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Exciton-polariton wakefields in semiconductor microcavities

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ABSTRACT

We consider the excitation of polariton wakefields due to a propagating light pulse in a semiconductor microcavity. We show that two kinds of wakes are possible, depending on the constituents fraction (either exciton or photon) of the polariton wavefunction. The nature of the wakefields (pure excitonic or polaritonic) can be controlled by changing the speed of propagation of the external pump. This process could be used as a diagnostic for the internal parameters of the microcavity.

1. Introduction

Semi-conductor microcavities, designed to increase the light-matter coupling, consist of a pair of distributed Bragg mirrors confining a electromagnetic mode and one (or several) quantum well with an exciton resonance [1,2]. In the strong coupling regime, where the exciton–photon coupling overcomes the losses, a new type of elementary excitations, called exciton polaritons (or cavity polariton), is formed [3]. Polaritons are therefore a coherent superposition of semi-conductor excitations (excitons) with light (photons).

One of the crucial aspects for the rapid development of research in the field of semi-conductor microcavities stems from the fact that polaritons may undergo Bose–Einstein condensation [4,5] – which has been experimentally observed in a series of works [6–12] –, putting together the fields of quantum optics and Bose–Einstein condensates (BEC) [13]. The most important differences with respect to the usual atomic BECs are: i) the possibility of condensation to occur at higher temperatures, as a consequence to the very small polariton mass (typically, $m \sim 10^{-5} m_e$) and ii) the fact of polariton BEC being a driven-dissipative phenomenon [14, 15], where the condensate is a steady state resulting from the pumping-loss balance. Nevertheless, it has also been shown that a thermodynamic regime for polariton condensates created by a nonresonant pump (i.e. pumped far away with respect to the condensation energy minimum) is possible, allowing for the definition of a temperature and a chemical potential for the system [16]. Another important property of polariton BECs concerns superfluidity. The phase transition expected for two-dimensional polaritons is a Berezinskii–Kosterlitz–Thouless (BKT) transition toward a superfluid state [17] and not a true BEC. Such a phase transition has not yet been observed in CdTe-based and GaN-based structures, as a consequence of a strong structural disorder leading to the formation of a glass phase [18] and to condensation in the potential minima of the disorder potential [19]. Signatures of BKT transition have nevertheless been reported in cleaner, GaAs-based samples [20]. Some interesting features related to the nonlinearity of the system, such as amplification [21] and optical bistability [22] have been investigated. Moreover, recent theoretical studies and experimental observation of topologically stable half-solitons [23–25] and half-vortices [26,27] in spinor polariton condensates have allowed the study of the dynamics and many-body properties of topological defects in the presence of external fields [28–30].

In the most usual experimental configurations, the external pump is fixed, occupying a well defined region of the planar cavity. However, if we allow the pump to move, the occurrence of new time-dependent phenomena can be expected in exciton-polaritons, even below the condensation threshold. Such an experimental configuration can be achieved with a pulsed laser, which propagation velocity can be controlled by changing the incidence angle [29]

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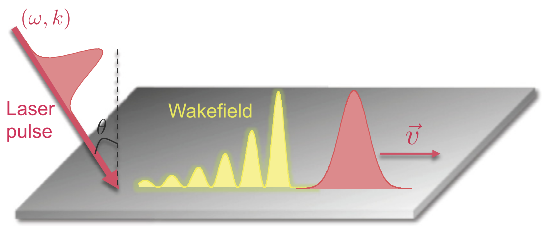

Fig. 1. (Color online.) Schematic representation of the experimental setup to produce wakefields. A laser pulse of frequency ω and momentum k is sent into the cavity (represented here by the shadowed rectangle). The velocity of the pulse is controlled with the incidence angle θ, such that $v = c\sin\theta$. The wakefield is behind the pulse.

(see Fig. 1 for a schematic representation). In this work, we study the properties wakefields in semi-conductor microcavities excited by a light pulse. Wakefields are universal phenomena which can be produced by the motion of a boat in the surface of a lake, or by a laser pulse propagating in a gas, having important technological implications in the case of laser-plasma acceleration [31]. When an intense electromagnetic pulse hits the plasma, it produces a wake of plasma oscillations through the action of a ponderomotive force. Electrons trapped in the wake can then be accelerated up to very high (relativistic) energies, providing an alternative yet efficient way of accelerating charged particles [32,33]. Acoustic wakefields produced by a Bose–Einstein condensate moving across a thermal (non-condensed) gas have also been considered [34,35]. Recently, wakefield excitation in metallic nanowires has also been investigated and pointed out as mechanism to produce energetic ultra-violet (XUV) radiaton [36]. In the present work, we show that a similar process can occur in a gas of excitons and exciton-polaritons (we should refer to the latter as "polaritons"), depending on the point of the dispersion that the system is pumped.

The structure of this paper is the following. In Section 2, we establish the basic equations of our problem. We start from the coupled photon–exciton wave equations and derive the energy spectrum. We then include the external pump term and derive the appropriate wakefield equations. In Section 3, we derive the wakefield solution for an exciton gas, by assuming that the lower polariton branch is pumped in the exciton-dominated part (high-wavevector). In Section 4, we derive the general form of polariton wakefields, produced if the pumped is tuned near the bottom of the polariton dispersion. Finally, in Section 5, we state some conclusions.

