2000). In addition, the single-atom decoherence rate associated with spontaneous emission from the Rydberg state and interaction with black-body radiation can be quite low. However, the applications of these ideas for controlled quantum state manipulation are challenging to implement. First, the precise strength of the interaction (and therefore the resulting state) is very sensitive to the interatomic separation. Therefore, controlling the motion and the position of individual atoms to submicron precision is required. Second, during the excitation period, the large interactions are accompanied by strong mechanical forces on the trapped atoms. Thus, the internal states of the trapped atoms become entangled with the motional degrees of freedom, resulting effectively in an additional source of decoherence. As a result, extremely tight confinement of atoms and ultralow temperatures are required. Such experimental problems most probably preclude the direct applications of the above ideas

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[4]The ideas along this direction have been independently developed by the ENS group of M. Brune, J.-M. Raimond, and S. Haroche.
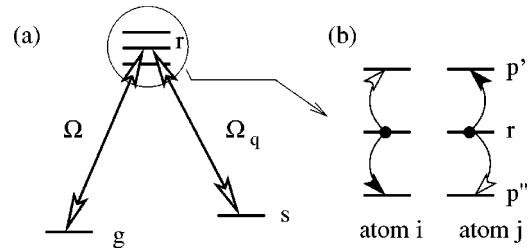

FIG. 7. Dipole blockade in Rydberg atoms: (a) Relevant single-atom energy levels and laser fields for dipole blockade. (b) Excitation exchange due to resonant dipole-dipole interactions.

for manipulating many-atom ensembles. At least some of these problems however can be mitigated using the techniques that we now describe.

Consider an ensemble of N identical multistate atoms (Fig. 7) distributed randomly in a volume V in which all atoms are initially trapped and prepared in a specific sublevel ($|g\rangle_i$, $i = 1, \ldots, N$) of the ground-state manifold. We assume modest atomic densities, such that interactions between atoms can safely be neglected whenever they are in the sublevels of the ground state. This also implies long coherence lifetimes—up to a few seconds—associated with these storage sublevels. However, when excited into the Rydberg states $|r\rangle_i$ the atoms interact strongly with each other due to the presence of resonant dipole-dipole interactions, resulting in the energy shifts of doubly excited states. The key idea is to use these interactions and associated level shifts to block the optical transitions into states with more than a single excitation. In this case, an entire ensemble can behave as a mesoscopic few-level system. This is the essence of the "dipole blockade" phenomenon (Lukin *et al.*, 2001). Before proceeding we note that this phonemenon closely resembles similar mesoscopic effects that are extensively studied in nanoscale solid-state devices (Altshuler *et al.*, 1991). It is remarkble that the tools and techniques from atomic physics can now be used to probe similar physics.

Figure 7 illustrates some of the basic properties of this technique for a relevant subset of Rydberg states. As a specific interaction mechanism we consider excitation "hopping" via resonant dipole-dipole interactions between Rydberg atoms. Suppose that the energy separation between the optically excited Rydberg state $|r_i\rangle$ and the pair of the sublevels of different parity $|p_i'\rangle$, $|p_i''\rangle$ is adjusted (using, e.g., electric fields) such that the transition energies are equal, $E_r - E_{p'} = E_{p''} - E_r$, as shown in Fig. 7(b). In this case, any pair of atoms excited in the $|r_i\rangle|r_j\rangle$ states would undergo a hopping transition into the states $|p_i'\rangle|p_j''\rangle$ and $|p_i''\rangle|p_j'\rangle$. Such a transition corresponds to exchange of virtual (microwave frequency) photons and results in a splitting of the two-atom states.

The relevant process is described by the Hamiltonian

$$\hat{V}_d = \hbar \sum_{i>j} \kappa_{ij} |r_i\rangle |r_j\rangle (\langle p_i'| \langle p_j''| + \langle p_j'| \langle p_i''|) + \text{H.c.}, \quad (18)$$

where $\hbar \kappa_{ij} \sim \wp_{rp'} \wp_{rp''}/r_{ij}^3$, with $\wp_{kl}$ being the dipole matrix elements for the corresponding transitions and r_{ij} the distance between the two atoms. In general, this in-
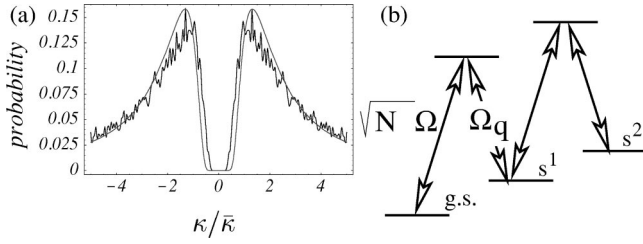
FIG. 8. Physics of dipole blockade: (a) Energy gap in a collective spectrum. (b) Manipulation of multiple atomic excitations.

