Electromagnetically induced transparency and slow light with n-doped GaAs

T. Wang,[1] R. Rajapakse,[1] and S. F. Yelin[1,2]

[1]*Department Of Physics, University of Connecticut, Storrs, CT 06269*
[2]*ITAMP, Harvard-Smithsonian Center for Astrophysics, Cambridge, MA 02138*

(Dated: September 27, 2005)

The suitability of ground exciton system in semiconductors is studied for use in schemes based on EIT, such as " slow " or " stored " photons or coherent nonlinear optics. We match the desired properties of a system explain EIT with the known physical realities of a semiconductor system, and suggest, in particular, two suitable schemes using donor impurities in GaAs. In addition to generic properties, we also focus on the influence of many neighboring levels and continuum levels, and on the effect of strong hole-mixing.

I. INTRODUCTION

Electromagnetically induced transparency (EIT) [1–3] was recently used to reduce the group velocity of propagating light pulses [4, 5] and to reversibly map them into stationary spin excitations in atomic media [6–9] and for information processing [10]. In particular, an incoming wavepacket can be coherently changed into spin coherence by turning off the control field while the pulse is propagating through the medium. Turning the control field back on reverses the process: the dark-state polariton is adiabatically restored to an optical excitation. The storage time is only limited by the decay of the Raman coherence, which could be very long, of the order of milliseconds [7, 8, 11]. Note that the efficiency of the procedure and the maximum storage interval is determined by the decoherence time of the spin states and by the selection rules that allow for a clean Raman coupling forming a so-called Λ configuration of atomic levels. In addition to fundamental interest, these efforts are stimulated by possibilities to develop a quantum memory for photonic states [12, 13]. During the storage, we can use some electromagnetic fields or even phonons to manage the coherence [9, 14]. This is in essence multiwave mixing with the fields being not applied simultaneously [10]. These techniques may allow for technological advances involving quantum communication over long distances [15]. Whereas the first steps are now being taken toward realization of these techniques in atomic systems [16, 17], the corresponding approaches in solid state would be extremely attractive in view of long-term goals of integration and applications. EIT and laser without inversion (LWI) in semiconductor has been analyzed [18, 19] and EIT in solid has been realized with rare-earth doped matrix [20]. At the same time,much effort has been recently devoted to optically controlled excitonic transitions in semiconductor quantum well with biexcitons [21] and "artificial atoms" quantum dots [22–25]. Because of the fast decay of coherence, the biexciton system in quantum wells is not appropriate for quantum information storage and processing. Although the use of such artificial atoms for quantum control is attractive, it suffers from a number of serious difficulties. First of all,the optical (and electron spin) transition frequencies exhibit large dot-to-dot fluctuations. Secondly,the relatively low symmetry associated with self-assembled quantum dots seriously allow only for unpolarized transitions. Transitions with definite polarizations are essential for a high degree of quantum control in processes involving two-photon Raman transitions and electronic spin coherence.

In semiconductors, free excitons can be bound to donor and acceptor impurity sites, forming bound excitons. In this paper, we explore quantum optical control of electronic and spin degrees of freedom associated with impurity bound excitons in semiconductors. Such excitations hold the promise to be the much better "artificial atoms", because when a bound exciton (BE) decays radiatively to a neutral impurity state, its inhomogeneous linewidth is characteristically narrow [26, 27]. The underling reason for this is that although the impurities might be embedded into the host lattice in slightly different relative locations, the weakly bound electronic states can be delocalized over several hundreds of lattice sites. In such a case, displacements will not shift atomic energies, resulting in very homogeneous optical transition frequencies and g factors. Moreover, in such a case,the symmetry of the host lattice will determine the symmetry of the eigenstates and dipole transitions in the system, and thus can be, for the most semiconductors, assumed to be tetrahedral T_d symmetry. This high symmetry allows for a wide range of well-defined selection rules to hold. This opens the door for studying quantum optical processes in such systems. These two important properties will be used to develop and explore new techniques to control light propagation using electron spin, and controlling electron spin and nuclear spin using light in semiconductors. Specifically, we will analyze the feasibility of EIT via donor bound excitons in GaAs, and its use for controlling light propagation. Before proceeding we note that our work involves electronic degrees of freedom that are associated with localized electrons of dilute ensemble impurities.While pioneering experiments of Awschalom and co-workers [28] have shown exceptional stability of the electronic spin states of free electrons in bulk semiconductors, extremely long T_2 times for optical transitions were indeed demonstrated [26, 27]. Indeed, the coherence time of an electron for an impurity atom can be even longer than that of free electrons in intrinsic semiconductors, because the density of the available final states, into which excitons are scattered as a result of electron-phonon scattering is lower in the $0D$ system.

