A theoretical description of polaritons in quasi-one-dimensional microcavities is given where the nonlinear effects due to Coulomb exciton-exciton interaction and phase space filling are included explicitly. Analytical results are obtained for the power and time dependences of the optical response including anomalous amplification at short delay times and scattering between quasi-one-dimensional subbands from the lower as well as from the upper polariton branches. Spin quantum beats for the long delay times are shown to be consistent with exchange splitting of optically active excitons induced by uniaxial strain. The power and exchange splitting dependences of this response are obtained. The nonlinear approach describing the optical effects for both time scales are similar and include the exchange interaction for longer delay times.

I. INTRODUCTION

Semiconductor planar microcavities are produced by embedding a semiconductor dielectric layer with multiple quantum wells as active material between two high reflectance Bragg mirrors. The interaction between the photon modes and the quantum well excitons can give rise to strong coupling between these excitations and anticrossing of their energies. The coupled excitations, called microcavity polaritons, are quasi-two-dimensional due to the translational invariance of the embedded quantum well and to the perpendicular confinement.

Recently, several studies have reported strong nonlinear emission from these planar microcavities in the strong-coupling regime. In particular, photoexcited polaritons at the inflection point \( k_0 \) of the lower polariton branch scatter into two macroscopically occupied states, the probe \( k = 0 \) and an idler \( 2k_0 \). This occurs when conservation of energy and momentum are satisfied for these three points, \( 2E_{k_0} = E_0 + E_{2k_0} \). A high optical amplification at \( k = 0 \) is observed at short delay times between the pump and the probe signals (a few ps). This phenomenon results from the exciton scattering in the quantum well, and it provides a powerful and detailed probe into exciton dynamics and their coupling with light. Recently, a theoretical approach for these effects in layered cavity systems was developed.\(^{2}\)

Additional insight can be gained into the dynamics of these interacting systems by modifying the photon modes. This can be done by using one-dimensional cavities in the shape of photonic wires formed by lateral fabrication of the layered cavities on the scale of the wavelength of light. This leads to discretization of the photon modes, and of the corresponding polaritons, into one-dimensional subbands.\(^{3}\) In recent work\(^{4}\) nonlinear polaritons have been reported for such photonic wire cavities where many different scattering channels were observed. There it was found that the scattering processes involve different sub-branches. The optical gain occurred at short delay times and at energies similar to the case for a planar cavity.

In related experiments on these photonic wires it was reported that a different type of nonlinear emission occurred.\(^{5}\) Optical amplification was observed at longer delay times, on the order of tens of ps, for the probed state \( k = 0 \) when a pump was applied at \( k = k_0 \) and there is no macroscopically occupied idler state. It was found that there were pronounced oscillations of polariton spin states as a function of delay time and excitation power. These experiments suggested that there is an uniaxial in-plane strain due to lateral patterning of the photonic wire. As a result an exchange splitting between the spin-up and -down polariton states occurs. As the excitation power was increased, the polariton ground state became macroscopically occupied, and the spin coherency in the system was maintained.

In this paper we address these types of experiments in photonic wires. We use a Hamiltonian that includes exciton-exciton interactions and phase space filling effects due to higher excitation powers in the system.\(^{6}\) To explain the quantum spin beats, we include an exchange splitting between the spin-up and -down polariton states. Microscopic equations are derived and solved analytically. The dependence of the optical response of the photonic wire on the pump excitation power, delay times, and polariton broadening are obtained. Good agreement between the experimental data\(^{4,5}\) and the theoretical results is obtained.

