ELECTROMAGNETICALLY INDUCED TRANSPARENCY

One can make opaque resonant transitions transparent to laser radiation, often with most of the atoms remaining in the ground state.

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Electromagnetically induced transparency is a technique for eliminating the effect of a medium on a propagating beam of electromagnetic radiation. EIT may also be used, but under more limited conditions, to eliminate optical self-focusing and defocusing and to improve the transmission of laser beams through inhomogeneous refracting gases and metal vapors, as figure 1 illustrates. The technique may be used to create large populations of coherently driven uniformly phased atoms, thereby making possible new types of optoelectronic devices.

To attain transparency or, at the least, to improve transmission, one applies two laser wavelengths whose frequencies differ by a Raman (nonallowed) transition of a medium. Figure 2 consists of several examples to which the concepts of EIT apply. Figure 2a depicts a prototype three-state system. Most of the experiments described in this article were done in a three-state lead-vapor system of this type. Figure 2b shows a system in which transparency is to be created at an energy that is above the ionization potential of an atom. Figure 2c illustrates a case in which, instead of overcoming absorption, the intent is to reduce the refractive index to as close to unity as possible. To the extent that this can be done, the effect of the medium on the propagating beams is eliminated.

Figure 2a also defines our terminology for the two frequencies that are sent into the medium. Because it is used to probe the absorption from the ground state, one frequency is called the probe frequency and is denoted by ω_p. The second frequency, which couples states $|2\rangle$ and $|3\rangle$, is called the coupling frequency and is denoted by ω_c.

How EIT works

If one asks for a classical explanation of how one may eliminate the effect of a medium on a propagating beam, the answer is that the electrons must be stopped from moving at the frequencies of the applied fields. If the electrons do not move, then they do not contribute to the dielectric constant. Nonmovement will occur if, at each applied frequency, the electron is driven by two sinusoidal forces of opposite phase.

But atoms must be treated quantum mechanically, and in quantum mechanics we deal with probability amplitudes and the expected value of the electron coordinate. In quantum mechanical terms, but in the spirit of the classical explanation, what happens is that the probability amplitude of state $|3\rangle$ of figure 2a is driven by two terms of equal magnitude and opposite sign. One driving term is proportional to the probability amplitude of the ground state $|1\rangle$. The other term is oppositely phased and proportional to the probability amplitude of state $|2\rangle$. These driving terms have the same frequency ω_p and balance such that the probability amplitude of state $|3\rangle$ and the expected value of the amplitude of the sinusoidal motion at each of the applied frequencies is zero.

There is still a subtle question: How does the coherence of the $|1\rangle \rightarrow |2\rangle$ transition atoms become phased so that, together with the applied field at ω_c, it will cause the necessary cancellation? There are two ways in which that may happen. One is that if the applied lasers are monochromatic, the coherence may have an arbitrary phase and, with time, the incorrect component of the phase is damped out. But when we use pulsed lasers, there is not enough time for that to happen. Instead, as explained below, we now know how to initiate the coherence of the $|1\rangle \rightarrow |2\rangle$ transition atoms with the correct phase and to maintain that phase thereafter.

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Early work

In an early experiment, Klaus Boller, Atac Imamoglu and I demonstrated EIT in optically opaque strontium vapor.[1] Figure 3a shows a partial energy-level diagram of strontium vapor. To have a broad absorption linewidth, we chose the upper state of the strontium atom to lie above the first ionization potential. Such an atom decays by autoionization, a process many times faster than radiative decay. Figure 3b plots transmission versus frequency for a probe laser beam propagating alone in the medium. The ratio of transmitted energy to incident energy is about exp(−20). Figure 3c shows the transmission as a function of the probe frequency in the presence of a fixed-wavelength coupling laser. Now the ratio of transmitted energy to incident energy is about 40%.

The physical effect that is the essence of EIT is called coherent population trapping. It was discovered in 1976 by Gerardo Alzetta and his coworkers at the University of Pisa in Italy.[2] In their experiments, the hyperfine states of sodium were the equivalent of states $|1\rangle$ and $|2\rangle$ in figure 2a. Using an elegant technique, the experimenters showed that when the spacing of a multiple of the modes of a multimode laser was coincident with the spacing of the hyperfine states, the fluorescence from state $|3\rangle$ was sharply reduced. The probability amplitudes of each atom (which determine the population of the group of atoms) were driven into a coherent superposition of states $|1\rangle$ and $|2\rangle$, and state $|3\rangle$ was empty. This process is population trapping.