We begin by considering two generic transitions in donor bound excitons embedded in a semiconductor such as GaAs [29–31]. Note that some experimental results on coherent population trapping have been reported by K.M.C. Fu et al. [32]. To compare with their work, we emphasize the effects of the hole mixing on EIT in bound exciton system. And we also include more related energy levels in our calculation and found that neighboring levels may make the transparency window asymmetric. Also included in our calculations are the dispersion properties which shows how the group velocity may change under the hole mixing.

This paper is organized as follows: In Sec. II we analyze the two sample systems for realizing EIT. Further calculations including other levels and continuum states are discussed in Sec. III. Discussion and conclusion follow in Sec. IV.

II. EIT WITH BOUND EXCITONS

The basic configuration for EIT in a Λ system (Fig. 1a) includes three levels and two lasers with the requirement that the two lasers must have different frequencies and/or polarizations to make sure that each laser only acts on one transition. In the case that the two transitions have similar transition frequencies, different polarizations are required, which means transitions with good experimental selection rules are very important. However, we note that since the probe pulse is usually very weak in EIT system, inadvertently coupling $|1\rangle$-$|3\rangle$ transition by the probe pulse can be tolerated. Note that in this paper, we use the notations in ref. [29] to denote the transitions and levels in our impurity doped semiconductor system. In EIT, one of the most important parameters is the width of the transparency window [33]. The probe will experience transparency only when its spectrum is within the transparency window. A good EIT system requires the smallest possible inhomogenous broadening between all the involved levels as compared with the Rabi frequency of the coupling laser, and a long decay/decoherence time between the metastable states. The latter may be the most important parameter to achieve EIT and light storage, because it determines the depth of the transparency window and the storage time of the stored light.

In expecting the possibility of integration of semiconductors, we try to find appropriate EIT systems in the semiconductors with the above mentioned general rules in mind. For example, because of hole mixing in acceptor states, acceptor bound excitons are excluded from our consideration for EIT. Thus we only consider the transitions for donor bound excitons in semiconductors. The transitions between the neutral donor state, D^0, and the exciton bound to it, D^0X is schematically shown in

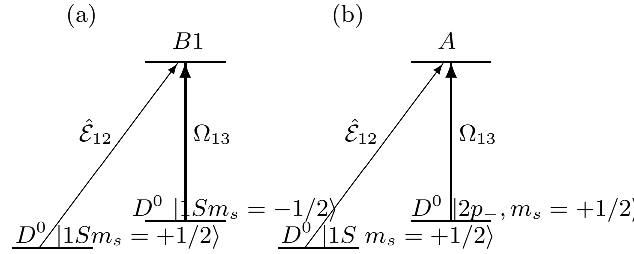

FIG. 1: Three-level atomic systems coupled with a probe field $\hat{\mathcal{E}}_{12}$ and one coupling field Ω_{12}. The actual energy level diagrams in n-GaAs for (a) the principal transitions, and (b) the TES transitions, in magnetic field at $12T$. In (a) the probe transition has π polarization and the coupling laser has σ polarization, which is unpolarized (see text); while in (b), both probe and coupling laser have σ polarizations.

