Propagation and storing of light in optically modified atomic media

Jarosław Zaremba
Institute of Physics
Nicolaus Copernicus University
ul. Grudziądzka 5/7
87 100 Toruń, Poland
E-mail: zaremba@fizyka.umk.pl

Abstract. Coherent interactions of laser light with atomic ensembles allow one to modify dispersive properties of a medium and lead to new optical phenomena. Studies of the controlled light propagation and storing in such media have recently become a dynamically developing field of research motivated both by the fundamental character of the processes and by potential applications. This article briefly reviews basic theoretical approach to the dynamics of the propagation of laser pulses in optically modified media. The method and the physical processes are discussed that allow one to slow down the group velocity of laser pulse to zero (stopping of light), to transfer the state of a light pulse to atomic coherences and to restore the pulse. The interpretation of these phenomena in the formalism of dark-state polaritons is presented. Examples of possible coherent manipulations on a stored light are also discussed.

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1. Introduction
Propagation of light pulses in atomic media is a fundamental optical phenomenon the properties of which are determined by the details of light-atoms interaction. A typical effect accompanying the propagation is the attenuation of a light beam due to the absorption of light by atoms when the radiation frequency matches the frequency of a particular atomic transition. The optical properties of a dielectric medium can be characterized by the frequency dependent electric susceptibility the imaginary part of which describes the absorptive while its real part - the dispersive properties of the medium which, via the refractive index, govern the group velocity of the pulse. Typically, in a resonance region the absorption spectrum curve presents a strong maximum while the dispersion curve is a decreasing function of frequency (steep anomalous dispersion).

In the past few decades modern optical techniques based on coherent interaction of atoms with laser radiation leading to quantum interference [1] made it possible to influence and modify this typical behaviour of a light pulse propagation by steering in a controlled way dispersive properties of an atomic sample. This is achieved by interactions of atoms with additional laser fields which ”dress” a medium. As a result its dispersive properties may by strongly modified and the optical response of the medium may be significantly changed leading to new atomic and optical coherence phenomena (for a review of various phenomena of such a kind see, e.g., [2]).
In particular, it becomes possible to make an initially opaque medium transparent for resonant radiation and to control and manipulate the propagation of laser pulses modifying their group velocities and even to "stop" and then retrieve them. The phenomenon which forms the basis for such possibilities has been termed electromagnetically induced transparency (EIT) in 1990 [3] and since its observation in 1991 [4] has become a point of departure for a discovery of many interesting physical effects and potential applications ranging from nonlinear optics to quantum information [5], [6], [7].

The intention of the present lecture is to give a concise introductory presentation of basic physical ideas underlying EIT-based modifications of the group velocity of a laser pulse, of "slow" and "stopped" light and of a theoretical analysis of these processes. We refer readers for much more detailed discussion to the comprehensive reviews [8], [9], [10], [11]. It has to be stressed that much work has been published in this field and the literature of the subject is still growing.

The paper is organized as follows. In section 2 physical grounds of the EIT are presented and modification of the group velocity is discussed in a lambda system. In section 3 the idea of light storage is discussed and the picture of dark-state polaritons is introduced. More complicated situation of four level atomic system and possibilities of coherent manipulations on the stored light are presented in section 4.

2. EIT and dispersive properties of a lambda system

A generic system in which the EIT can be discussed is a three-level atomic system coupled to two co-propagating laser fields in the lambda configuration. In such a configuration two lower long living states \( b \) and \( c \) (e.g., Zeeman sublevels of a hyperfine structure in the atomic ground state) are coherently coupled to a common excited state \( a \): a weak probe or signal field 1 of the amplitude \( \epsilon_1 \) and the frequency \( \omega_1 \) is tuned to the transition \( b \rightarrow a \) while a strong control laser field 2 of the amplitude \( \epsilon_2 \) and the frequency \( \omega_2 \) drives the transition \( c \rightarrow a \). The coupling \( b \rightarrow c \) is the dipole forbidden transition. The lambda configuration is shown in Figure 1. The combined effect of the interaction with two fields is the so-called dark superposition of the lower states. The atoms in such a dark state are decoupled from the upper state, the light is not absorbed and the population is trapped in lower states [12]. Thus, if the control laser field were absent the probe beam would be resonantly absorbed, switching the control field on makes the system transparent for the probe radiation: this is the electromagnetically induced transparency effect.