In 1986, Olga Kocharovskaya and Yakov Khanin at the

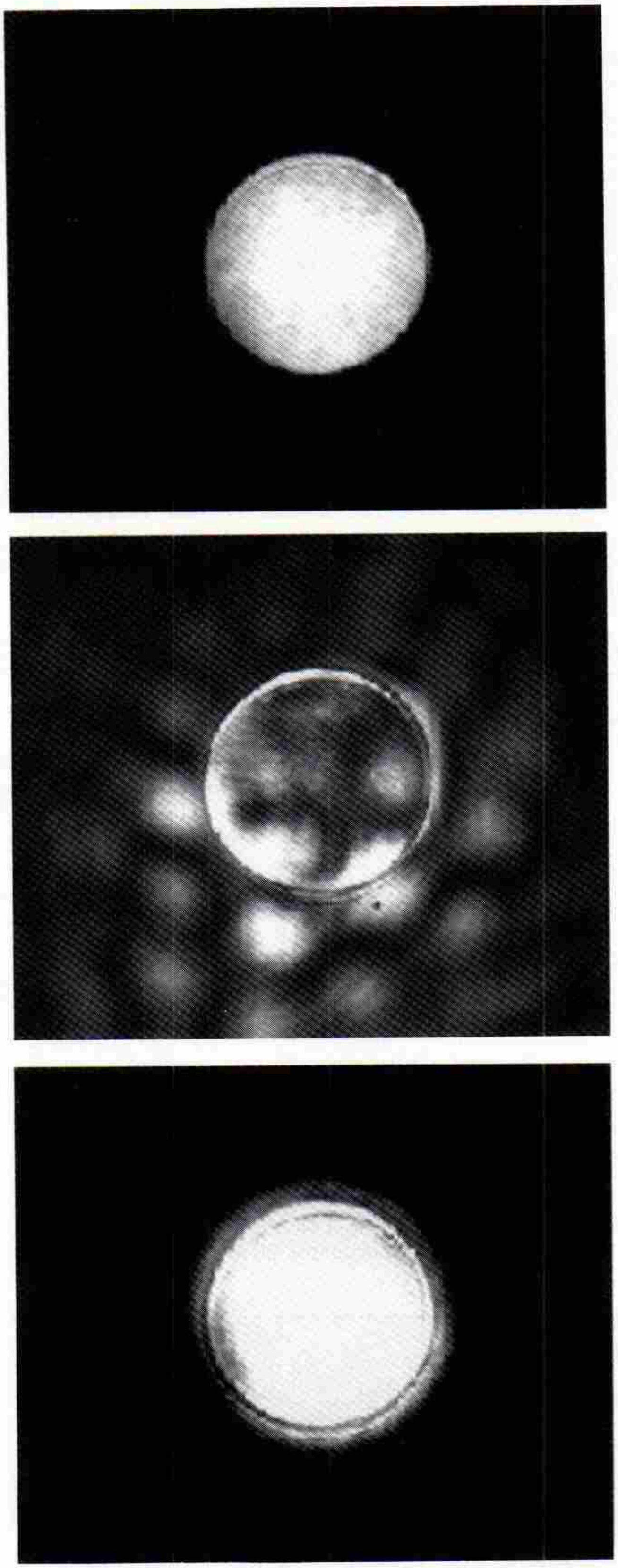

FIGURE 1. ELIMINATION OF OPTICAL SELF-FOCUSING in lead vapor through electromagnetically induced transparency. **a:** Low-intensity probe laser beam that has passed through a 3.2 mm aperture followed by lead vapor. **b:** Probe intensity is increased by a factor of about 10 000; the laser beam filaments and breaks apart in an uncontrolled manner. **c:** The coupling laser is turned on and the image of the aperture is nearly restored. (Adapted from M. Jain, ref. 11.)

Institute of Applied Physics in what is now Nizhni Novgorod, Russia, suggested that a mode-locked train of pulses with a pulse repetition rate equal to the hyperfine splitting of an atom would, in the spirit of the Italian work, establish a population-trapped state and thereby "coherently bleach" the medium. M. B. Gornyi and his colleagues at the Kalinin Leningrad Polytechnical Institute described how, in a dense medium, that could be done by using two monochromatic lasers with different frequencies.[3]

The impetus for work on EIT was the realization that the emission and absorption line shapes of an atom need not be the same. In late 1988, Kocharovskaya and Khanin (working together) and I (working independently) suggested related techniques for making three-state lasers that could operate without the need for a population inversion. Suggestions by Marlan Scully at Texas A&M University and Imamoglu at the University of California, Santa Barbara, followed rapidly and spurred a considerable effort aimed at developing a new understanding of what had been an axiom of laser physics.[4]

Quantum interference

To understand the use of quantum interference, consider first the simple example of atomic hydrogen. Here, instead of using two lasers to establish transparency, one uses a probe laser and a DC field. Figure 4 shows the essential states of atomic hydrogen (neglecting the Lamb shift) with and without a DC field present. The effect of the DC field is to mix and split the otherwise degenerate $|2s\rangle$ and $|2p\rangle$ states. Now, suppose that probe radiation of frequency ω_p is applied at the frequency of the (zero field) $|1s\rangle \rightarrow |2s\rangle$ transition and, therefore, at equal detunings from the mixed states. Also, suppose for the moment that there is no radiative decay. Because the applied frequency is above one resonance and below the other, the contributions of these resonances to the refractive index cancel and, at this frequency, there is no electron motion and no time-varying dipole moment. Now allow for radiative decay. Though not obvious, it turns out that there is still no dipole moment at frequency ω_p. The cause is a more subtle quantum interference that takes place through the interaction with the vacuum field. Quantum mechanics requires that two paths that result in the same end product must interfere.

In general, the dipole-allowed and nonallowed transitions are at different energies, and we must use a second laser instead of a DC field to connect state $|2\rangle$ to state $|3\rangle$. (See figure 2a.) With both lasers tuned sufficiently close to resonance, the Hamiltonian and the quantum interference are the same as for atomic hydrogen. Again, the expected value of the probability amplitude of state $|3\rangle$ is zero, and there is no interaction of the atom with either of the applied fields.[5]

Line shape

If the frequencies of the probe and coupling laser are tuned so that their difference is exactly equal to the Raman transition resonance, then, as explained above, the