Fig. 2. The donor ground state D^0 consists of the single electron loosely bound to the positively charged donor. We treat the lowest state of the donor exciton D^0X as an excited state, which consists of two electrons on the donor level and one hole in the valence band. The transitions from D^0X to the ground state of D^0 are called the " principle transitions ", and those from D^0X to the excited D^0 states are called the *" two electron states "* (TES). While the model for D^0X remains to be resolved completely, the simple spherical approximation works well in many cases [29]. However, our group theoretic model for T_d symmetry matches the experimental data in [29] for more transitions than the spherical approximation, especially in the case of L=0 and L=1 [34]. Thus in the following we will base our analysis on selection rules which match both experimental data and calculations based on T_d symmetry. Another problem for realizing EIT with bound excitons is that the energy levels are highly degenerate. So polarization selection does not work, and thus electric and/or magnetic field may be needed to separate the levels from each other to avoid unwanted coupling. In the following, magnetic field will be used to separate the levels.

As an specific example, the first system is shown in Fig. 1a which is at the magnetic field of $12T$ in Faraday configuration at liquid-He temperature. The Λ system is composed of ground levels, $D^0\ |-1/2\rangle$ and $|1/2\rangle$, and excited level $D^0X\ B_1$, where B_1 is the first excited state of D^0X with orbital angular momentum $L = 1$, $m_L = 1/2$ and hole quasispin $m_j = 1/2$. The major reason why this level is selected is because B_1 and B_1^* transitions have good selection rules from experimental data [29]. While A_1, which has angular momentum $L = 1$, $m_L = 0$ and $m_j = -1/2$, was identified as having good selection rules, we noticed that this might not be the fact if we look at the experimental data for A_1 of Fig. 18a of ref. [29]. Another reason is that both B_1 and B_1^* transitions are strong, so weaker lasers may be used to lessen the unwanted coupling with other levels (see next section). Since B_1 and B_1^* transitions have σ_- and π polarization, the pump

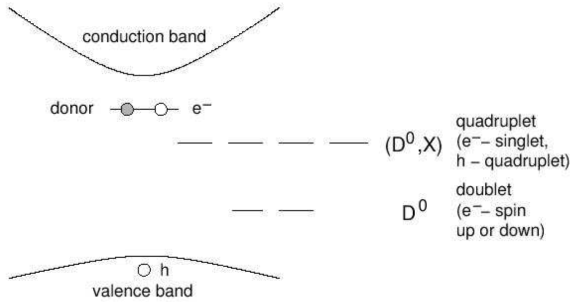

FIG. 2: Energy levels for bound excitons in GaAs. Cartoon of the ground and excited state of a shallow donor-bound exciton. In the ground state, there is one donor electron (grey)with spin up or down resulting in the doublet D^0. In the excited state, an additional electron is excited from the valence band into the donor level, creating a second electron and a hole (both white). The electrons form a singlet state, the hole a quadruplet state, together forming the fourfold degenerate state D^0X.

laser and probe beam are perpendicular to each other. In this case, we identify D^0 1S $|m_s = -1/2\rangle$ as $|2\rangle$, $|m_s = 1/2\rangle$ as $|3\rangle$ and $|B_1\rangle$ as $|1\rangle$. For numerical simulation of this system, we refer the reader to the next section.

