Beam splitting and Hong-Ou-Mandel interference for stored light

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Storing and release of a quantum light pulse in a medium of atoms in the tripod configuration are studied. Two complementary sets of control fields are defined, which lead to independent and complete photon release at two stages. The system constitutes a new kind of a flexible beam splitter in which the input and output ports concern photons of the same direction but well separated in time. A new version of Hong-Ou-Mandel interference is discussed.

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I. INTRODUCTION

Light propagation and storing in laser driven atomic media have recently become a subject of numerous studies, stimulated both by a fundamental character of the discussed problems and by possible applications. A generic system for such studies is a medium of three-level atoms interacting with two laser beams (signal and control) in the $\Lambda$ configuration. The essence of the process of light storing is a coherent mapping of the signal pulse into an atomic excitation, described by an atomic coherence, due to a switch-off of a control laser beam. By switching the control field again one restores the trapped pulse, preserving all the phase relations. This means in fact first writing down an information conveyed by the photons and later reading it back. The whole process can be effectively interpreted in terms of a joint medium+field excitation called dark state polariton. For recent reviews on light storage see, e.g., Refs [1, 2, 3].

Admitting additional fields and thus additional active atomic states enriches the dynamics of the process and gives new possibilities of its control. In particular by adding another upper state one obtains a double $\Lambda$ system [4] in which it is possible to make the medium transparent simultaneously for two signal pulses or to split a signal pulse into two ones of different frequencies [5, 6]. Channelization of information is also possible in the case of inverted $Y$ systems [7]. Coupling the $\Lambda$ system with a side level makes it possible to release a pulse of a frequency different from the original one or to temporarily prevent the signal from being released [8]. Recently the dynamics of light propagation in 5-level atomic systems in the $M$ configuration has also been discussed [9].

A special example is the so-called tripod atomic system in which three low-lying levels are coupled with an upper level by a signal field and two control ones. By a manipulation on control fields one can steer the propagation, storing and release of the signal, which may be stored in the form of two atomic coherences. Light propagation and slowdown in such a medium have been studied in Refs [10, 11] while the problem of light storing has been discussed in detail in our recent paper [12]. An attempt of releasing the trapped pulse by an arbitrary combination of two control fields results in general in splitting it into two parts one of which was leaving the sample while the other one remained trapped and could be released by switching on a new set of control fields. The evolution of such a system could be effectively described in terms of a couple of polaritons.

The previously used classical description of dividing the pulse into two or more parts may be essentially enriched by treating the pulse quantum-mechanically and by asking new questions about the photon statistics, especially important in the case of nonclassical light states, e.g., Fock states with a small photon number. Such an approach is also necessary in the consideration of quantum information storage in an atomic memory. Wang et al. [13] demonstrated a possibility of time splitting of a single photon by first storing it in one coherence, than by a sequence of transferring a part of the excitation into the other coherence and of a photon release from the latter. An additional transfer procedure based on a fractional adiabatic population transfer (F-STIRAP) had to be performed before each release procedure.

In this paper we investigate light propagation and storage of a quantum signal field in a medium of atoms in the tripod configuration. We discuss mapping of a quantum wave packet into two atomic coherences which means in particular a channelization of a single photon. In our approach, being more general than that of Ref. [13], at the release stage the two control fields may be twice switched on and off simultaneously. If the control pulses are complementary, which means that their amplitude and phase relations are properly chosen, the whole of the stored signal is divided into two parts, which are released at two separated time instants. In the case of a single trapped photon this leads to its time-entangled state, but without a need of an additional transfer stage of Ref. [13]. We discuss in detail the case in which two time-separated photons are trapped at the same place of the sample due to an action of two complementary sets of control pulses. They are later released by a different pair of sets of such pulses. We examine the possibility of photon coalescence, i.e. a release of both photons at the same release stage. By changing localizations and shapes of the stored photon wavepackets and the details of the control fields we obtain an analogue of the Hong-Ou-Mandel interferometer [14, 15, 16, 17] working on stored light. In contradistinction to the original version of such an interferometer we have to do with photons propagating always in the same spatial direction but well resolved in time.