![Lambda and Tripod Configurations](image)

**Figure 1.** Lambda (left) and tripod (right) configurations. 1 denotes signal field, 2 and 3 denote control fields and black circles the initial population.

The optical response of the medium to the probe laser field is characterized by its complex susceptibility \( \chi = \chi' + i\chi'' \). To determine the complex susceptibility we analyze the density matrix of the three-level atom interacting with both laser fields which fulfills the von Neumann
equation completed with phenomenological relaxation terms

\[ i\hbar \dot{\rho} = [H, \rho] + R\rho \]  

(1)

where \( \rho = \rho (z,t) \) denotes the density matrix of the atom in the position \( z \), \( R \) describes the relaxation processes, and \( H \) is the Hamiltonian which in the rotating wave approximation reads

\[ H = h\omega_a |a⟩⟨a| + h\omega_b |b⟩⟨b| + h\omega_c |c⟩⟨c| + \left( -\hbar\Omega_1 e^{i(\omega_1 t - k_1 z)} |b⟩⟨a| - \hbar\Omega_2 e^{i(\omega_2 t - k_2 z)} |c⟩⟨a| + h.c. \right) \]  

(2)

where \( \Omega_1 = \frac{1}{\pi} d_{ba} \epsilon_1 \) and \( \Omega_2 = \frac{1}{\pi} d_{ca} \epsilon_2 \) denote the Rabi frequencies of the corresponding couplings and \( d_{ba} \) and \( d_{ca} \) are the matrix elements of the electric dipole operator. We assume that Rabi frequencies are real.

Making the substitution \( \rho_{ab} = \sigma_{ab} e^{-i\phi_1} \), \( \rho_{ac} = \sigma_{ac} e^{-i\phi_2} \), \( \rho_{cb} = \sigma_{cb} e^{-i(\phi_1 - \phi_2)} \), \( \sigma_{ii} = \sigma_i \) where \( \phi_i = \omega_i t - k_i z \) one obtains the system of optical Bloch equations which describe the evolution of the slowly varying matrix elements of the density matrix \( \sigma \).

Making the perturbative approximation with respect to the weak signal field 1 but treating the strong control field 2 exactly, assuming that the control field is in resonance with the corresponding atomic frequency (\( \omega_2 = \omega_{ac} \)) and that the atom is initially in the state \( b \), one obtains the low intensity limit of the equations of motion in the form:

\[ \sigma_{ab} = - (\gamma_{ab} - i\Delta_1) \sigma_{ab} + i\Omega_1 \sigma_{cb} + i\Omega_2 \sigma_{ac} \]  

(3)

\[ \sigma_{cb} = - (\gamma_{cb} - i\Delta_1) \sigma_{cb} + i\Omega_2 \sigma_{ab} \]  

(4)

where \( \Delta_1 = \omega_1 - \omega_{ab} \) is the detuning of the signal field from atomic resonance and \( \gamma_{ik} \) are the constants characterizing the relaxations of the corresponding atomic coherences. The stationary solution of this system of equations allows one to find the susceptibility \( \chi \) of the atomic system. To do this one calculates the polarization of the atomic system (\( N \) is the atomic density of the sample)

\[ P = N \langle d \rangle = Tr (\rho d) \]  

(5)

which, on the other hand, is related to the signal field \( \epsilon_1 \) by the relation \( P = \epsilon_0 \chi \epsilon_1 \) (\( \epsilon_0 \) denotes the permittivity of the vacuum). One obtains:

\[ \chi = \frac{2Nd_{ba}}{\epsilon_0 \epsilon_1} \sigma_{ab} \]  

(6)

The plots of the real and imaginary parts of the complex susceptibility as a function of a normalized detuning of the signal beam are shown in Figure 2. At resonance, the imaginary part is equal to zero which corresponds to cancellation of absorption of the signal beam, for which strong absorption would normally be expected. The medium has thus become, due to an interaction with a strong coherent control field, transparent for a resonant radiation. This is the electromagnetically induced transparency effect and the region of a weak absorption in the vicinity of resonance represents the transparency window for incoming radiation. Also, the dispersive properties of a medium illustrated by real part of susceptibility are greatly modified: a steep normal dispersion is observed inside the transparency window. Such a behaviour is completely different from that occurring in dielectric media with typical dispersive properties: the optical response of the system is significantly and qualitatively modified by coherent excitation in the lambda configuration.