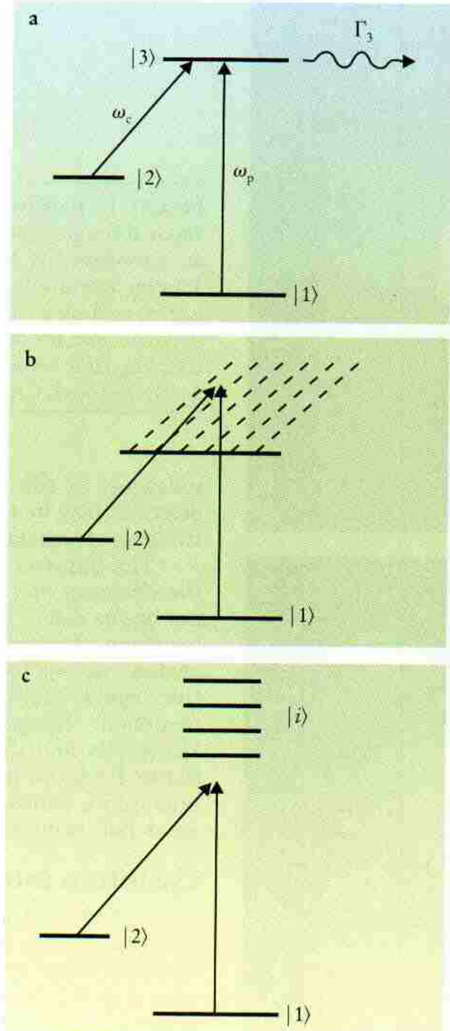

FIGURE 2. SEVERAL EIT SCHEMES. **a:** A three-state system in which the upper state decays with a rate Γ_3 to states outside the system. **b:** Transparency in the continuum. **c:** Use of EIT to modify the refractive index of a medium.

atom and the radiation field will not interact. But if the frequency of either laser is changed so that the Raman resonance condition is no longer satisfied, there will be an interaction. Just as the absorption versus probe frequency line shape of a two-state atom is the well known Lorentzian, the three-state atom with a monochromatic coupling laser (or DC field) has its own characteristic line shape. In the ideal case of zero $|1\rangle \to |2\rangle$ transition linewidth, a weak probe beam and no inhomogeneous broadening, the line shape $L(\delta\omega_p)$ is

$$L(\delta\omega_p) = \frac{4}{\pi}\,\frac{\delta\omega_p}{\Omega_c^2 - 4\delta\omega_p^2 + 2j\gamma_{13}\delta\omega_p} \qquad (1)$$

The quantity $\delta\omega_p = \omega_p - \omega_0$ is the detuning from line center; γ_{13} is the linewidth of the $|1\rangle \to |3\rangle$ transition; and Ω_c is the Rabi frequency of the coupling laser (the Rabi frequency is a measure of the interaction strength of the electric field and the atomic transition). With a transition matrix element μ and an amplitude of the sinusoidal optical field E, the Rabi frequency is $\mu E/\hbar$.

The imaginary part of the line shape is proportional to the medium's loss and is shown in figure 5a; it is zero at line center and varies quadratically as a function of ω_p. The real part, when small, is proportional to the difference of the refractive index from that of free space $(n-1)$ and is shown in figure 5b. Because EIT lets us observe the dielectric constant very close to a resonance transition, the slope of the refractive index is much steeper than usual. As described below, this slope leads to both slow group velocities and the possibility of new types of measurement apparatus.

Experiments that use EIT often require a probe intensity comparable to that of the coupling laser. When that is the case, the quantum interference is more easily understood from the viewpoint of adiabatic preparation. Here, the coupling laser is turned on before the probe laser. That forces the eigenvector of the population-trapped state to coincide with the ground state of the atom. The probe and coupling fields are then increased sufficiently slowly that the system remains in this eigenstate.[6] Even if the decay rate of state $|3\rangle$ is fast in comparison to both Rabi frequencies, if the difference of the field envelopes changes sufficiently slowly, the population will remain trapped in states $|1\rangle$ and $|2\rangle$, and the interaction with the electromagnetic field will remain small.

One should also note the special role of matched pulses. When the pulse envelopes $f(t)$ and $g(t)$ are the same, then no matter how fast the pulses vary, the population-trapped eigenstate of the system will remain unchanged. (See the box below on adiabatic preparation.)

Before discussing propagation, we must understand the time scale on which a population-trapped state may be established in a single atom. This time scale depends on whether one, or both, of the lower states are initially populated. In the experiments of Alzetta's group,[2] these

Adiabatic Preparation

We use a notation system in which the Rabi frequencies Ω_p and Ω_c are constants and the temporal and spatial dependence of the probe and coupling laser pulse shapes are $f(z,t)$ and $g(z,t)$, respectively.

With both lasers tuned close to resonance, the Hamiltonian for the three-state atom of figure 2 with a decay rate Γ_3 from state $|3\rangle$ is

$$H = -\frac{1}{2}\begin{bmatrix} 0 & 0 & \Omega_p f(t) \\ 0 & 0 & \Omega_c g(t) \\ \Omega_p f^*(t) & \Omega_c g^*(t) & j(\Gamma_3/2) \end{bmatrix}$$

As may be verified by inspection, an eigenvector with zero eigenvalue is $[\Omega_c g^*(t), -\Omega_p f^*(t), 0]$. It is this eigenvector that represents the population-trapped state.

To use adiabatic preparation, the coupling laser pulse is applied with the probe pulse still zero. The population-trapped eigenvector is then [1,0,0] and is the same as the ground state of the atom. If both fields are then changed sufficiently slowly, the atom will remain in this eigenstate thereafter.[6]

Because it might be more intuitive to apply the lasers in a sequence that accesses the population at $t = 0$, the type of preparation described here is sometimes called counterintuitive. But from the point of view of quantum interference, it is not counterintuitive. One could ask, Would you make a tunnel, (that is, an interference) and then go through it, or would you first go through it, and then make it?

states were the hyperfine states of sodium, and state $|2\rangle$ was as populated as state $|1\rangle$. In this case the time scale for establishing a population-trapped state is several radiative decay times. In experiments done at Stanford by Klaus Boller and his coworkers at the University of Kaiserslautern in Germany, state $|2\rangle$ was empty.[1] Here, for a small probe, a population-trapped state is established on the time scale of the greater of $1/\Omega_c$ or γ_{13}/Ω_c^2. This makes it possible to use higher-power pulsed lasers, which are necessary to overcome the Doppler linewidth of the $|1\rangle \rightarrow |2\rangle$ transition.