The above scheme assumes the ground states of the impurity atoms as the two metastable states of the Λ system. Since the coupling and probe lasers have π and σ_+ polarizations, they are perpendicular to each other, which might be disadvantageous for some application, such as integration. With this in mind, we come to the second scheme to realize EIT as shown in Fig. 1b with Λ system composed of levels $|2p_-, m_s = +1/2\rangle$, D^0 $|m_s = +1/2\rangle$ and A. A is the ground D^0X state with $L = 0$ and $m_j = -3/2$. In this case, since both the pump laser and probe beam have σ_+ polarizations, they are parallel to each other. An additional benefits using the two electron satellite levels is that they are better resolved than the ground states of D^0. Thus the unwanted coupling to neighboring levels is less likely to happen and the asymmetry happens in the first scheme may not happen any more as proven by numerical simulations. Since the dipole moments for these transitions are about 10 times smaller than the principle transitions used in the first scheme, the intensity of the coupling laser need to be 100 time bigger than that for the first scheme if other parameters remain the same. One problem for this scheme is that the population may be trapped to D^0 1S $m_s = -1/2$ state due to optical pumping. This can be overcome by an additional laser to repump the population back to the Λ system. One might also think of a scheme that makes use of level A and TES states. Our calculation shows that unwanted optical pumping may be strong. In addition, because TES states are excited states, initial population within this system is small compared with the previous ones, so we will not analyze this system any further.

III. EFFECTS FROM NEIGHBORING LEVELS

At magnetic field of $12T$, the neighboring levels only perturb the Λ system slightly, because the detuning for the neighboring levels is at least at the order of $0.1\,meV \sim 20GHz$, which is much larger than the Rabi frequency of the coupling laser, which is around $1GHz$ (see below). However, there exists heavy hole-light hole mixing [35, 36]. This makes the decay rate of the Raman coherence ρ_{23} much faster than the exciton recombination rate [37]. For example, σ^+ excitons can be changed to σ^- excitons and 'dark' excitons. $\sigma^\pm$ excitons are excitons coupled to light with $\sigma^\pm$ polarizations and 'dark' excitons are excitons that do not couple to light by optical dipole transition. Suppose the probe light, which couple $|1\rangle$-$|2\rangle$ has σ_+ polarization, then both σ^- (hereafter, 'wrong' excitons) and dark excitons will be dark to the probe light. Thus 'Λ' system using bound excitons should include 3 extra levels (Fig. 1b) to take into account these decoherence. All above are taken into account through the optical Bloch equations including $|1\rangle$-$|7\rangle$ including level $A1$.

Figure 3 shows the effects of hole mixing as well as the effects of ground state decoherence Γ_{23} on the susceptibility $\chi = \chi 1 + i\chi 2$. In the calculations, for simplicity, the hole mixing is assumed to be the same between all holes. Although this is not realistic, we are considering the general effects of hole mixing here. The imaginary part of susceptibility is proportional to absorption coefficient of probe light, from which the transparency window can be found. It can be seen from Fig. 3 (a)-(d) that hole mixing reduces transparency window width dramatically. In particular, the transparency window width without hole mixing is $2GHz$ and when the hole mixing rate is $100 \times 10^9 S^{-1}$, then the width reduces to $0.1GHz$. We note that the real part of the susceptibility becomes steeper as the hole mixing rate gets larger and its pattern outside the transparency window changes qualitatively. We should emphasize that in ideal EIT, no population is on the excited state $|1\rangle$. But we see that hole mixing, which by itself only perturbs $|1\rangle$, does affect the dynamics dramatically. Since the transparency window is reduced and the $\chi 1$ gets steeper, the group velocity becomes much slower than if there were no hole mixing. However, localization of exciton wave function by impurity atom makes the hole mixing suppressed [26, 27]. In this case, polariton effect may be considered as a source of broadening [27] for the bound exciton states and the hole mixing is replaced by this broadening.

As a check for our calculation, when the decoherence for the ground states, $|2\rangle$ and $|3\rangle$, is turned on, Fig. 3(e) shows that even at zero detuning, there is no perfect transparency. Note that for very pure semiconductors, the ground state coherence time could be very large $\sim 60mS$ [11], which makes less absorption in the transparency window. Fig. 3(f) demonstrates that a large spontaneous decay shows qualitatively the same effect as hole mixing.