II. QUANTUM POLARITONS IN THE TRIPOD CASE

In order to investigate quantum effects in light storing and propagation both the signal and the medium have to be described quantum-mechanically. Some formulæ of this section, pertaining to quantum operators, resemble those presented in our previous paper [12] for classical pulse. However, they are rederived below not only to make the presentation complete but also because they will serve to a discussion of physical effects connected with a quantum field, in particular with single photon states. In particular polaritons which in the present formulation are quantum-mechanical operators will play a central role in describing photon correlation effects.
Consider a one-dimensional medium of atoms in the tripod configuration (Fig. 1) including an upper state $a$ and three stable lower states $b$ (initial), $c$ and $d$. The sample is irradiated with three collinear laser beams. The quantized signal field couples the states $b$ and $a$ and is written as:

$$\epsilon(z,t) = \epsilon^+(z,t) + \epsilon^-(z,t) = \Sigma_{k} g_k a_k \exp[i(kz - \omega t)] \exp[-i(k_{ab}z - \omega_{ab}t)] + h.c.,$$

where $a_k$ is the annihilation operator of a photon with wave number $k$, $g_k = \sqrt{\frac{\hbar\epsilon_0}{2\epsilon N}}$. $V$ is the quantization volume, $\epsilon_0$ - the vacuum electric permittivity and $k_{a\beta} = \omega_{a\beta}/c = (E_a - E_\beta)/(\hbar c)$, $\alpha, \beta = a, b, c, d$. The two control fields $\epsilon_{23}(t) \cos(k_{23}z - \omega_{23}t)$ are treated classically and are supposed to be strong enough for the propagation effects to be neglected; they couple respectively the states $c$ and $d$ with $a$. The medium excitation is described in the formalism of the second quantization by the flip operators $\sigma_{a\beta} = |a><\beta| \exp(ik_{a\beta}z)$, $k_{a\beta}$ being the corresponding wave vector. The medium is treated in a continuous way; this is why $\frac{N}{L} \int dz$ in the above equation has replaced the sum over $N$ atoms ($L$ is the length of the sample).

The interaction hamiltonian in the rotating-wave approximation reads

$$H = \frac{N}{L} \int dz (\epsilon^+(z,t) - \epsilon^-(z,t)) \Sigma_{k} g_k \exp[i(kz - \omega t)] + h.c. - \frac{1}{2} d_{ac} \sigma_{ac} \exp[-i(k_{ac}z - \omega_{ac}t)]\epsilon_2(t) \exp[i(k_2z - \omega_2t)] + h.c. - \frac{1}{2} d_{ad} \sigma_{ad} \exp[-i(k_{ad}z - \omega_{ad}t)]\epsilon_3(t) \exp[i(k_3z - \omega_3t)] + h.c.,$$

where $d_{a\beta}$ are the dipole moment’s matrix elements, assumed for simplicity to be real and positive.

In our analysis in terms of polaritons it is enough to restrict oneself to the deterministic part of the Heisenberg evolution (relaxation- and noise-free). The equations of motion for the flip and signal field operators in the first-order approximation with respect to the signal field, in resonance conditions read

$$i\dot{\sigma}_{ba}(z,t) = -\frac{d_{ba}}{\hbar} \epsilon^+(z,t) - \Omega_2(t) \sigma_{bc}(z,t) - \Omega_3(t) \sigma_{bd}(z,t),$$

$$i\dot{\sigma}_{bc}(z,t) = -\frac{d_{bc}}{\hbar} \epsilon^+(z,t),$$

$$i\dot{\sigma}_{bd}(z,t) = -\frac{d_{bd}}{\hbar} \epsilon^+(z,t),$$

$$\left(\frac{\partial}{\partial z} + \frac{\kappa^2}{4}\right)\epsilon^+(z,t) = \frac{\hbar \epsilon}{\kappa} g^2 d_{ba} \sigma_{ba}(z,t),$$

where $g$ is the value of $g_k$ for the central field frequency of the signal field, and $\Omega_2(t) = d_{ac} \epsilon_2(t)/(2\hbar)$ and $\Omega_3(t) = d_{ad} \epsilon_3(t)/(2\hbar)$ are the Rabi frequencies for the control fields.