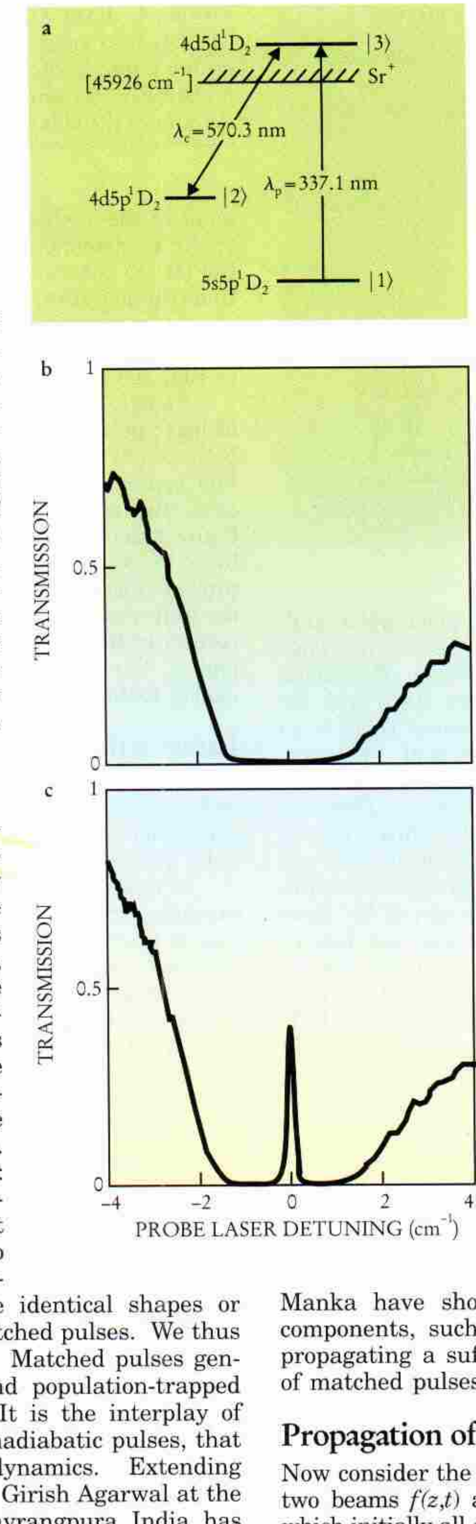

FIGURE 3. ELECTROMAGNETICALLY INDUCED TRANSPARENCY IN STRONTIUM VAPOR.[1] **a:** Partial energy-level schematic. The frequency of a laser of wavelength λ_c is held fixed. Because this laser couples states $|2\rangle$ and $|3\rangle$, it is termed the coupling laser. The frequency of a probing laser of wavelength λ_p is varied. **b:** Transmission versus probe laser detuning with the coupling laser absent. The minimum transmission in this figure is exp(−20). **c:** Transmission versus probe laser frequency with the coupling laser present. The ratio of the transmitted to the incident light is now about 40%. Under the conditions of this experiment, almost all of the atoms remain in the ground state $|1\rangle$.

Atoms drive fields

To understand propagation in an optically (or refractively) thick medium, one must understand not only how the electromagnetic fields drive the atoms, but also how the atoms drive the fields. If the fields—or, more exactly, the difference of the field envelopes $f(t) - g(t)$—varies sufficiently slowly, the atom populations follow the time variation of the fields. But, if the fields vary differently, and sufficiently rapidly, the atoms cannot follow the variation. We then find a striking result: Population-trapped atoms cause arbitrarily shaped optical pulses that are applied at the medium input to generate pulses that, after a characteristic propagation distance, have identical shapes or envelopes.[7] We call such pulses matched pulses. We thus have a basic nonlinear reciprocity: Matched pulses generate population-trapped atoms and population-trapped atoms generate matched pulses. It is the interplay of these processes, especially with nonadiabatic pulses, that leads to interesting propagation dynamics. Extending these ideas to the quantum regime, Girish Agarwal at the Physical Research Laboratory in Navrangpura, India, has shown the matching of photon statistics.[7]

For EIT to work, it must be possible to self-consistently establish the population-trapped state throughout the optically opaque medium. This must occur not only in the direction in which the beam is propagating, but also across the transverse profile, if the beam is not to be distorted by diffractive effects.

Consider the case in which the coupling laser pulse is long enough to be taken as independent of time. As described above, the coupling laser creates the interference for the probe pulse, which propagates at a group velocity determined by the slope in figure 5b. Athos Kasapi and his colleagues at Stanford University have observed well-formed pulses that traveled 10 cm in lead vapor at a velocity of $c/160$. As the pulse rises, energy flows into the population-trapped state, and as it falls, this energy is returned. For a sufficiently slowly rising and falling probe pulse, the probability amplitude of state $|3\rangle$ is small, and even if its decay rate is fast, little energy is lost.[8]

As the strength of the probe pulse becomes larger, an interesting effect occurs: As first shown by Rainer Grobe, Fock Hioe and Joe Eberly, a coupled pulse pair, which they have termed an adiabaton, propagates with a common group velocity.[9] If the time variation of the pulses is slow enough, this pulse pair appears to be stable and may propagate for a long distance. But Michael Fleischhauer and Aaron Manka have shown that, because of its nonadiabatic components, such a pulse pair is not stable and, after propagating a sufficient distance, will evolve into a pair of matched pulses.[10]

Propagation of matched pulses

Now consider the case in which the pulse envelopes of the two beams $f(z,t)$ and $g(z,t)$ are the same at $z = 0$, and in which initially all of the atoms are in the ground state. (The pulses need not have equal Rabi frequencies Ω_p and Ω_c.)