teraction does not affect the singly excited collective states (e.g., $\hat{V}_d|\mathbf{r}^1\rangle=0$) but leads to a splitting of the levels when two or more atoms are excited. In particular, the collective eigenstate with atoms i and j being excited and all other atoms in the ground or storage states is

$$
\begin{aligned}
|\pm^2_{\mathbf{ij}}\rangle = \frac{1}{\sqrt{2}}[&|g_1,\ldots,r_i,\ldots,r_j,\ldots,g_N\rangle \\
&\pm(|g_1,\ldots,p'_i,\ldots,p''_j,\ldots,g_N\rangle \\
&+|g_1,\ldots,p'_j,\ldots,p''_i,\ldots,g_N\rangle)/\sqrt{2}],
\end{aligned}
\tag{19}
$$

the energies of the $\pm$ states being split by $\hbar\kappa_{ij}$. It follows that for an ensemble contained in a finite volume V, the manifold of doubly excited states has an energy gap of order $\hbar\bar{\kappa}=\wp_{rp'}\wp_{rp''}/V$ [Fig. 8(a)]. Physically this gap corresponds to the minimal interaction energy of any pair of excited Rydberg atoms confined in a volume V. Thus if $\bar{\kappa}$ is much larger than the linewidth γ_r of the Rydberg state, resonant excitations from the singly to the doubly excited states are strongly suppressed. This is the essence of the dipole blockade. Note that for a system confined to a volume of $\sim 10~\mu\text{m}^3$ and excited to Rydberg states ($n\sim 50$), typical dipole-dipole interactions correspond to an energy gap in the range of 100 MHz, whereas coherence times limited by radiative relaxation and black-body radiation correspond to γ_r <10 kHz.

The remarkable consequence of the "dipole blockade" is that the states that are involved in strong atomic interactions (i.e., the states with two and more atoms excited) are never populated. Hence, the present approach avoids mechanical interaction between atoms and leaves the qubit degrees of freedom decoupled from atomic motion. Finally, the precise magnitude of the level shift (which depends upon the specific atomic configuration) is not important as long as it is large enough to inhibit transitions. Since the atomic motion is never coupled to the qubits, the temperature is only limited by the requirement that the atomic distribution should not change significantly on the time scale of the gate operation T (the "frozen" gas approximation). This is the case even for temperatures as high as a few mK.

The manipulation of atoms in an ensemble can be accomplished by light fields of different frequencies and polarizations, which illuminate the entire ensemble and excite all atoms with equal probability. As a consequence only symmetric collective states are involved in

the process. As before, we label these symmetrized states as $|\mathbf{s}^{n_s},\mathbf{r}^{n_r}\rangle$ with n_s and n_r being the number of atoms in states $|s\rangle$ and $|r\rangle$, respectively. Consider first the situation in which the optical field with Rabi frequency Ω is tuned to the transition from the ground state to the Rydberg sublevel r. If the atoms were noninteracting, then each of them would undergo independent Rabi oscillations, thereby transferring all atoms from the ground state to the excited state and back with a period of $2\pi/\Omega$. In the case when there is a strong interaction between Rydberg atoms, a light field tuned to the single-atom resonance frequency can excite the transition between the ground state and the first collective state $|\mathbf{r}^1\rangle$. But when the splitting of the states $|\pm^2_{\mathbf{ij}}\rangle$ is large, successive transitions into these higher states are strongly inhibited. Hence, if the atomic system is initially in its ground state $|\mathbf{g}\rangle$, the evolution is given by the two-level dynamics

$$
\begin{bmatrix} |\mathbf{g}(t)\rangle \\ |\mathbf{r}^1(t)\rangle \end{bmatrix} = \begin{bmatrix} \cos\theta(t) & -i\sin\theta(t) \\ i\sin\theta(t) & \cos\theta(t) \end{bmatrix} \begin{bmatrix} |\mathbf{g}(0)\rangle \\ |\mathbf{r}^1(0)\rangle \end{bmatrix},
\tag{20}
$$

where $\theta(t)=\sqrt{N}\int_0^t\Omega(\tau)/2\,d\tau$. The ensemble displays Rabi oscillations only between the ground and the singly excited Rydberg state with a collective Rabi rate $\sqrt{N}\Omega$.