II. ELASTIC POLARITON SCATTERING IN PHOTOニック WIRES

Angle-resolved pump-probe experiments were performed in photonic wires similar to those in planar cavities.\(^{4}\) In these quasi-one-dimensional structures the photon degree of freedom becomes quantized, and the energy is approximately \( E_{i}^{\text{phot}} = (E_{i}^{e} + \hbar^2 c^2 \pi^2 i^2 / \epsilon L^2 + \hbar^2 c^2 k^2 / \epsilon)^{1/2} \) with \( i = 1, 2, 3, \ldots \), \( c \) is the speed of light, \( L \) is the length of the wire, \( \epsilon \) is the dielectric constant, and \( k \) is the quasi-one-dimensional wave vector of the photons.\(^{3}\) Thus the polariton state is characterized by a photon subband index \( i \) and the continuous one-dimensional (1D) wave vector \( k \). Pump-probe experiments\(^{5}\) show that again a high optical gain is observed at \( k = 0 \) and at the idler state \( 2k_0 \), but for photonic wires the energies can be located on different subbranches from the pump. Here the energy conservation condition for the polaritons, \( \epsilon_{n,0} + \epsilon_{i,2k_0} = 2 \epsilon_{m,k_0} \), must be satisfied. The
The starting Hamiltonian includes Coulomb exciton-exciton interactions and by a continuous transformation for the excitons with operators $a_{X,0}$ the energy conservation. In the photonic wire, the three branches, which is forbidden in the planar cavity because of the energy conservation. In the photonic wire, the three states, $k = 0$, $k_0$, and $2k_0$, can be located on a number of lower and upper subbranches. For example, in one case studied in Ref. 4 the pump was at $k_0 = 1.4 \mu m^{-1}$, $m = 3$ in the lower polariton branch (LPB), the probe was at $k = 0$, $n = 1$ LPB, and the idler was at $2k_0$, $l = 1$ in the upper polariton branch (UPB).

The observed optical amplification of the polariton modes in a planar cavity has been described theoretically using an interacting polariton model including only scatterings within the LPB. Here we use a similar approach for the photonic wires. One of our particular interests will be scattering between the LPB and the UPB made possible in these systems. The starting Hamiltonian includes Coulomb exciton-exciton interactions and phase space filling at increased excitation powers. The polariton energies are characterized by an integer $i$ corresponding to the subbranches and by a continuous 1D wave vector $k$.

We consider only the process $(m,k_0; m,k_0) \rightarrow (n,k = 0; l,2k_0)$. Such a treatment is appropriate because all three states are macroscopically occupied. The interacting Hamiltonian for the excitons with operators $B_{i,k}$ and photons with operators $a_{j,q}$ is given by

$$H = \frac{1}{2} \sum_{ijlm,kl'} V_{ijlm,kl'} \left[ \frac{A^2}{\hbar} B_{i,k+q}^+ B_{j,k- \ell}^+ B_{l,k} B_{m,k} \right] - \sum_{i,m,kl'} \frac{\Omega R}{\hbar} a_{X,0}^2 \left[ a_{X,i+q}^+ a_{X,j- \ell}^+ a_{X,l} a_{X,m} + H.c. \right],$$

(1)

where $V_{ij}$ is the Coulomb interaction, $n_{sat}$ is the saturation density, $a_X$ is the 2D Bohr radius, $\Omega_R$ is the Rabi splitting energy, and $A$ is the surface area. Performing the Hofield transformation $L_{i,k} = X_{i,k} B_{i,k} - C_{i,k} a_{i,k}$ for the LPB and $U_{i,k} = C_{i,k} B_{i,k} + X_{i,k} a_{i,k}$ for the UPB, we can write the Hamiltonian in terms of the polariton operator $L_{i,k}$ and $U_{i,k}$. We also have to add the Hamiltonian from the external sources $H = \sum_{i,k} g_k \left[ \dot{X}_{i,k} U_{i,k}^+ - C_{i,k} L_{i,k}^+ \right]$, where $g_k = g_0$ or $g_{k_0}$ for the probe and pump. Then we can write the equations for the dynamics of the probe, pump, and idler states. Here we give explicitly the case when the pump and probe are at the LP subbranches and the idler is on an UP subbranch.

$$L_{m,k_0} = \frac{i}{\hbar} \left[ (E_{m,k_0} + i\gamma) L_{m,k_0} + 2V_{nml,kl'} g_k L_{m,k_0}^+ L_{l,k_0} \right],$$

(3)

$$E_{n,k_0} = \frac{i}{\hbar} \left[ (E_{n,k_0} + i\gamma) U_{i,k_0} + V_{nml,kl'} g_k L_{m,k_0}^+ L_{l,k_0} \right],$$

(4)

$$V_{nml,kl'} = V_{0} \left[ a_{X,n}^2 X_{m,k} + C_{l,k}^2 \left( \frac{\xi}{nA} X_{m,k} C_{l,k} C_{l,k_0} \right)^2 \right],$$

(5)

where $\xi = -a_X g_0 D_{m,0} / \sqrt{A}$ is the probe source and $D_{m,0} = -g_k \chi_{m,k_0} (t) C_{m,k_0} / \sqrt{A}$ is the pump source. Dots indicate time derivatives. The polariton linewidth due to additional losses in the cavity is given by $\gamma$. Similar equations can be written for other cases when, for example, all three states are in the LPB. The differences will be in the expressions for the energy and the matrix element $V_{nml,kl'}$.