It is interesting to compare our results with the experimental data (see Fig.3 of ref. [32]). First of all, they use A as $|1\rangle$ while we prefer to use B_1, which has angular momentum $L = 1$, as $|1\rangle$. The reason for this is simple: the oscillator strength for our transitions are larger than those in ref. [32] and thus the laser intensity could be smaller and the unwanted coupling could be reduced. Nevertheless, we tried to fit their data. Figure 4 shows our fit to the experiment [32]. The broadening of the excited with their method, $\gamma_{3a,3b}$, has the same effect as hole mixing of ours, to the first order of optical Bloch equations. And we also see that the group velocity is not very large by reading the $\chi 1$ from fig. 4a. As a final detail, our calculations show that the asymmetry in Fig.3 of ref. [32] may be caused by coupling between D^0 1S +1/2 and $A1$, which has orbital angular momentum $L = 1, m_L = 0$, and hole quasispin $m_j = -1/2$. We also found that without the hole mixing, the asymmetry is becoming much smaller. Also shown in fig. 4 is the case when the ground state broadening is $4\mu S$, which is still much larger than the value measured in experiment [11]. We see that the change of $\chi 1$ within the transparency window is much larger than the previous case and the group velocity is much smaller correspondingly. In a word, the possibility of achieving slow light is very high.

In the above, we assumed resonant coupling laser. Since the decay of the excited state B_1 is large with lifetime being $\sim 1nS$ [38], one might think of detuning the coupling laser to relief this. This is the case as is shown in ref. [39] for manipulating single quantum dot by STIRAP. Although STIRAP and EIT use the same underlining physics, the trapped state, their emphasis are different: STIRAP concerns population transfer which could be done even with single photon detuning [39]; EIT concerns the information in the probe light and the manipulation of it by coupling laser. If the detuning of coupling laser is too large, then the group velocity can not be manipulated effectively by dynamically changing the intensity of the coupling laser, which makes large detuning not feasible for quantum information using slow light. So we assumed resonant coupling laser.

Another consideration involves the influence of the transitions from bound electronic states into an electron-hole continuum,which is important because of its low binding energies $\sim 1meV$. Although issue of the effect of continuum states on the dark-resonances has been a subject of spirited debates in the past [40], the detailed study [41] shows that a system with two discrete levels only coupled to pure continuum states can have EIT and light storage, and if the lasers are detuned from the continuum, then the continuum will only gives Ac Stark shift $\delta E = -\int_0^\infty d\epsilon \frac{|\Omega|^2}{D}$ [39, 41].

Now let us estimate the intensity of the coupling lasers. To this end, we need the dipole moments of the transitions. The dipole moments of B_1 and B_1^* transitions are about 100 *Debye* [38, 42]. With such, the estimated intensity of the coupling laser is $I \sim 0.5W/cm^2$, which is consistent to ref. [32]. Following the above line, we found that scheme 2 which makes use of TES states need coupling laser intensity about 100 times larger that for the principle transition.