As the matched pulses enter the first layer of atoms, they will not, on the time scale of the Rabi frequency, force the atoms into a population-trapped state. That may be understood from either the quantum interference or adiabatic viewpoints. Both viewpoints require that the coupling laser be applied first, either to establish the quantum interference or to cause the population-trapped eigenstate to initially coincide with the ground state.

So how does EIT work? Surprisingly, the medium sets itself up to become transparent: Because all of the atoms are in the ground state, the probe pulse has a

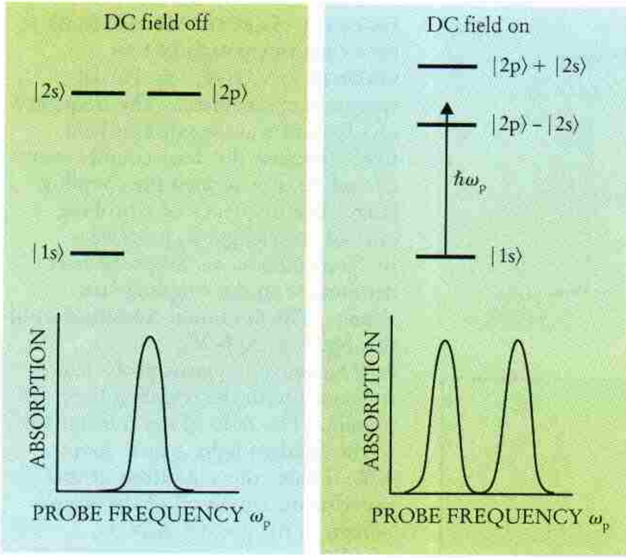

FIGURE 4. BARE AND MIXED STATES OF ATOMIC HYDROGEN. The DC field mixes the bare $|2s\rangle$ and $|2p\rangle$ states and creates quantum interference at the probe frequency. At line center, a probe beam sees zero absorption and not, as one might expect, the sum of the tails of Lorentzian lines.[14]

slower group velocity than the coupling laser pulse and, after propagating a small distance, lags behind it. That is all that is required for adiabatic preparation. The atoms are now driven to the population-trapped state and, for all time thereafter, in this first layer of atoms, there is no further interaction with the electromagnetic field. This process is repeated throughout the remainder of the medium.

Based on a numerical study (see figure 6), Zhen-Fei Luo and I have shown that for an EIT pulse pair to propagate through a medium, the number of photons per unit area in the coupling laser pulse must exceed the number of atoms in the laser path weighted by their oscillator strength. This photon number does not depend on the linewidth and absorption of the medium, nor on whether the objective is to reduce the loss or to modify the refractive index. It is an important limitation on EIT: For short pulses, the required pulse energy varies linearly with the product of the atom density and the length of the medium.[7]

It has now been demonstrated that EIT works sufficiently well that images can be transmitted through otherwise highly absorbing and refracting media. The reason is that EIT automatically allows for different transverse profiles of the incident laser beams. At any (transverse) point in the medium, for arbitrary phases and amplitudes of the electromagnetic fields, the phase and amplitude of the $|1\rangle \rightarrow |2\rangle$ coherence self-adjusts to establish the population-trapped state. That is not the case for other types of transparency-creating mechanisms, as discussed in the box below on EIT and SIT. (For a discussion of the limits of EIT, see the box on page 41.)

Figure 1 shows how one can use a copropagating pair of laser pulses to eliminate optical self-focusing.[11] Here, both lasers are tuned to the high-frequency side of the line center, and in the same spirit as the on-resonance case, they establish a population-trapped state. Part a of figure 1 shows the image of an aperture that is illuminated by a weak probe beam. In part b the intensity of the probe is increased about ten thousand times. That causes the refractive index to become spatially dependent and results in filamentation and destruction of the beam. In part c, the coupling laser is applied and the image is nearly restored.

EIT and SIT

There is another, well-known mechanism for creating transparency in a two-state atom: self-induced transparency, or SIT. Here, a single pulse whose area (the integral of its Rabi frequency over time) is 2π is applied to the medium. Such a pulse causes the atoms to cycle smoothly from the ground to the excited state and then back again. In principle, such a pulse can propagate over long distances. There are several distinctions between EIT and SIT. EIT, but not SIT, is nearly unaffected by spontaneous decay from state $|3\rangle$. Because there are very few atoms in state $|3\rangle$, it does not matter if they decay. The nature of the SIT process makes it difficult to transmit an image, or even a beam, with a Gaussian transverse profile. That is because if the pulse area on the beam axis is 2π, then it will be less than 2π off the axis.

Lasers without inversion and nonlinear optics

From the point of view of EIT, the essence of the laser-without-inversion concept is straightforward: Atoms in a population-trapped state do not interact with the radiation field, and are not to be counted in the population balance.

If new atoms are moved into states $|1\rangle$ or $|2\rangle$—for example, by optical or electron pumping—the atoms cause loss or gain, respectively, at the probe wavelength. The inversion condition in the ordinary laser is then replaced by a condition on the rates of population transfer into states $|1\rangle$ and $|2\rangle$. One might still ask, What about the Einstein A and B coefficients? The answer is that these coefficients carry the assumption that linewidths substantially exceed the Rabi frequencies. Lasers require a departure from equilibrium, but this departure need not occur between the upper and lower laser states.