It is convenient to express the solutions of Eqs (3) in terms of two field+medium excitations called polaritons: $\Psi(z,t)$ (the dark-state polariton) and $Z(z,t)$

$$\Psi(z,t) = \exp(-i\chi_2)\{\epsilon^+(z,t) \cos \theta - \frac{d_{bc}}{\hbar\epsilon} \sin \theta \exp(i\chi_3) \sin \phi \sigma_{bc}(z,t) + \exp(i\chi_3) \sin \phi \sigma_{bd}(z,t)\},$$

$$Z(z,t) = \frac{\hbar}{d_{bc}} \exp(i\chi_2) \sin \phi \sigma_{bc}(z,t) - \exp(i\chi_3) \cos \phi \sigma_{bd}(z,t),$$

where the mixing angles $\theta$ and $\phi$ are defined as $\theta = \kappa/\Omega$, $\tan \phi = |\Omega_3/\Omega_2|$ with $\chi_2 = \arg(\Omega_2)$, $\chi_3 = \arg(\Omega_3)$, $\Omega = \sqrt{|\Omega_2|^2 + |\Omega_3|^2}$, $\kappa^2 = N \epsilon^2 |d_{ab}|^2 / \hbar^2$. $\chi$ satisfies the equation $\dot{\chi}_2 = \sin^2 \theta \dot{\chi}_2$. For simplicity we assume that $\dot{\chi}_2 = \dot{\chi}_3$.

The field and medium components of the polariton $\Psi$ are

$$\epsilon^+ = \exp(i\chi_2) \Psi \cos \theta,$$

$$\exp(i\chi_2) \cos \phi \sigma_{bc} - \exp(i\chi_3) \sin \phi \sigma_{bd} = \frac{d_{bc}}{\hbar\epsilon} \exp(i\chi_2) \Psi \sin \theta.$$

The evolution equations for the two polaritons read

$$\left(\frac{\partial}{\partial z} + \frac{\kappa^2}{4}\right)\Psi = \exp(-i\chi_2) \sin \theta \dot{\phi} Z,$$

$$\frac{\partial}{\partial z} Z = i\dot{\chi}_2 Z - \exp(i\chi_3) \sin \theta \dot{\phi} \Psi.$$

The operators $\Psi(z,t)/(g\sqrt{L})$ and $Z(z,t)/(g\sqrt{L})$ fulfill typical bosonic commutation relations in the first-order approximation, in which $\sigma_{bb} = 1$, $\sigma_{ac} = \sigma_{dd} = \sigma_{ad} = 0$, ([$\Psi(z)$, $\Psi^\dagger(z')$]) = [Z(z), Z^\dagger(z')] = g^2 L \delta(z-z'), [\Psi(z), Z^\dagger(z')] = 0, so those operators may be considered annihilation operators of the joint field+medium excitation. The essence of the approach used in this paper will be treating a medium excitation due to the stored photon as a superposition of two excitations, properly suited to the pulse releasing control fields (see Eq. (9) below).
If the control fields are proportional, i.e. $\phi = 0$ and $\chi_2 = \chi_3 = 0$ the evolution equations for the polaritons are decoupled. The polariton $\Psi$ propagates through the medium without changing its “shape”, i.e. the solution is

$$\Psi(z, t) = \Psi(z - c \int_0^t \cos^2 \theta(t') dt', t = 0),$$

while the polariton $Z = \text{const}$ keeps both its shape and localization.

The above equations constitute a generalization of the approach of Fleischhauer et al. [2] for a four-level system and at the same time of our previous results [12], formulated in the language of a quantum signal field. Due to such an approach we will be able to discuss the effect of the process in terms of photons rather than in the language of splitting a classical pulse.