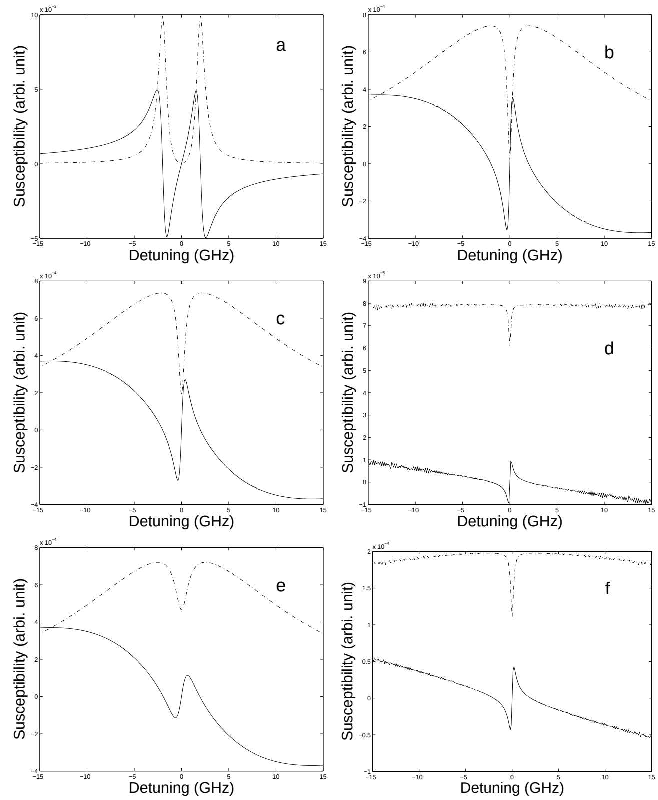

FIG. 3: Real part (solid line) and imaginary part (dash-dotted) of susceptibility as function of probe detuning for resonant coupling field. Resonant coupling laser has Rabi frequency $\Omega_{13} = 2 \times 10^9 S^{-1}$, spontaneous decay rates from bright bound excitons are $1 \times 10^9 S^{-1}$. (a) without hole mixing $\Gamma_h = 0S^{-1}$ and ground state decoherence rate $\Gamma_{23} = 0.01 \times 10^9 S^{-1}$, (b) hole mixing rate is $\Gamma_h = 10^9 S^{-1}$. (c) $\Gamma_h = 10 \times 10^9 S^{-1}$ and $\Gamma_{23} = 0.1 \times 10^9 S^{-1}$, (d) $\Gamma_h = 100 \times 10^9 S^{-1}$ and $g_{23} = 0.1 \times 10^9 S^{-1}$. (e) $\Gamma_h = 100 \times 10^9 S^{-1}$ and $\Gamma_{23} = 0.5 \times 10^9 S^{-1}$. (f) $\Gamma_h = 0S^{-1}$ and $\Gamma_{23} = 0.5 \times 10^9 S^{-1}$ with $\Gamma_{12} = 100 \times 10^9 S^{-1}$. It can be seen from (a)-(d) that hole mixing reduce transparency window width dramatically (from the imaginary part of χ). Also, although the pattern of real susceptibility, and thus the index of refraction, do not change qualitatively within the transparency window, its patten outside the transparency window changes a lot. With the decoherence for the ground states, $|2\rangle$ and $|3\rangle$, (e) shows that even at zero detuning, there is no perfect transparency. (f) demonstrates that a large spontaneous decay shows qualitatively the same the effect as hole mixing.

IV. DISCUSSION AND CONCLUSION

Since the doping density in our proposed system is low ($\sim 10^{14} cm^{-3}$), bound-exciton bound-exciton interaction is not important. So microscopic many-particle theory

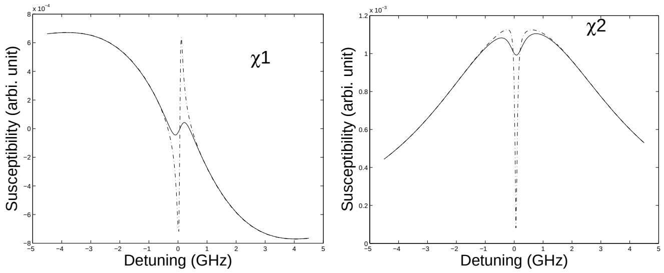

FIG. 4: (a) the real part ($\chi 1$) and (b) the imaginary part ($\chi 2$) of susceptibility for the experiment in ref. [32]. The fitting parameters used are $\Omega_{13} = 0.065GHz$, $\Omega_{12} = 0.0016$, $\gamma_{13} = 1nS^{-1}$, $\gamma_{12} = 0.08\gamma_{13}$. $\gamma_{23} = (4nS)^{-1}$ is used for the fit to the experimental data in ref. [32] (solid curve), while dashed curve is used for the case that the broadening between ground states is much smaller $\gamma_{23} = 4\mu S$. Asymmetry in $\chi 2$ due to neighboring levels is clearly seen.

to include exciton-exciton correlation as used in ref. [21] may be not necessary. We also note since the inhomogeneous broadening is small, we did not include it in our calculation, although this can be done by following the lines in ref. [43].