There is now an extensive literature on lasing without inversion, and there are many different ways to create nonequilibrium situations in which inversion is not required.[4] One experiment by Gunasiri Padmabandu and his coworkers at the Houston Advanced Research Center, works by combining optical pumping and EIT.[12]

Nonlinear optical techniques and devices are now extensively used to generate new wavelengths and to make possible new types of measurement. Examples include harmonic generators and frequency converters, optical parametric oscillators and generators of squeezed light. All of these devices benefit from large nonlinear susceptibilities. But most often, the only way to increase the nonlinear susceptibility is to approach an atomic transition to the ground state. But, as such a transition is approached, the medium exhibits a rapidly increasing refractive index and becomes opaque.

By using the techniques of EIT, even when the product of atom density and length is large, one may closely approach a resonant optical transition, and do so with a refractive index that has a destructive interference and approaches unity at line center. (See figure 5.) The question is then, Will the nonlinear coefficient experience a destructive or a constructive interference at line center? The answer is closely related to the laser-without-inversion problem. If, as part of the nonlinear optical frequency

Can Laser Beams Go Through Walls?

The question sometimes arises: Will electromagnetically induced transparency allow optical beams to be transmitted through complex solids or through kilometers of turbulent molecular gases at atmospheric pressure? The answer is no to the first part and probably no to the second.

As we now understand it, there are at least four limitations on EIT:

▷ As discussed in the main text, the number of photons in the optical pulse must exceed the oscillator-strength-weighted number of atoms or molecules in the laser path.

▷ The peak power of the lasers must be sufficient that the transmission width for EIT exceeds the linewidth of the Raman transition. For broad transitions in solids, field strengths that exceed the breakdown strength of the material would be necessary.

▷ When the Rabi frequencies of both fields are comparable, so as to create large coherences, then both pulse lengths must be short as compared to the dephasing time of the Raman transition.

▷ We do not yet understand how to handle the situation in which many lower states are populated. For example, can the refractive index as caused by many populated rotational levels in molecules be simultaneously reduced by a single pair of applied fields?

There is an additional caution: Almost all of this article has focused on the ideal three-state atom of figure 2a. Peter Lambropoulos has noted that the continuum transparency of figure 2b will not work quite as well. Alexi Sokolov and I have examined the control of refractive indices in molecular hydrogen in which the detunings from the upper states are very large, and we have found that here the refractive indices of the generated comb of Stokes and anti-Stokes sidebands, though substantially reduced, are not equal to that of the vacuum.[18]

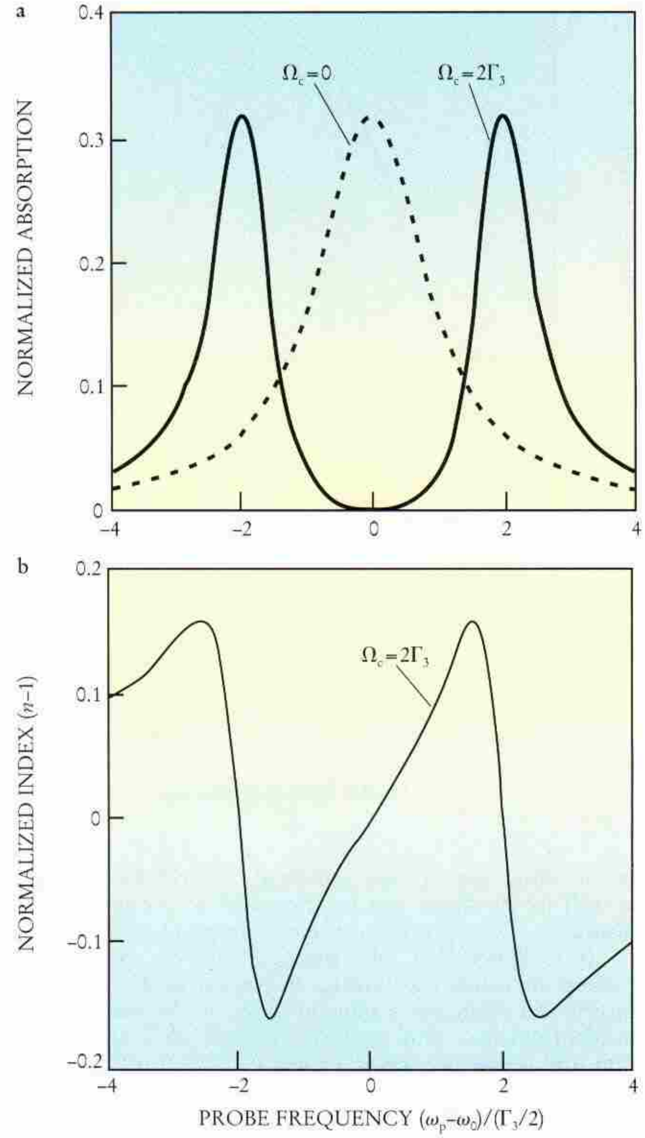

FIGURE 5. IMAGINARY (a) AND REAL (b) PARTS OF THE SUSCEPTIBILITY of a probe frequency ω_p in the presence of a strong-coupling field ω_c. The dashed curve in **a** is the imaginary part of the susceptibility in the absence of the coupling field. The steep slope of the real part of the susceptibility results in the slow group velocity that is characteristic of the preparation phase of EIT. Ω_c is the Rabi frequency of the coupling laser; Γ_3 is the upper-state decay rate.

chain, one accesses the same upper state that one would pump to make an inversionless laser, then the interference is constructive.[13]

The first experiments that demonstrated this constructive interference in the nonlinear susceptibility, with at the same time a destructive interference in the linear susceptibility, were done in atomic hydrogen at the University of Toronto by Kohzo Hakuta and Boris Stoicheff.[14] Recently, Maneesh Jain and his colleagues at Stanford have used EIT to overcome a vanishingly small transmission of exp(−100 000) to create a situation in which the nonlinear polarization is as large as the linear polarization and, therefore, in which nearly complete energy conversion occurs in a single coherence length. Because the product of dispersion and length is so small, this work opens up the possibility of frequency converters and optical parametric amplifiers with very large bandwidths.