The dipole blockade technique can be used to generate superpositions of collective spin states in an ensemble and to perform quantum gate operations between qubits encoded in such ensembles. For example, using the ideas described above, the system can be driven into superpositions of collective states $\alpha_0|\mathbf{g}\rangle+\alpha_1|\mathbf{r}^1\rangle$. The single-quantum excitation can now be stimulated into a storage sublevel (e.g., $|\mathbf{s}^1\rangle$) by a π pulse $[\int\Omega_q(\tau)d\tau=\pi]$ and we can associate a qubit with the state $\alpha_0|\mathbf{g}\rangle+\alpha_1|\mathbf{s}^1\rangle$. As illustrated in Fig. 8(b), this procedure can be generalized to the generation of higher-order collective states: one can prove that a synthesis of arbitrary superpositions $|\boldsymbol{\Psi}_n\rangle=\Sigma_{m=0}^n\alpha_m|\mathbf{s}^m\rangle$ ($n\leq N$) is possible using a sequence of properly timed pulses (Lukin *et al.*, 2001). Entanglement of excitations stored in different ensembles can be done by first transferring them into photons, followed by subsequent trapping into distinct hyperfine states of a single ensemble. A conditional phase shift can then be induced via controlled excitation of different sublevels into the Rydberg states.

The ideas of dipole blockade originate from earlier work on manipulation and entanglement of single atoms (Brennen *et al.*, 1999; Jaksch *et al.*, 2000; Lukin and Hemmer, 2000). Related ideas involving dipolar molecules have also been recently proposed by DeMille (2002). Ideas for using blockade mechanisms for spin squeezing (Bouchoule and Molmer, 2002), robust generation of nontrivial atomic states (Unanyan and Fleischhauer, 2002), and single atom/photon sources (Saffman and Walker, 2002) have also been discussed recently.
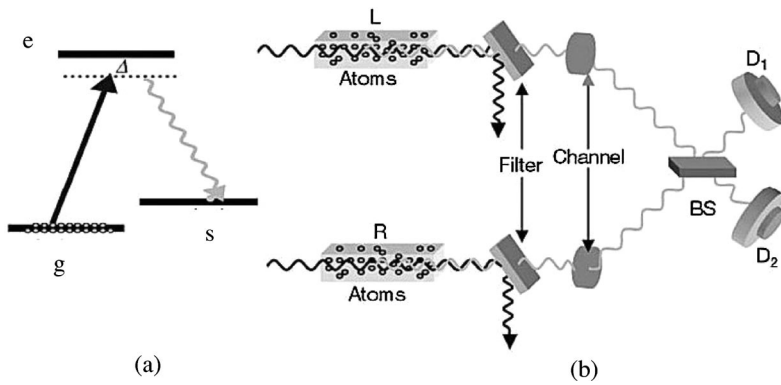
FIG. 9. Setup for generation of nonclassical ensemble states using Raman scattering. (a) The relevant level structure of the atoms in the ensemble. (b) Schematic setup for generating entanglement between the two atomic ensembles L and R.

C. Atomic state manipulation using atom-photon correlations

The "dipole blockade" technique described above applies to a system with a strongly nonlinear response, which allows one to control the state of the atomic ensemble at a level of single quanta of excitation. In some sense this represents an example of "ultimate" control over a quantum system. Although one can hope that this, or related techniques of this kind, may some day become practical, experimental efforts along these lines are only beginning.

At the same time, many examples of nonlinear phenomena in various areas of physics, and in particular in optics, are well known. As a rule, however, nonlinearities are not sufficiently strong to have an observable effect at the level of single quanta. It is natural to ask if phenomena of this kind can be used for efficient entanglement manipulation. It turns out that this is the case, although the useful manipulations are probabilistic in nature.

We now present an example illustrating how such a weakly nonlinear process [Raman scattering (Raymer *et al.*, 1985)] can be used for manipulation of atomic quantum states (Duan *et al.*, 2001; Andre *et al.*, 2002). We then indicate how these ideas can be used to create a robust entanglement of atomic ensembles via realistic (i.e., absorbing) channels. The latter is especially important for applications involving long-distance quantum communication.