In the above schemes, magnetic field of $12T$ is used. Can we get rid of the magnetic field? Since strain can split heavy/light holes, well controlled strain may be used to realizing the EIT system [34].

In conclusion, electromagnetically induced transparency (EIT) and slow light are analyzed in a n-doped GaAs system. Two Λ systems are identified in this system, where shallow impurity bound exciton states are used as excited states with lower states being metastable states. We found that hole mixing and ground state broadening may dramatically reduce the transparency window and reduces the group velocity of the probe light. Further experiments toward reducing the broadening of ground states are expected.

ACKNOWLEDGMENTS

The authors gratefully acknowledge the discussions with Y. Yamamoto, and the financial support from NSF.

[1] K. J. Boller, A. Imagmoğlu, and S. E. Harris, Phys. Rev. Lett. **66**, 2593 (1991).
[2] S. E. Harris, Physics Today **50**, 36 (1997).
[3] M. D. Lukin and A. Imamoglu, Nature **413**, 273 (2001).
[4] M. M. Kash, V. A. Sautenkov, A. S. Zibrov, L. Hollberg, . M. D. L. George R. Welch, Y. Rostovtsev, E. S. Fry, and M. O. Scully, Phys. Rev. Lett. **82**, 5229 (1999).
[5] L. V. Hau, S. E. Harris, Z. Dutton, and C. H. Behroozi, Nature **397**, 594 (1999).
[6] M. Fleischhauer and M. D. Lukin, Phys. Rev. Lett. **84**, 5094 (2000).
[7] C. Liu, Z. Dutton, C. H. Behroozi, and L. V. Hau, Nature **409**, 490 (2001).
[8] D. F. Phillips, A. Fleischhauer, A. Mair, R. L. Walsworth, and M. D. Lukin, Phys. Rev. Lett. **86**, 783 (2001).
[9] A. S. Zibrov, A. B. Matsko, O. Kocharovskaya, Y. V. Rostovtsev, G. R. Welch, and M. O. Scully, Phys. Rev. Lett. **88**, 103601 (2002).
[10] T. Wang, M. Koštrun, and S. F. Yelin, Phys. Rev. A **70**, 053822 (2004).
[11] A. M. Tyryshkin, S. A. Lyon, A. V. Astashkin, and A. M. Raitsimring, Phys. Rev. B **68**, 193207 (2003).
[12] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Phys. Rev. Lett. **78**, 3221 (1997).
[13] M. D. Lukin, S. F. Yelin, and M. Fleischhauer, Phys. Rev. Lett. **84**, 4232 (2000).
[14] R. Walsworth, S. F. Yelin, and M. D. Lukin, Opt. Phot. News. **13**, 50 (2002).
[15] L. M. Duan, M. D. Lukin, J. I. Cirac, and P. Zoller, Nature **414**, 413 (2001).
[16] A. Kuzmich, W. P. Bowen, A. D. Boozer, A. Boca, C. W. Chou, L. M. Duan, and H. J. Kimble, Nature **423**, 731 (2003).
[17] C. H. V. der Wal, M. D. Eisaman, A. Andre, R. L. Walsworth, D. F. Phillips, A. S. Zibrov, and M. D. Lukin, Science **301**, 196 (2003).
[18] D. E. Nikonov, A. Imamoğlu, and M. O. Scully, Phys. Rev. B **59**, 12212 (1999).
[19] T. Li, H. Wang, N. Kwong, and R. Binder, Opt. Exp. **24**, 3298 (2003).
[20] A. V. Turukhin, V. S. Sudarshanam, M. S. Shahriar, J. A. Musserm, B. S. Ham, and P. R. Hemmer, Phys. Rev. Lett. **88**, 023602 (2002).
[21] M. C. Phillips, H. Wang, I. Rumyantsev, N. H. Kwong, R. Takayama, and R. Binder, Phys. Rev. Lett. **91**, 183602 (2003).
[22] P. Michler, A. Kiraz, C. Becher, W. V. Schoenfeld, P. M. Petroff, L. Zhang, E. Hu, and A. Imamoglu, Science **290**, 2282 (2000).
[23] J. Kim, O. Benson, H. Kan, and Y. Yamamoto, Nature **397**, 500 (1999).
[24] X. Li, Y. Wu, D. G. Steel, D. Gammon, T. H. Stievvator, D. S. Katzer, D. Park, C. Piermarocchi, and L. J. Sham, Science **301**, 809 (2003).
[25] R. I. Dzhioev, V. L. Korenev, I. A. Merkulov, B. P. Zakharchenya, D. Gammon, A. L. Efros, and D. S. Katzer, Phys. Rev. Lett. **88**, 256801 (2002).
[26] D. Karaiskaj, M. L. W. Thewalt, T. Ruf, M. Cardona, H.-J. Pohl, G. G. Deviatych, P. G. Sennikov, and H. Riemann, Phys. Rev. Lett. **86**, 6010 (2001).
[27] H. Oohashi, H. Ando, and H. Kanbe, Phys. Rev. B **54**, 4702 (1996).
[28] J. M. Kikkawa, I. P. Smorchkova, N. Samarth, and D. D. Awschalom, Science **277**, 1284 (1997).
[29] V. A. Karasyuk, D. G. S. Beckett, M. K. Nissen, A. Villemaire, T. W. Steiner, and M. L. W. Thewalt, Phys. Rev. A **49**, 16381 (1994).
[30] F. A. J. M. Driessen, H. G. M. Lochs, S. M. Olsthoorn, and L. J. Giling, J. Appl. Phys. **69**, 906 (1991).
[31] W. Ruhle and W. Klingenstein, Phys. Rev. B **18**, 7011 (1977).