Experiments and proposals for nonlinear optics with continuous-wave (cw) lasers are also proceeding: Philip Hemmer at Hanscom Air Force Base has shown how to use population trapping to enhance optical phase conjugation, and Imamoglu and his students have suggested a technique for ultrasensitive phase measurement.[15]

Magnetometers and isotopes

The steep slope of the real part of the refractive index as a function of frequency (figure 5) at a point of vanishing absorption suggests applications to optical interferometry. One example is a magnetometer suggested by Scully and Fleischhauer: If a material exhibiting EIT is placed in one arm of an interferometer, then a change in the magnetic field will cause a Zeeman shift of the nonallowed transition.[16] This results in an unusually large change in path length and improved measurement sensitivity.

There are numerous other possibilities. Min Xiao at the University of Arkansas has shown how cw lasers can be used in a Doppler-free configuration to make a ladder system with unusual dispersive properties.[17] Kasapi has shown how EIT may be used to allow a low-abundance isotope to be seen "behind" an otherwise absorbing isotope.[17] At NIST, in Boulder, Colorado, Alexander Zibrov and his coworkers have experimentally demonstrated that,

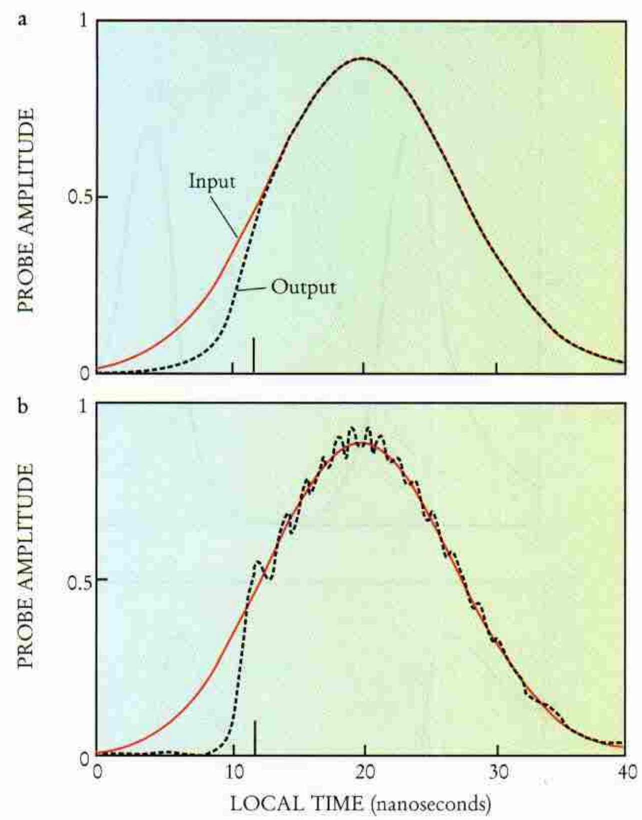

FIGURE 6. ESTABLISHING THE POPULATION-TRAPPED STATE. If one applies matched pulses to an optically thick medium, the pulses will self-consistently establish the population-trapped state, and for all time thereafter will be transmitted without further loss. Both parts of this figure show the pulse at the input (solid curve) and at the output (dotted curve) of a medium. In both parts, the product of atom density and length is the same. But in **a**, the loss to a probe if alone is exp(−5), while in **b**, the loss to a probe if alone is exp(−5000). The large tick mark in each figure denotes the time at which the energy requirement is satisfied. Roughly, at earlier times the medium is being prepared, and at later times it is transparent. (Adapted from S. E. Harris and Z.-F. Luo, ref. 7.)

by including appropriate pumping, the refractive index, as well as its slope, can be increased at a point of zero absorption. This verifies an early proposal of Scully's.[16]

It is likely that EIT concepts will be extended to systems in which the Raman excitation is of a collective nature—for example, a phonon mode or the resonance of an ideal plasma. But in these systems, as in the continuum and refracting media examples of figure 2, EIT will make possible transparency for one, but not both, of the propagating laser beams.[18]