If at the beginning the control fields corresponding to a given constant mixing angle $\phi^0$ and given constant $\chi_2^0$ and $\chi_3^0$ are strong enough ($\theta \approx 0$), the medium is transparent and an incoming photon enters the medium. Switching the control field off ($\theta = \frac{\pi}{2}$) results in the medium becoming opaque and the photon becomes trapped in a coherent superposition of the atomic coherences. The propagation and storing are described in terms of the polaritons $\Psi^0$ and $Z^0$. The former polariton has evolved from being almost purely electromagnetic to become purely atomic while the latter one has not taken part in the evolution. The probability amplitudes that the photon has actually been trapped released at the two stages, which would correspond to changing the transmission and reflection rate of a usual beam splitter.

Also the released photons have the same direction but are separated in time. By changing the mixing angle $\phi$ and at the same time of our previous results [12], here formulated in the language of a quantum signal field. Due to the above results are much more general than those of Ref. [13] in which the authors proposed a time splitting of a photon by first storing the pulse in one coherence $\sigma_{bc}$ (here this corresponds to $\phi^0 = 0$), than by pumping a part of the excitation into the other coherence $\sigma_{bd}$ and finally by a pulse release from the latter coherence (here $\phi^1 = \frac{\pi}{2}$). The pumping procedure (F-STIRAP) was to be separated from the release stages which required using light pulses transverse to both signal and control fields; the latter should be realized in ultracold gases. Here, due to using simultaneously pairs of control fields of the same shape pumping and releasing stages coincide.

### III. BEAM SPLITTER AND HONG-OU-MANDEL - TYPE INTERFEROMETER

The above-described behavior of the atom+field allows one to use the medium and the system of control fields as an effective and flexible beam splitter working on stored light. The technical difference between the usual device and that of ours is that incoming photons may arrive from only one direction but at one or two different time instants. Also the released photons have the same direction but are separated in time. By changing the mixing angle $\phi$ and the phases $\chi_2^1$ and $\chi_3^1$ of the control fields one can smoothly regulate the amplitudes and phases of the pulse components released at the two stages, which would correspond to changing the transmission and reflection rate of a usual beam splitter.

The operation of photon storing can be performed at two stages (input ports). First we trap a first portion of incoming photons by switching on and off a pair of control fields characterized by the parameters $\phi^0, \chi_2^0$ and $\chi_3^0$. After the this part of the trapping operation has been finished we may trap a second portion of photons by applying the control fields of the mixing angle $\frac{\pi}{2} - \phi^2, \chi_2^2 + \pi$ and $\chi_3^2$. Note that the second trapping operation does not affect the coherences due to the first one. Thus a photon from the first portion is described by the polariton field $\Psi^0$ and...
that from the second portion - by \( Z^0 \). A special case is assumed in which \( \phi^0 = 0 \) which means that photons have been first trapped in the coherence \( \sigma_{bc} \) and next - in \( \sigma_{bd} \).

The release operation consists also of two stages (output ports) connected with two pairs of control fields separated in time, the first of which is characterized by the mixing angle \( \phi^1 \) and \( \chi_2^1 \) and the second by \( \chi_2^1 + \pi \) and \( \chi_3^1 \). The output includes thus two portions of photons, separated in time, and the photon release is complete.

In the case of trapping and releasing of two photons we have a new realization of the Hong-Ou-Mandel effect. In its original formulation it concerns a 50-50 photon beam splitter with exactly one incoming photon in each of the two input ports (incoming photon directions). Due to interference phenomena the probability amplitudes for obtaining exactly one photon in each of the two output ports (outgoing photon directions) cancel out.