Useful insight can be gained by considering the steady-state solution for the polariton occupations, which can be obtained by substituting $L_{m} = \bar{L}_m e^{i\omega_m}$, $U_{l} = \bar{U}_l e^{i\omega_l}$, and $\bar{L}_l = \bar{L}_l e^{i(\omega_0 - \omega_l)}$. We find, as in planar cavities, that there is a threshold pump power expressed as $|\bar{L}_m|^2 \sim \gamma V$. Also, the gain is blueshifted by $\delta \omega_m = V |\bar{L}_m|^2 \sim \gamma$. The parameter $\gamma$ is a measure of the losses in the cavity, while $V = V_{nml,kl'}$ comes from the nonlinear nature of the scattering processes. In the photonic wire $V$ will be different for each process and thus the threshold pump density will be also different.

Solutions for the probe and idler state also can be obtained as a function of time. This is done by turning the system of Eqs. (2)-(4) into two independent second-order differential equations for the probe and idler states. An expression for the probe state at higher excitation power when the energy for the wave mixing is conserved is given by

$$\bar{L}_n = \frac{2i}{\hbar} \left[ (E_n + i\gamma) L_n + V^2 |L_n|^2 \right] + \frac{\gamma + \gamma}{\hbar} \left[ (E_n + i\gamma)(E_n + 3i\gamma) + \gamma \right] \frac{\bar{L}_n}{\hbar^2},$$

(7)

(2) The solution is easily found to be
Having the explicit formula for the polariton state, one can perform the Fourier transformation

$$L_n \sim e^{-2\gamma t/\hbar} \left( 1 - \frac{1}{2} (1 + i \delta) e^{iE_n/\hbar - \xi t/\hbar} \right. $$

$$\left. - \frac{1}{2} (1 - i \delta) e^{iE_n/\hbar + \xi t/\hbar} \right),$$

(8)

$$\delta = \frac{E_n}{\sqrt{E_n^2 + \xi^2}},$$

(9)

$$\xi^2 = \gamma^2 + V^2 |L_m|^4.$$  

(10)

FIG. 1. (a) Optical gain versus energy for scattering between pump and probe in the lowe polariton branch and idler in the upper polariton branch. (b) Optical gain as a function of pump-probe delay time for three values of the polariton linewidth.

III. EXCHANGE SPLITTING OF POLARITONS

The previous section dealt with polaritonic scattering processes in “photonic wires” for which a high optical gain is found at short delay times. Those processes are described in terms of energy and momentum conservation. Recently, related experiments on photonic wires also found optical gain at much larger delay times, on the order of tens of ps.\(^5\) To suppress the polariton-polariton scattering channel involving the probe \(k = 0\), pump \(k = k_0\) and idler \(2k_0\) states, the pump excitation was applied at different \(k = k_0'\) in which case there is not a macroscopically populated idler state. The observed gain at \(k = 0\) is much smaller than the gain for the polariton amplifier described in the previous section. In addition, quantum spin oscillations for the spin-up and -down polariton states as a function of delay time and excitation power were observed.

The spin sensitivity of this polariton scattering is attrib-
uated to an exchange splitting between the optically active exciton states in the structure. It was suggested\(^4\) that the uniaxial strain was introduced by the lateral patterning in those structures. The strain gives rise to a reduction of the symmetry of excitons in the quantum well. The energy of the optically active heavy-hole states for the quantum well excitons can be described as\(^7\)

\[
E_1 = \frac{c_x}{2} + \frac{1}{2}(c_x + c_y),
\]

\[
E_2 = \frac{c_x}{2} + \frac{1}{2}(c_x + c_y),
\]

where \(c_{x,y} \) are the electron-hole exchange interaction parameters as a function of the quantum well width. An ideal quantum well has a \(D_{2d} \) symmetry, which has a fourfold rotation-reflection axis, and thus \(c_x = -c_y \). Therefore the optically active states are degenerate. When the symmetry is broken, as is the case when there is strain, then \(c_x \neq c_y \) and an exchange splitting between the states appears. From the pump-probe experiments in Ref. 5 an exchange splitting of 130 \(\mu\)eV was estimated using the period of the beats.