We consider again a cloud of identical three-level atoms (Fig. 9) that are initially prepared in the ground state $|g\rangle$. A sample is excited by an off-resonant laser pulse that induces a two-photon Raman transition into the states $|s\rangle$, corresponding to flipped atomic spins that are accompanied by emission of the so-called Stokes photons (wavy arrow in Fig. 9). It is important to emphasize that atomic spin flips and Stokes emission events are correlated: for each emitted photon, there is a corresponding flipped spin. We are particularly interested in the forward-scattered Stokes light that is copropagating with the laser. A photon emitted in this mode is uniquely correlated with the excitation of the symmetric collective spin-wave mode S, given by $S \equiv (1/\sqrt{N_a})\Sigma_i |g\rangle_i\langle s|$, where the summation is taken over all the atoms. To be specific, let us assume that the pump pulse duration t_p is short so that the mean photon number in the forward-

scattered Stokes pulse is much smaller than 1. The state of the system after the pulse can be written in the following form:

$$|\phi\rangle = |\mathbf{g}\rangle|0\rangle + \sqrt{p_c}S^\dagger a^\dagger|\mathbf{g}\rangle|0\rangle + o(p_c), \qquad (21)$$

where $p_c = 4g^2NL/c|\Omega_p|^2/\Delta^2 t_p$ is the small excitation probability and $o(p_c)$ represents terms with more excitations, whose probabilities are equal to or smaller than p_c^2. Here Ω_p is a Rabi frequency of the pump field and Δ is a single-photon detuning. Note that the excitation probability is proportional to the number of atoms (i.e., the product of the linear density and the length) in the ensemble.

Equation (21) indicates that whenever a single Stokes photon propagating in the forward direction is detected, the state of the atomic ensemble will be given by $S^\dagger|0_a\rangle$. That is, the detection of the single photon "projects" the ensembles into a nonclassical state with a single quantum in a well-defined spin-wave mode. This is an example of probabilistic quantum state manipulation. Before proceeding, we remark that a large fraction of light is always emitted in directions other than forward, due to the spontaneous nature of the process. However, whenever the number N of the atoms is large, such events will mostly populate different spin-wave modes, whereas the contribution to the population in the symmetric collective mode will be small. As a result, the use of atomic ensembles results in a large signal-to-noise ratio, which enhances the efficiency of the scheme. The stored spin-wave quanta can be released by conversion into the photon wave packet with well-defined properties, such as shape and duration.

Figure 9(b) illustrates how the Raman scattering scheme described above can be used to generate entanglement between two atomic ensembles, left (L) and right (R), separated by an absorbing channel (Duan *et al.*, 2001). The two pencil-shaped ensembles are illuminated by the synchronized classical pumping pulses. The forward-scattered Stokes pulses are collected and coupled to optical channels (such as fibers) after the filters, which are polarization- and frequency-selective filters of the pumping light. The pulses after the transmission channels interfere at a 50%-50% beam splitter, with the outputs detected, respectively, by two single-photon detectors D_1 and D_2. The basic idea is that in such a configuration a detector click implies that one quantum

of spin excitation has been created in two ensembles, but it is fundamentally impossible to determine from which of two ensembles the photon arrived. In this case, the measurement projects the state of the system into an entangled state of two ensembles.

After two laser pulses excite both ensembles, the whole system is described by the product state for the left and the right subsystems: $|\phi\rangle_L \otimes |\phi\rangle_R$, where $|\phi\rangle_L$ and $|\phi\rangle_R$ are given by Eq. (21) with all the operators and states distinguished by the subscript L or R. When the forward scattered Stokes light from both ensembles is combined at the beam splitter and a photodetector click in either D_1 or D_2 measures the combined radiation from two samples: $a_+^\dagger a_+$ or $a_-^\dagger a_-$ with $a_\pm = (a_L \pm a_R)/\sqrt{2}$. Depending on the detector click, we find the projected state of the ensembles L and R is a maximally entangled state of the form

$$|\Psi_\varphi\rangle_{LR}^\pm = (S_L^\dagger \pm e^{i\varphi} S_R^\dagger)/\sqrt{2}|0_a\rangle_L|0_a\rangle_R. \tag{22}$$

The most remarkable feature of the above process is that it can be made robust with respect to imperfections and losses during the optical propagation. In particular, when the total losses in both left and right optical paths are equal, the resulting state will still always be given by (22); the loss will affect only the overall probability to get a click.[5] Hence the techniques for probabilistic manipulation of the atomic ensembles may have an interesting application for long-distance quantum communication in realistic (lossy) photonic channels, where absorption leads to the exponential loss of signal after propagation over long distances. The idea described above illustrates a general principle of the so-called quantum repeater (Briegel *et al.*, 1991)—a technique for correcting errors (Bennett *et al.*, 1991; Knill *et al.*, 1998) in quantum commmunication.