In our realization of the input ports single photons, separated in time, are trapped, one in each of the two superpositions of the atomic coherences, connected with complementary sets of control fields. In the output ports photons are released, again at two stages, from two other superpositions of the two coherences. The analogue of photon incoming from two orthogonal directions is their storing at two storing stages. The analogues of reflection and transmission on a traditional beam splitter are photon release at the first or second release stage while the parameters characterizing the control fields and the atomic system determine the analogues of the reflection and transmission coefficient and phase relations. Using the above-mentioned transformations of polaritons at the storing stage and having in mind that the incoming polaritons \( \Psi^0, Z^0 \) (outgoing polaritons \( \Psi^1, Z^1 \)) become almost purely electric field operators \( e^1, e^2, (e^3, e^4) \) for \( t \to -\infty \) \( (t \to \infty) \) we may give the net result for the transformation of the field operators

\[
\left( \begin{array}{ll} R_{31} & R_{32} \\ R_{41} & R_{42} \end{array} \right) \left( \begin{array}{l} e^1 \\ e^2 \end{array} \right) \to \left( \begin{array}{l} e^3 \\ e^4 \end{array} \right).
\]

where

\[
\begin{align*}
R_{31} &= \cos \phi^1 \cos \phi^0 \exp[i(\chi_2^1 - \chi_2^0)] + \sin \phi^1 \sin \phi^0 \exp[i(\chi_3^1 - \chi_3^0)], \\
R_{32} &= \cos \phi^1 \sin \phi^0 \exp[i(\chi_2^1 - \chi_2^0)] - \sin \phi^1 \cos \phi^0 \exp[i(\chi_3^1 - \chi_3^0)], \\
R_{41} &= \sin \phi^1 \cos \phi^0 \exp[i(\chi_2^1 - \chi_2^0)] - \cos \phi^1 \sin \phi^0 \exp[i(\chi_3^1 - \chi_3^0)], \\
R_{42} &= \sin \phi^1 \sin \phi^0 \exp[i(\chi_2^1 - \chi_2^0)] + \cos \phi^1 \cos \phi^0 \exp[i(\chi_3^1 - \chi_3^0)].
\end{align*}
\]

Let the respective wave packets of the stored excitation be \( f_1(z) \) and \( f_2(z) \), their shape being identical with the shape of the initial photon wavepackets and their localization depending on the time instants of the switch-off of the control fields. The maximum value (unity) of the overlap of the packets \( s \equiv \int f_1^*(z)f_2(z)dz \) corresponds to the situation in which two photons with wavepackets of the same shape have been stored exactly at the same place inside the sample.

The field operators corresponding to the input ports are \( \Psi^0(j) = \int f_1^*(z)\Psi^0(z)dz, \ Z^0(j) = \int f_2^*(z)Z^0(z)dz \). The corresponding operators in the output ports are \( \Psi^1(j) = \int f_1^*(z)\Psi^1(z)dz, \ Z^1(j) = \int f_2^*(z)Z^1(z)dz \).

The key commutation relations are \( [\Psi^0(1),\Psi^0(2)] = [Z^0(1),Z^0(2)] = g^2L, \ [\Psi^0(1),\Psi^1(2)] = [Z^0(1),Z^1(2)] = g^2Ls \), with analogous relations for \( \Psi^1 \) and \( Z^1 \).

The quantum state of the medium after trapping has been accomplished is constructed as due to a creation of two excitations due to complementary control fields and characterized by two possibly different wave packets

\[
|\zeta> = \frac{1}{g^2L}\Psi^0(1)Z^0(2)|0>.
\]

where the "vacuum" state \( |0> \) means all atoms in the state \( |b> \). The state corresponding to two photons being released at the first stage (photon coalescence) is constructed as due to a creation of two excitations corresponding to the first release stage characterized by possibly different wave packets

\[
|\zeta_1> = \frac{1}{g^2L}\sqrt{\frac{1}{1 + |s|^2}}\Psi^0(1)\Psi^1(2)|0>.
\]

where a normalizing factor for the two-particle states has been included. Note that the state corresponding to the photon coalescence at the second stage is constructed in an analogous way, using the \( Z \) operators instead of \( \Psi \).