It was observed that at higher excitation power there are spin quantum beat as functions of the delay time between pump and probe and excitation power. The experiments were done for polaritons from the LP subbranch \(i = 1 \), and they suggested that there is a macroscopic population of each spin state at \(k=0 \) and \(k=k'_0 \). The scattering was attributed to a stimulated transfer from the reservoir at the excitation point \(k'_0 \).

To describe these processes we start with the linear Hamiltonian

\[
H = H_0 + H_{\text{exch}},
\]

\[
H_0 = \sum_{i,k} \epsilon_{i,k} P_{i,k}^+ P_{i,k},
\]

\[
H_{\text{exch}} = (\Delta \chi_{i,k}^2) e^{i\phi_{i,k}} P_{i,k}^+ P_{i,k} + e^{-i\phi_{i,k}} P_{i,k} P_{i,k}^+.
\]

where \(H_0 \) is the free polariton Hamiltonian\(^2\) with energy \(\epsilon_{i,k} \) and \(H_{\text{exch}} \) is the exchange Hamiltonian\(^8\) with \(\Delta \) the exchange splitting. Note that here we keep only the polariton operators for the LPB. \(H \) can be diagonalized by the transformation

\[
P_{i,k} = e^{-i\phi_{i,k}/2} (\alpha_{i,k} + \beta_{i,k}) / \sqrt{2} \quad \text{and} \quad P_{i,k}^+ = e^{i\phi_{i,k}/2} (\alpha_{i,k} - \beta_{i,k}) / \sqrt{2}.
\]

Then the Hamiltonian becomes

\[
H = \sum_{i,k} \left[ (\epsilon_{i,k} - \Delta \chi_{i,k}^2) \alpha_{i,k}^+ \alpha_{i,k} + (\epsilon_{i,k} - \Delta \chi_{i,k}^2) \beta_{i,k}^+ \beta_{i,k} \right].
\]

The contribution of the pump and probe sources must also be added \(H_{\text{ext}} = (\alpha_{i,k} / \sqrt{A}) [G_{01}(t) C_{01} P_{i,k}^+ + G_{10}(t) C_{10} P_{i,k} + G_{01}(t) C_{01} P_{i,k}^+] \), where \(G_{01}(t) \neq 0 \) when the probe is with \(\alpha_{i,k}^+ \) polarization.

The nonlinear Hamiltonian describing the polariton dynamics is that in Eq. (1). First the Hopfield transformation is performed keeping only the contribution from the lower polariton subbranches. Then the transformation is done to change to the new states \(\alpha_{i,k} \) and \(\beta_{i,k} \).

Now we consider interactions between only the pump at \(k'_0 \) and the probe polaritons at \(k=0 \) \((i = 1) \). This simplification can be made when a discrete set of polariton modes is coherently and macroscopically occupied as was the case in the experiments. This allows us to neglect scattering channels that do not have macroscopically occupied modes as initial or final states. Then the equations of motion for the \(\alpha \) polariton states are

\[
\dot{\alpha}_0 = -\frac{i}{\hbar} \left[ (\epsilon_0 + \Delta \chi_{00}^2 + \Delta \chi_{01}^2 + \epsilon_1 + \epsilon_2) \alpha_0 + V(\alpha_{k_0}^+ \alpha_{k_0} + \beta_{k_0}^+ \beta_{k_0}) \alpha_0 \right]
\]

\[
+ V(\alpha_{k_0}^+ \beta_{k_0} + \beta_{k_0}^+ \alpha_{k_0}) \beta_0 + f_{0a},
\]

\[
\dot{\alpha}_k = -\frac{i}{\hbar} \left[ (\epsilon_k + \Delta \chi_{k0}^2 + \epsilon_1 + \epsilon_2) \alpha_k \right]
\]

\[
+ V(\alpha_{k_0}^+ \alpha_{k_0} + 2 \beta_{k_0}^+ \beta_{k_0}) \alpha_k \]

\[
+ V(\alpha_{k_0}^+ \beta_{k_0} + \beta_{k_0}^+ \alpha_{k_0} + F_{k_0}),
\]