References

1. K.-J. Boller, A. Imamoglu, S. E. Harris, Phys. Rev. Lett. **66**, 2593 (1991). J. E. Field, K. H. Hahn, S. E. Harris, Phys. Rev. Lett. **67**, 3062 (1991).
2. G. Alzetta, A. Gozzini, L. Moi, G. Orriols, Nuovo Cimento B **36**, 5 (1976).
3. O. A. Kocharovskaya, Ya. I. Khanin, Sov. Phys. JETP **63**, 945 (1986). M. B. Gornyi, B. G. Matisov, Yu. V. Rozhdestvenskii, Sov. Phys. JETP **68**, 728 (1989).
4. O. A. Kocharovskaya, Ya. I. Khanin, Sov. Phys. JETP Lett. **48**, 630 (1988). S. E. Harris, Phys. Rev. Lett. **62**, 1033 (1989). M. O. Scully, S-Y. Zhu, A. Gavridiles, Phys. Rev. Lett. **62**, 2813 (1989). A. Imamoglu, S. E. Harris, Opt. Lett. **14**, 1344 (1989). There is also earlier work: V. G. Arkhipkin, Yu. I. Heller, Phys. Lett. **98A**, 12 (1983). Review articles on lasers without inversion are O. Kocharovskaya, Phys. Rep. **219**, 175 (1992); P. Mandel, Contemporary Physics **34**, 235 (1993); M. O. Scully, Quantum Opt. **6**, 203 (1994).
5. For some reviews and papers that discuss coherent population trapping and quantum interference, see the following. P. L. Knight, M. A. Lauder, B. J. Dalton, Phys. Rep. **190**, 1 (1990). B. D. Agap'ev, M. B. Gornyi, B. G. Matisov, Sov. Phys. Usp. **36**, 763 (1993). E. Arimondo, *Progress in Optics*, E. Wolf, ed., Elsevier Science, Amsterdam (1996), p. 257. D. A. Cardimona, M. G. Raymer, C. R. Stroud Jr, J. Phys. B **15**, 55 (1982). A. Imamoglu, Phys. Rev. A **40**, 2835 (1989). G. S. Agarwal, Phys. Rev. A **55**, 2467 (1997).
6. J. Oreg, F. T. Hioe, J. H. Eberly, Phys. Rev. A **29**, 690 (1984). U. Gaubatz, P. Rudecki, M. Becker, S. Schiemann, M. Külz, K. Bergmann, Chem. Phys. Lett. **149**, 463 (1988). C. E. Carroll, F. T. Hioe, Phys. Rev. Lett. **68**, 3523 (1992).
7. S. E. Harris, Phys. Rev. Lett. **70**, 552 (1993). G. S. Agarwal, Phys. Rev. Lett. **71**, 1351 (1993). S. E. Harris, Phys. Rev. Lett. **72**, 52 (1994). E. Cerboneschi, E. Arimondo, Phys. Rev. A **52**, R1823 (1995). S. E. Harris, Z.-F. Luo, Phys. Rev. A **52**, R928 (1995). J. H. Eberly, A. Rahman, R. Grobe, Phys. Rev. Lett. **76**, 3687 (1996). G. Vemuri, K. V. Vasavada, Opt. Commun. **129**, 379 (1996).
8. A. Kasapi, M. Jain, G. Y. Yin, S. E. Harris, Phys. Rev. Lett. **74**, 2447 (1995). A. Kasapi, G. Y. Yin, M. Jain, S. E. Harris, Phys. Rev. A **53**, 4547 (1996). S. E. Harris, J. E. Field, A. Kasapi, Phys. Rev. A **46**, R29 (1992).
9. R. Grobe, F. T. Hioe, J. H. Eberly, Phys. Rev. Lett. **73**, 3183 (1994). J. H. Eberly, M. L. Pons, H. R. Haq, Phys. Rev. Lett. **72**, 56 (1994). J. H. Eberly, Quantum Semiclassical Opt. **7**, 373 (1995).
10. M. Fleischhauer, A. S. Manka, Phys. Rev. A **54**, 794 (1996). I. E. Mazets, B. G. Matisov, Quantum Semiclassical Opt. **8**, 909 (1996).
11. M. Jain, A. J. Merriam, A. Kasapi, G. Y. Yin, S. E. Harris, Phys. Rev. Lett. **75**, 4385 (1995). R. R. Moseley, S. Shepherd, D. J. Fulton, B. D. Sinclair, M. H. Dunn, Phys. Rev. Lett, **74**, 670 (1995).
12. G. G. Padmabandu, G. R. Welch, I. N. Shubin, E. S. Fry, D. E. Nikonov, M. D. Lukin, M. O. Scully, Phys. Rev. Lett. **76**, 2053 (1996).
13. S. E. Harris, J. E. Field, A. Imamoglu, Phys. Rev. Lett. **64**, 1107 (1990). M. Jain, H. Xia, G. Y. Yin, A. J. Merriam, S. E. Harris, Phys. Rev. Lett. **77**, 4326 (1996). A related earlier work is S. P. Tewari, G. S. Agarwal, Phys. Rev. Lett. **56**, 1811 (1986).
14. K. Hakuta, L. Marmet, B. P. Stoicheff, Phys. Rev. Lett. **66**, 596 (1991).
15. J. Donoghue, M. Cronin-Golomb, J. S. Kane, P. R. Hemmer, Opt. Lett. **16**, 1313 (1991). P. R. Hemmer, D. P. Katz, J. Donoghue, M. Cronin-Golomb, M. S. Shahriar, P. Kumar, Opt. Lett. **20**, 982 (1995). H. Schmidt, A. Imamoglu, Opt. Lett. **21**, 1936 (1996).
16. M. O. Scully, M. Fleischhauer, Phys. Rev. Lett. **69**, 1360 (1992). M. O. Scully, Phys. Rev. Lett. **67**, 1855 (1991). A. S. Zibrov, M. D. Lukin, L. Hollberg, D. E. Nikonov, M. O. Scully, H. G. Robinson, V. L. Velichansky, Phys. Rev. Lett. **76**, 3935 (1996).
17. M. Xiao, Y.-Q. Li, S.-Z. Jin, J. Gea-Banacloche, Phys. Rev. Lett. **74**, 666 (1995). J. Gea-Banacloche, Y.-Q. Li, S.-Z. Jin, M. Xiao, Phys. Rev. A **51**, 576 (1995). A. Kasapi, Phys. Rev. Lett. **77**, 1035 (1996).
18. S. E. Harris, Phys. Rev. Lett. **77**, 5357 (1996). S. J. van Enk, J. Zhang, P. Lambropoulos, Phys. Rev. A **50**, 2777 (1994). S. E. Harris, A. V. Sokolov, Phys. Rev. A (1997), to be published. ■