Dispersive properties of a medium are characterized by the frequency dependent refractive index \( n = n(\omega) \). The group velocity of a light pulse propagating in such a medium is defined as:

\[ v_g = \frac{d\omega}{dk} \]  

(7)
The dispersion relation \( kc = \omega n(\omega) \) leads to the following formula for the group velocity:

\[
v_g = \frac{c}{n(\omega) + \omega \frac{dn(\omega)}{d\omega}}
\]  

(8)

In the region of a steep normal dispersion the derivative \( \frac{dn(\omega)}{d\omega} \) can attain great values which will result in an important decrease of \( v_g \) of the pulse. Thus the pulse propagating inside the sample will be delayed with respect to the pulse propagating at the same distance in vacuum - this is the manifestation of a "light slow down" effect. When the light pulse enters the medium it exhibits a spatial compression: the front edge of the pulse is already decelerated while its back edge, still outside the medium, propagates in vacuum with the speed \( c \) much greater than the group velocity inside the medium. This picture of propagation requires that the spectral width of the pulse should be smaller than the bandwidth of the transparency window which increases proportionally to the square of the electric field intensity of the control field. However, the larger is the transparency window, the smaller the value of the derivative of the refractive index inside it.

3. Stored light and dark-state polaritons

The transparency of the medium in EIT can be modified in time via modifications of the intensity of the control field. Such a control of this phenomenon opens new possibilities of steering the propagation of light pulses. Switching off the control field makes the medium opaque and the signal pulse that has entered it cannot propagate: it is stored in the atoms of the medium. However, it should be stressed that it is not absorbed by atoms - the process is reversible and when the control field is switched on again and the transparency conditions are reestablished, the pulse reappears and is retrieved. This phenomenon is sometimes termed "stopping of light" which is an illustrative although not correct name: the pulse is rather temporarily stored in the form of an atomic coherence and can be released after some time. This effect has been demonstrated in an ultracold sample of sodium atoms [13], in a gaseous sample of warm rubidium atoms [14] and in a solid state [15].

To discuss the dynamics of propagation in EIT medium we consider a quasi one-dimensional sample of three-level atoms. The position of an atom is described by the continuous variable \( z \) and thus the atomic density matrix is \( z \)-dependent. The above Bloch equations should be completed with the Maxwell propagation equation for the signal field \( \epsilon_1 \) or for the corresponding
Rabi frequency $\Omega_1$, which reads in the slowly varying envelope approximation (SVEA, see, e.g., [16])

$$\frac{\partial \Omega_1}{\partial z} + \frac{1}{c} \frac{\partial \Omega_1}{\partial t} = -i\kappa^2 \sigma_{ba}$$

(9)

where $\kappa^2 = \frac{N|d_{ab}|^2 \omega_1}{2\hbar \epsilon_0}$. The r.h.s term of Eq (9) is a source term which describes the polarization of the medium and must be determined from the Bloch equations. One thus obtains the coupled Bloch-Maxwell equations which are the fundamental tool to analyze the propagation phenomena in dielectric media and the solution (usually numerical) of these equations allows one to obtain a detailed information on the process.

A good insight into the dynamics of light propagation may be obtained from the formulation of its description in terms of a so-called dark-state polariton being a combination of signal field and atomic coherence. Such an approach to problems of stopped light has been introduced by Fleischhauer and Lukin [17]. Note that in their original paper the polaritons were introduced as quasi-particles in a fully quantum approach, nevertheless in the present semiclassical description the corresponding objects will also be called polaritons.

To discuss the light storage we will consider the propagation of a weak signal pulse of envelope $\epsilon_1(z, t)$ in the lambda system described in the previous chapter. For simplicity, the propagation effects for the strong control field are neglected, i.e., its envelope $\epsilon_2 = \epsilon_2(t)$ is $z$-independent and one assumes that both fields are in resonance with atomic transitions. If one additionally makes perturbative approximation with respect to the signal field (the control field is treated exactly), neglects the relaxations and makes the adiabatic approximation $\dot{\sigma}_{ab} = 0$ then from Eqs (3), (4) and (9) there follows the equation satisfied by the dark-state polariton:

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

(10)

The dark-state (or simply dark) polariton is defined as

$$\Psi(z, t) = \cos \theta(t) \Omega_1 - \kappa \sin \theta(t) \sigma_{bc}$$

(11)

where the mixing angle $\theta$ is defined as:

$$\tan \theta(t) = \frac{\kappa}{\Omega_2}$$

(12)

It is a shape-preserving solution of the wave equation (10)

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

(13)

and propagates in the sample with the group velocity

$$v_g = c \cos^2 \theta = \frac{1}{1 + \tan^2 \theta}$$

(14)

This velocity depends, via the mixing angle $\theta$, on the intensity of the control field and can thus be controlled by varying this field in time.