\[
V = V_0 \left( \frac{\alpha_{k_0}^2 + \alpha_{k_0}^3}{V_0^2} + \frac{\alpha_{k_0}^2 + \alpha_{k_0}^3}{V_0^2} \right),
\]

\[
V' = \frac{i}{2} \left[ V_0 \left( \frac{\alpha_{k_0}^2 + \alpha_{k_0}^3}{V_0^2} + \frac{\alpha_{k_0}^2 + \alpha_{k_0}^3}{V_0^2} \right) \right],
\]

where \(f_{0a} = -\left[ g_0(t) e^{i\phi_{02}} + g_0(t) e^{i\phi_{01}} \right] C_{0a} \chi_{k_0} / \sqrt{2} \) and \(f_{0b} = -\left[ g_0(t) e^{i\phi_{02}} - g_0(t) e^{i\phi_{01}} \right] C_{0b} \chi_{k_0} / \sqrt{2} \). The corresponding equations for the \(\beta \) excitations are obtained by substituting \(\alpha \) with \(\beta \).

An applied cw-optical pump with an in-plane wave vector \(k_0 \) drives the polariton polarizations \(\langle \alpha_{k_0}(t) \rangle \) and \(\langle \beta_{k_0}(t) \rangle \).

First we obtain the steady-state solution by substituting \(\langle \alpha_{k_0} \rangle = |\alpha_{k_0}| e^{i\omega_{0a}} \) and \(\langle \alpha_0 \rangle = |\alpha_0| e^{i\omega_{0a}} \), and then the solution for \(k=0 \) can be obtained in terms of the pump intensity. The solution for \(k=0 \) is

\[
\tilde{a}_0 = \frac{V \tilde{f}_{0a} \left[ h \omega - (\epsilon_0 + \Delta \chi_{00}^2 + \epsilon_1 + \epsilon_2) \right] f_{0a}}{(h \omega - E_1) (h \omega - E_2)},
\]

\[
\tilde{b}_0 = \frac{V \tilde{f}_{0b} \left[ h \omega - (\epsilon_0 + \Delta \chi_{00}^2 + \epsilon_1 + \epsilon_2) \right] f_{0b}}{(h \omega - E_1) (h \omega - E_2)},
\]

\[
N = N_{k_0} + N_{k'_0},
\]

\[
\tilde{N} = N_{k'_0} - N_{k_0},
\]

\[
E_1 = \epsilon_0 + \epsilon_1 + \epsilon_2 + \sqrt{\Delta \chi_{00}^2 + \Delta \chi_{01}^2 + \Delta \chi_{11}^2},
\]

\[
E_2 = \epsilon_0 + \epsilon_1 + \epsilon_2 - \sqrt{\Delta \chi_{00}^2 + \Delta \chi_{01}^2 + \Delta \chi_{11}^2},
\]

For a probe beam with \(\sigma_+ \) polarization,
For a probe beam with $\sigma_-$ polarization,

$$P_{01} \sim \frac{\Delta \sigma_0^2}{(\hbar \omega - E_1)(\hbar \omega - E_2)}.$$

(28)

For a probe beam with $\sigma_+$ polarization,

$$P_{01} \sim \frac{\Delta \sigma_0^2}{(\hbar \omega - E_1)(\hbar \omega - E_2)}.$$

(29)

These results indicate that the energy position of the optical response depends on the excitation intensity that is transferred from the pump. These expressions also show that there is no threshold excitation power, which is different from the case for the short delay time amplification discussed in the previous section. Therefore the optical gain here will be smaller. It is also evident that $P_{01}$ and $P_{0\bar{1}}$ are out of phase by $\pi$ for the different spin-polarization projections, which is consistent with the experiments in Ref. 5.