Calculating the projections with the use of the above mentioned commutation relations yields the probability amplitude of photon coalescence at the first release stage

\[
<\zeta_1|\zeta> = (1 + |s|^2)\{\cos \phi^0 \cos \phi^1 \exp[-i(\chi_2^0 - \chi_2^1)] + \sin \phi^0 \sin \phi^1 \exp[-i(\chi_3^0 - \chi_3^1)]\} \\
\{\sin \phi^0 \cos \phi^1 \exp[-i(\chi_2^0 - \chi_2^1)] - \cos \phi^0 \sin \phi^1 \exp[-i(\chi_3^0 - \chi_3^1)]\}
\]

(14)
If in particular $\phi^0 = 0$, which means that exactly one photon has been trapped in each coherence, then the photon coalescence probability at the first release stage is

$$P_{\text{coal}}(1) = \frac{1}{4} (1 + |s|^2) \sin^2 2\phi^1. \quad (15)$$

If the wave packets overlap ($s = 1$) we obtain for $\phi^1 = \frac{\pi}{4}$ (equal amplitudes of the releasing control fields) that $P_{\text{coal}}(1) = \frac{1}{2}$. The same value is obtained for the coalescence probability at the second stage. This means that the situation is impossible in which exactly one photon is released at each stage. This is analogue of a symmetric beam splitter with exactly one photon in each input port. A reduction of the overlap integral leads in a continuous way to the situation in which the probability of the latter situation grows to $\frac{1}{2}$; this is an analogue of changing the length of the arms of the standard HOM interferometer. For example if the wave packets $f_{1,2}$ are Gaussians with standard deviations $\delta_{1,2}$ separated by a distance $a$, the probability of releasing a single photon at each stage (photon noncoalescence), being our analogue of coincident registration of single photons in the usual realization, is

$$P_{\text{noncoal}} = \frac{1}{2} [1 - \frac{2}{\delta_1^2 + \delta_2^2} \exp(-\frac{a^2}{2(\delta_1^2 + \delta_2^2)})]. \quad (16)$$

This probability as a function of the distance $a$ exhibits a minimum called the Mandel dip, reaching zero for equal packet widths (see Fig. 2), which corresponds to an ideal photon coalescence. On the other hand in the case of coinciding packet centers ($a = 0$) the probability equals $\left(\delta_1 - \delta_2\right)^2 / [2(\delta_1^2 + \delta_2^2)]$ (see Fig. 3). Note that it is only for equal packets’ widths that the coalescence probability drops to zero.

The effect can also be controlled by choosing the phases of the control field. From Eq. (14) it follows for example that for $\phi^0 = \phi^1 = \frac{\pi}{4}$ the coalescence probability at the first stage is

$$P_{\text{coal}}(1) = \frac{1}{4} (1 + |s|^2) \sin^2 \left(\chi_0^2 - \chi_2^0 - \chi_1^0 + \chi_3^1\right). \quad (17)$$

Thus the probability of the photons’ coalescence strongly depends not only on the shape of the wave packets and the relative amplitudes of the control fields but also, in an oscillatory way, on the relative phases of the latter.

IV. CONCLUSIONS

We have analyzed the propagation and storing of photons in an atomic medium in the tripod configuration. An adiabatic evolution of the system can be described in terms of a couple of polaritons. Using two complementary sets of control fields at the storing stage one can stop two photons arriving at different time instants. Applying two other complementary sets of the control fields at the release stage one retrieves two time-entangled photons. The field operators of the outgoing photons are expressed by those of the incoming ones, depending on the parameters of the control fields. We have thus obtained a kind of a beam splitter, operating on a stored light, in which the input and output ports correspond to photons arriving and leaving at different time instants. Our analogues of the transmission and reflection coefficients as well phase relations can be regulated on demand. We have demonstrated the effect of the Hong-Ou-Mandel-type interference and analyzed its dependence on the shapes and positions of the trapped photon wavepackets and on the parameters of the control fields.

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Figures

FIG. 1: Level and coupling scheme of the tripod system.

FIG. 2: The probability of photon noncoalescence as a function of the packet’s separation $a$ for $\delta_1 = \delta_2$ solid line, and $\delta_1 = 3\delta_2$, dashed line.

FIG. 3: The probability of photon noncoalescence as a function of the ratio $\delta_2/\delta_1$ of the packets’ widths for $a = 0$, solid line, and $a = 2\delta_1$, dashed line.
$a[\delta_1]$

probability