The dark polariton is built of two physically different components: the signal field (term with $\Omega_1$) and the atomic coherence (term with $\sigma_{bc}$). The relative weights of these components depend on the control field intensity and change adiabatically. Thus, although the polariton does not change its shape during the propagation, these components change considerably. When the control field is switched on ($\theta \approx 0$) the medium is transparent, the polariton is "purely optical": $\Psi(z, t) = \Omega_1(z, t)$, and the signal pulse propagates inside the sample. Switching the
control field off \((\Omega_2(t) \to 0, \theta(t) \approx \frac{\pi}{2}, v_g \to 0)\) makes the medium opaque, the polariton decelerates and becomes "purely atomic": \(\Psi(z,t) \to -\kappa \sigma_{bc}(z,t)\) and the pulse is "stopped". At the release stage at which the control field is switched on again, the atomic component of the polariton decreases while its field component increases; the polariton becomes "optical" again which means that the signal pulse is reconstructed and retrieved. Such an evolution of the form-stable polariton explains the mechanism of the stopping of light. The signal field is mapped into the atomic coherence which is created during the switching the control field off. The pulse is thus stored inside the medium. The process of light storing is reversible and the information about the signal is read out from atomic coherence at the release stage. In this ideal scenario the coherence has not been disturbed by any relaxation, the control field at the release stage is identical with that at the storage stage and the released signal is a replica of the stored one. If the adiabatic approximation were not valid, the description of the evolution only in terms of the dark-state polariton would not be sufficient and another superposition of the field and atomic coherence called a bright-state polariton would be necessary to analyze the dynamics [18].

To describe the dynamics of the pulse storing we have used a semiclassical analysis in which the signal light has been treated as an classical electromagnetic wave. In fact, the original polariton interpretation of the process by Fleischhauer and Lukin used a second quantization formalism and a quantized light field [17]. The dark-state polariton should then be interpreted as a quasi particle corresponding to excitations of the composed atoms plus field system. Its evolution is essentially the same as that described above and the light storage and retrieval consists now in a reversible mapping of the quantum state of the signal into the atomic coherences. Such a picture forms a basis for potential interesting applications in quantum information processing: the information conveyed by photonic quantum state of the signal pulse can be stored (written down) in a ”quantum memory” [18] formed by collective atomic excitations described by an atomic coherence and later retrieved (read out) by switching the control field on again and releasing the ”stopped” signal pulse.

4. Coherent processing of stored light

Light storing in the form of atomic coherences opens new possibilities of controlled modifications of the properties of the released pulse by processing the atomic medium during the storage stage. Such an additional control of the process can be obtained in more advanced systems admitting additional fields and additional active atomic states. The dynamics is then essentially enriched and various systems have been studied in the literature. Here we give an example illustrating such possibilities, namely the case of a medium of four-level atoms interacting with three laser fields in the so-called tripod configuration. Three lower states \(b, c, d\) are coherently coupled with a common excited state \(a\): a signal field 1 drives the \(b \to a\) transition while two control fields 2 and 3 couple the state \(a\) with the states \(c\) and \(d\) respectively and there is no coupling among lower states. The tripod configuration is shown in Figure 1.

Light propagation and slowdown in such a system have been considered in [19], [20], [21] while its storing has recently been studied in [22] using a generalized polariton picture of the process. The medium which is opaque for the resonant signal beam may be made transparent by switching on the control fields. As in the case of a lambda system one can, by manipulations on control fields, modify dispersive properties of the medium and steer the propagation and storing of the signal pulse. However in this case two essentially new aspects appear. First, the signal is stored in two atomic coherences \(\sigma_{bc}\) and \(\sigma_{bd}\). It becomes thus possible to manipulate the stored light by applying additional interactions which modify these two coherences and, as a consequence, to modify the properties of the released pulse. Second, the signal pulse can be stored and retrieved by different sequences of control fields. In the following we will briefly discuss examples of both cases.
The detailed analysis of pulse propagation in an atomic medium in the tripod configuration requires solutions of the Bloch-Maxwell equations. The generalized polariton description of the dynamics is also possible and allows one to obtain a clear picture of a light storage and retrieval analogous to that of a lambda case [22]. However, a single dark-state polariton is not sufficient even in the case of an adiabatic evolution. The dark-state polariton is now of the form