[32] K.-M. C. Fu, C. Santori, C. Stanley, M. Holland, and Y. Yamamoto (2005), cond-mat/0504012.
[33] M. Fleischhauer and M. D. Lukin, Phys. Rev. A **65**, 022314 (2002).
[34] S. F. Y. R. Rajapakse, to be published.
[35] T. Uenoyama and L. J. Sham, Phys. Rev. Lett. **64**, 3070 (1990).
[36] K. Bott, O. Heller, D. Bennhardt, S. T. Cundiff, P. Thomas, E. J. Mayer, G. O. Smith, R. Eccleston, J. Kuhl, and K. Ploog, Phys. Rev. B **48**, 17418 (1993).
[37] S. A. Crooker, D. D. Awschalom, J. J. Baumberg, F. Flack, and N. Samarth, Phys. Rev. B **56**, 7574 (1997).
[38] G. Hooft, W. V. der Poel, L. molenkamp, and C. Foxon, Phys. Rev. B **35**, 8281 (1987).
[39] P. C. Chen, C. Piermarocchi, L. J. Sham, D. Gammon, and D. G. Steel, Phys. Rev. B **69**, 075320 (2004).
[40] K. Bergmann, H. Theuer, and B. W. Shore, Rev. Mod. Phys. **70**, 1003 (1998).
[41] T. Wang, L. Zhao, and S. F. Yelin, in preparation.
[42] J. R. Guest, T. H. Stievater, X. Li, J. Cheng, D. G. Steel, D. G. D. S. Katzer, D. Park, C. Ell, A. Thränhardt, G. Khitrova, et al., Phys. Rev. B **65**, 241310R (2002).
[43] E. Kuznetsova, O. Kocharovskaya, P. Hemmer, and M. O. Scully, Phys. Rev. A **66**, 063802 (2002).