The experimental results in Ref. 5 also show a variation in the position of the $k=0$ peak as a function of the delay time between pump and probe. We address this issue here. Note that Eqs. (17) and (18) for the $\alpha$ excitations together with the similar equations for the $\beta$ excitations can be solved simultaneously. We obtain two independent second-order differential equations for each excitation. The equation for $\alpha$ is

$$\dot{a}_0 - \frac{2i}{\hbar} (\epsilon_0 + VN + 2i \gamma) a_0 - \frac{1}{\hbar^2} \left( (\frac{VN}{\hbar})^2 - 2i \gamma VN \right) a_0 - (\epsilon_0 + VN + \Delta X_0^2 + i \gamma) (\epsilon_0 + VN - \Delta X_0^2 + 3i \gamma) a_0 + \frac{\epsilon_0 - \Delta X_0^2 + VN + 3i \gamma}{\hbar^2} a_0 - \frac{VN}{\hbar^2} f_\alpha - \frac{VN}{\hbar^2} f_\beta$$

(31)

and similarly for $\beta_0$. Then we readily obtain for the $\alpha$ mode

$$\epsilon_0 \sim -\gamma \theta^h \left[ e^{i \theta^h} \left( \cos \frac{\theta^t}{\hbar} - i \frac{\delta^h}{\theta^h} \sin \frac{\theta^t}{\hbar} \right) + 1 \right],$$

(32)

$$\theta = \sqrt{((\Delta X_0^2)^2 + (VN)^2)},$$

(33)

$$\delta = \epsilon_0 + VN.$$

(34)

The optical gain at $k=0$ can be calculated by performing a Fourier transformation for the excitations $\alpha$ and $\beta$. In Fig. 3 the energy of the peak line position for different excitation powers is given as a function of delay time. These time-dependent results have a similar appearance to those observed experimentally.5 The experiments in Ref. 5 found that the peak position for the $\sigma_-$ probe polarization was shifted by $\pi$ in time with respect to that for $\sigma_+$ probe.

In the calculations in Fig. 3 as the pump power is increased, energy oscillations start to appear as a function of the delay time and of excitation power. These oscillations are controlled by the exchange splitting and by the excitation density through the parameter $\theta$. At low excitation the oscillation period is determined by the exchange splitting $\Delta$, which is relatively very small. Thus we do not see oscillations at low powers, and indeed no oscillations are seen experimentally. This is shown in Fig. 3(a). As the excitation power is increased, more free excitons become available, the $VN$ factor becomes dominant, and pronounced oscillations are seen [Figs. 3(b) and 3(c)]. From this work there are two implications for experiment. One is that exchange splitting accounts for the out-of-phase beats in terms of the $\sigma_+$ and $\sigma_-$ probe polarizations. The second is that nonlinear exciton-exciton interaction and phase space filling are responsible for the period of the beats as a function of excitation power.

Finally, we notice that the role of the polariton linewidth $\gamma$ in this case is expressed only through the exponential decay of the energy line splitting, whereas for the short delay time optical amplification discussed above $\gamma$ is one of the parameters that determines the threshold power.

**IV. CONCLUSION**

We have presented a microscopic treatment for polariton effects in photonic wire systems. Results are obtained for the parametric wave mixing of a high optical gain at short delay times for which energy momentum is conserved for the three macroscopically occupied states, the probe, pump, and idler states. Scattering between the lower and upper polariton subbranches, which is not possible in planar cavity, is found
here. An analytical solution for the $k=0$ excitation is obtained as a function of time. It shows that the gain depends on the excitation power, on polariton broadening, and on delay time in a similar way as in the case for a planar cavity.

In addition, we treat the case where pump-probe experiments show optical gain at larger delay times where only two macroscopically occupied states are involved—the pump state at $k_0$ and the probe state at $k=0$—and these are the only macroscopically occupied states. The $k_0$ state here is different from that for the short delay time optical amplification where an idler state was macroscopically occupied. The most interesting feature in that work is that the optical gain is spin sensitive since oscillations between the two optically active states with circular polarizations $\sim 1$ and $\sim 2$ are seen. To describe this, we include the exchange splitting between the spin split bright polaritons. In structures, such as quantum wires, where there is uniaxial strain, a splitting between polaritons with spin up and down is expected to appear. Assigning the quantum beats to different spins seems reasonable since the decay time of the beats is relatively long. This is consistent with the fact that the decoherence time of the spin degree of freedom of excitons is much longer than the exciton decoherence time.9

The microscopic equations for the coherent pump-probe processes were derived from the same model with the inclusion of an exchange interaction for the spin quantum beats experiments. We have obtained solutions for the optical gain as a function of delay time, excitation power, and polariton broadening.

ACKNOWLEDGMENTS

The authors would like to thank G. Ramon and M. Bayer for discussions. This work was supported in part by ONR and by the DARPA QuIST program. One of us (L.M.W.) acknowledges support from NRL/NRC.