\[ \Psi(z, t) = \cos \theta \Omega_1 - \kappa \sin \theta (\cos \phi \sigma_{bc} + \sin \phi \sigma_{bd}) \]  

where the mixing angle \( \theta \) and the angle \( \phi \) depend on both control fields intensities

\[ \sin^2 \theta = \frac{\kappa^2}{\kappa^2 + \Omega^2} \quad \cos^2 \theta = \frac{\Omega^2}{\kappa^2 + \Omega^2} \quad \tan \phi = \frac{\Omega_3}{\Omega_2} \quad \Omega^2 = \Omega_2^2 + \Omega_3^2 \]  

The dark-state polariton is here a natural generalization of that for a lambda system. As in that case it is built of two physically different components: the signal field term with \( \Omega_1 \) and the atomic coherence term which, however, contains now two atomic coherences \( \sigma_{bc} \) and \( \sigma_{bd} \). For strong control fields \( \theta \approx 0 \) and the polariton is "purely electromagnetic" while for vanishing control fields \( \theta = \pi/2 \) and it becomes "purely atomic"; thus modifying the intensities of control fields in time one changes the relative weights of these two components. To describe fully the dynamics a second polariton of the form

\[ Z(z, t) = \sin \phi \sigma_{bc} - \cos \phi \sigma_{bd} \]  

is necessary, it is "purely atomic" composed solely of two atomic coherences and does not contain a field component. The two polaritons satisfy the coupled equations of motion that follow, under the assumption identical to those made in the case of a lambda system, from Bloch-Maxwell equations

\[ \left( \frac{\partial}{\partial t} + c \cos^2 \theta \frac{\partial}{\partial z} \right) \Psi(z, t) = \tan^2 \theta \cos \theta \Omega Z \]  

\[ \frac{\partial Z}{\partial t} = -\cos \theta \frac{1}{\Omega} \Psi \]  

Thus, the adiabatic evolution becomes quite complicated because, in general, the polariton \( Z \) can be created inside the atomic medium even if initially, before the pulse has entered the sample, \( \Psi = \Omega_1 \) and \( Z = 0 \). It is possible that the polariton \( Z \) remains different from zero even after releasing the pulse when \( \Psi = 0 \) inside the sample. It means that a part of atomic coherences in which the initial pulse has been mapped can remain trapped in spite of the fact that control fields have been switched on again.

The equations of motion are decoupled and \( Z \) does not change its shape during evolution if \( \phi = \text{const} \), i.e. if both control fields are identical, \( \Omega_2 = \Omega_3 \), or change in a proportional way. If additionally the initial value of \( Z \) is zero, which is the natural initial condition, the evolution is completely determined by the form-stable dark-state polariton \( \Psi(z, t) \) only (as in the case of a lambda system) which propagates with a group velocity \( v(t) = c \cos^2 \theta(t) \) where \( \cos \theta \) is now given by (16) and the coherences into which the pulse has been mapped are then proportional.

To illustrate the above discussion we will briefly demonstrate two examples; the detailed discussion of them has been given in [22]. In the first case the signal pulse is stored and retrieved by pairs of identical simultaneously switched-off and -on control pulses but the atomic coherences are modified in the storage stage. In the second case the signal pulse is stopped by two identical control fields but in the stage of releasing it the control fields are shifted in time.

Assume that identical control fields have been switched-off simultaneously \( (\dot{\phi} = 0) \) and a signal pulse has been stopped so that \( Z = 0 \) and two proportional coherences have been created. Then an additional interaction is switched on which modifies the coherences before the pulse is
released. If the states $b, c, d$ are Zeeman substates (e.g., of the hyperfine structure as it is in the case of rubidium atoms typically used in experiments on light storage) such an interaction may be introduced by a pulse of a magnetic field parallel to the direction of propagation. A magnetic field pulse introduces additional phase shift $\delta$ between atomic coherences which leads to a creation of $Z \neq 0$ polariton the value of which depends on this phase shift, i.e. on the pulse duration. If the control fields used to release the signal pulse are identical to those used to store it the polariton $Z$ does not change which means that the part of coherences constituting $Z$ does not leave the sample. As a consequence, the coherences cannot be turned back into the pulse completely and the released signal is not a replica of the stored one. Thus a controlled modification of the relative phase of coherences changes the resulting output signal which depends on $\delta$.