Before concluding we note that there exists a large body of literature on probabilistic quantum information processing using linear optics (see, e.g., Knill *et al.*, 2001; Pan *et al.*, 2001; Franson *et al.*, 2002). In a related context, it was recently suggested that photon states stored in ensembles can be used for efficient photodetection (Imamoglu, 2002; James and Kwait, 2002). Very recently, the idea of storing one photon from a parametric down converter in a fiber loop or ring resonator has also been discussed (Pittman *et al.*, 2002). The above scheme for entanglement generation originates from earlier proposals to entangle single atoms (Cabrillo *et al.*, 1990). An interesting connection to the studies of atom-photon correlations in quantum degenerate systems (Inouye *et al.*, 1999; Moore and Meystre, 2000) should be also noted. At the same time, the Raman excitation scheme

described above parallels theoretical and experimental studies on resonantly enhanced four-wave mixing (Hemmer *et al.*, 1995; Jain *et al.*, 1996; Lukin *et al.*, 1999). It was also shown theoretically that these techniques can be used for controlling quantum states of atoms (Andre *et al.*, 2002; Sorensen and Molmer, 2002).

IV. SUMMARY AND OUTLOOK

In summary, we have described a technique that allows one to transfer quantum states between light fields and metastable states of matter. The technique is based on trapping quantum states of photons in coherently driven atomic media, in which the group velocity is adiabatically reduced to zero. We have also outlined several approaches for manipulating quantum states of atomic ensembles. Specifically we discussed a "dipole blockade" phenomenon in which optical excitation of mesoscopic samples into Rydberg states can be used to control the state of ensembles at the level of individual quanta. At the same time we showed how atom-photon correlations in simple nonlinear optical processes (such as Raman scattering) can be used to effectively manipulate the ensemble states.

Before concluding, we note that a number of other interesting avenues are presently being explored that are closely related to the work described here. Controlled, coherent interactions between atoms form a basis for the rapidly developing field of nonlinear atom optics (Meystre, 2001; Anglin and Ketterle, 2002; Rolston and Phillips, 2002). Of particular relevance are effects involving strong correlations of atomic ensembles in traps (Pu and Meystre, 2000; Sørensen *et al.*, 2001) and optical lattices (Greiner *et al.*, 2002; Orzel *et al.*, 2001; Sørensen and Molmer, 1999) that produce nonclassical atomic states. Recently it has been shown that these systems open up very interesting prospects for studying complex phenomena involving quantum phase transitions in a highly controllable environment. Apart from quantum information and many-body science, a possible application of these techniques may involve atom interferometry with enhanced resolution (Holland and Burnett, 1993).

A large body of work on EIT has been motivated by developing new techniques for nonlinear optics (Harris *et al.*, 1990; Marangos, 1998; Lukin *et al.*, 2000). In particlular, resonantly enhanced photon-photon interactions enabled by EIT (Schmidt and Imamoğlu, 1996) could be useful for manipulating quantum states of light (Lukin and Imamoglu, 2000). It may also be possible to trap polaritons having a nonvanishing photonic component (André and Lukin, 2002) thereby creating very favorable conditions for nonlinear interactions. At the same time, "more traditional" applications such as frequency conversion and generation of intense, short pulses are being actively explored (Sokolov *et al.*, 2000). Very interesting efforts on using optical control of excitation in solid media should also be noted (Turukhin *et al.*, 2002).

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[5]Any asymmetry of the setup results in a correction to this state, including in particular, unequal amplitudes and phases. As shown by Duan *et al.* (2001), these errors can be efficiently corrected. We also note that the use of two-photon transitions involving Raman or Zeeman sublevels significantly reduces the sensitivity to phase errors due to, e.g., any uncertainty in position of the detectors.

Finally it should be emphasized again that even though the ideas and techniques for quantum control have already made an impact in various areas of physics and technology, the applications in quantum information science are still quite speculative in nature. Experiments in the coming years will undoubtedly shed new light on the opportunities and practical limitations for manipulating quantum states. Nevertheless, it is a very exciting area of physics and it seems reasonable to predict that at least some of the techniques described here are likely to find important applications, both in science and technology.

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