**Figure 3.** The schematic sequence of applied control fields and of a magnetic field pulse which introduces a phase shift $\delta$ between atomic coherences (left) and the released pulses at the end of the sample plotted for various $\delta$ (right). (In the case presented here the magnetic field modifies, for simplicity, only one of the atomic coherences; such a situation may occur, e.g., for a tripod system composed of Zeeman sublevels in rubidium - for details see [22].)

In Figure 3 the shape of the released signal pulse at the end of the sample is shown for different values of $\delta$ and the interference like behavior of its height is clearly seen (the plots were obtained from numerical solutions of the Bloch-Maxwell equations the details of which are not given in the present general qualitative description, see [22]). For $\delta$ being the odd multiple of $\pi$ the pulse is not released and remains trapped inside the medium in a form of atomic coherences into which it was mapped in the storing stage while for $\delta$ being a multiple of $2\pi$ it attains the maximum. Such a behavior is also reflected by the form of plots (not shown) of the spatial distribution of the corresponding $Z$ polariton after the released pulse has left the medium: their maximum values correspond to those values of $\delta$ for which the pulse is not released. The similar idea of changing the phase of atomic coherence was discussed and verified experimentally in [23] for the simpler case of light storing in a lambda system in which the pulse was stored in a single atomic coherence. The phase of this coherence was changed and the retrieved pulse was made to interfere with a reference beam. In the tripod case the situation is different: the relative phase of two atomic coherences is changed and the interference pattern in the retrieved pulses is due to a superposition of two components inside the medium.

Our second example concerns the case in which the control fields are identical at the storing stage and shifted in time at the release stage. This situation is presented in Figure 4. The
trapped signal pulse is now released in two portions and due to $\dot{\phi} \neq 0$ the $Z \neq 0$ polariton is created which does not leave the sample. Thus a part of the stopped pulse remains stored in atomic coherences and can be further released by another pair of control pulses. It is evident that detailed properties of the pulse leaving the sample can be controlled by a changing the parameters of the storage/release procedure.

Light storing and release in a medium of atoms in the tripod configuration give rise to new interesting phenomena in the case of the input signal being a photonic quantum state. If, for example, two time-separated photons are stored and released by switching on and off two sets of control fields, the tripod medium can behave as a beam splitter in time domain in which the input and output ports concern photons separated in time. It is possible to choose conditions of photon storing and release in which a version of the Hong-Ou-Mandel interference [24] for stored light is realized, i.e., the situation in which two photons are released only either in the first or in the second release stage. Detailed analysis of this effect as well as of steering statistical properties of stored nonclassical light is given in [25] and [26]. Another interesting situation is discussed in [27] in which a possibility has been suggested of producing time-entangled photon state by transferring a part of the excitation created by a trapped photon in one atomic coherence into the other one before a photon release. Certainly investigations concerning new possibilities of controlling quantum properties of the light storage processes are important not only because of their fundamental physical interest but also in the context of quantum information processing.

5. Summary

In summary, we have briefly reviewed the basic ideas of storage and release of laser pulses in atomic media. The phenomenon is based on the possibility of modifications of dispersive properties of a medium via interactions with additional "control" light fields. Due to such interactions an initially opaque medium becomes transparent for resonant radiation, this is the electromagnetically induced transparency. By modifying the control fields in time one can dynamically modify the dispersive properties of a medium and thus manipulate the propagation of light pulses in a controlled way. In particular, the pulse can be stopped (or rather stored)
inside the medium in a form of atomic coherences and later released. The simplest system in which such a scenario can be realized is a three level atomic system interacting with two laser fields in a lambda configuration and the clear physical picture of the process can be formulated in the dark-state polariton formalism. In more advanced systems which contain more laser fields and more active atomic states one has more external parameters to control the process. An example of such a system, namely the tripod configuration has been discussed and the examples of possible controlled manipulations on stored light have been presented.

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