2. Basic equations

The coupled dynamics of the photonic and excitonic fields, respectively represented by $\phi(\mathbf{r},t)$ and $\chi(\mathbf{r},t)$, can be described by the following equations [1,37,38]

$$i\hbar\frac{\partial\phi}{\partial t} = -\left[\frac{\hbar^2}{2m_\phi}\nabla^2 + \frac{i\hbar}{2\tau_\phi}\right]\phi + \frac{\hbar}{2}\Omega_R\chi + P, \tag{1}$$

$$i\hbar\frac{\partial\chi}{\partial t} = -\left[\frac{\hbar^2}{2m_\chi}\nabla^2 + \frac{i\hbar}{2\tau_\chi}\right]\chi + \frac{\hbar}{2}\Omega_R\phi + \alpha_1|\chi|^2\chi, \tag{2}$$

where we have considered the special case of zero detuning between the laser and the cavity mode. Here, $m_\phi \ll m_e$ and $m_\chi \leq m_e$ are the photon and exciton masses, respectively (m_e represents the electron mass), and Ω_R is the Rabi frequency measuring the strength of the coupling. The nonlinear term $\alpha_1 = 6E_b a_B^2/\mathcal{S}$, where $\mathcal{S}$ is the normalization surface, E_b is the exciton binding energy and a_B the corresponding Borh radius, accounts for the polariton–polariton contact interactions [39]. Here, we choose the parameters according to the typical experimental conditions, $m_\phi = 5\times10^{-5}m_e$, $m_\chi = 0.4m_e$ and $\hbar\Omega_R = 8$ meV. The quantities τ_ϕ and τ_χ are the lifetimes of cavity photons and excitons, with typical values $\tau_\phi = 10$ ps and $\tau_\chi = 400$ ps.

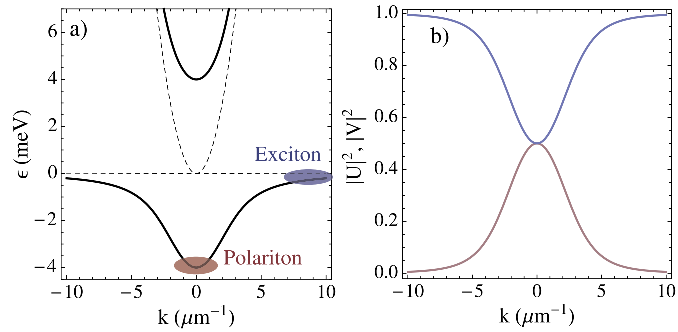

Fig. 2. (Color online.) Panel a): Lower and upper polariton branches in a semiconductor microcavity. We excite different wavevector ranges in the lower branch to generate the wake fields. Panel b): The Hopfield coefficients measuring the exciton ($|\mathcal{U}|^2$, upper curve) and photon ($|\mathcal{V}|^2$, lower curve) fractions. We have used typical parameters for a GaAs microcavity: $m_\phi = 5\times10^{-5}m_e$, $m_\chi = 0.4m_e$, and $\hbar\Omega_R = 8$ meV.

In Eq. (1), $P \equiv P(\mathbf{r},t)$ is the external pump, acting as a source of photons, which we choose to be resonantly tuned with respect to the lower polariton branch, as we specify below. We start by deriving the appropriate dispersion relations from the photon–exciton coupled field. The bare photonic and excitonic modes are readily obtained by neglecting the Rabi coupling Ω_R and the interaction term α_1,

$$i\hbar\frac{\partial\phi}{\partial t} = -\left[\frac{\hbar^2}{2m_\phi}\nabla^2 + \frac{i\hbar}{2\tau_\phi}\right]\phi, \tag{3}$$

$$i\hbar\frac{\partial\chi}{\partial t} = -\left[\frac{\hbar^2}{2m_\chi}\nabla^2 + \frac{i\hbar}{2\tau_\chi}\right]\chi, \tag{4}$$

for which we can find solutions of the form

$$\phi(\mathbf{r},t) = \phi_0\exp(i\mathbf{k}\cdot\mathbf{r} - i\omega_\phi t),$$
$$\chi(\mathbf{r},t) = \chi_0\exp(i\mathbf{k}\cdot\mathbf{r} - i\omega_\chi t). \tag{5}$$

Replacing in Eqs. (3)–(4), we obtain the dispersion relations

$$\omega_\phi = \frac{\hbar k^2}{2m_\phi} - i\gamma_\phi, \quad \omega_\chi = \frac{\hbar k^2}{2m_\chi} - i\gamma_\chi, \tag{6}$$

where the damping rates are given by $\gamma_\phi = 1/2\tau_\phi$ and $\gamma_\chi = 1/2\tau_\chi$. For a given value of the wavenumber k, we have $\omega_\phi \gg \omega_\chi$, as a result of the mass difference $m_\phi \ll m_\chi$. The polariton modes can then be obtained for a finite value of the Rabi field. Using again solutions of the form

$$(\phi,\chi)(\mathbf{r},t) = (\phi_0,\chi_0)\exp(i\mathbf{k}\cdot\mathbf{r} - i\omega t), \tag{7}$$

we can then derive the dispersion equation

$$(\omega-\omega_\phi)(\omega-\omega_\chi) = \frac{\tilde{\Omega}_R^2}{4}, \tag{8}$$

where $\tilde{\Omega}_R^2 = \Omega_R^2 - 4\gamma_\phi\gamma_\chi$ and ω_ϕ and ω_χ are determined by Eq. (6). Solving for ω, we get the two solutions $\omega_\pm$, such that

$$\omega_\pm = \frac{1}{2}\left[(\omega_\phi+\omega_\chi) \pm \sqrt{(\omega_\phi+\omega_\chi)^2 - 4\omega_\phi\omega_\chi + \tilde{\Omega}_R^2}\right]. \tag{9}$$

This corresponds to the lower (ω_-) and upper (ω_+) polariton branches, as illustrated in Fig. 2. The upper (U) and lower (L) polariton modes are given by $\psi_U = \mathcal{U}\chi + \mathcal{V}\phi$ and $\psi_L = \mathcal{V}\chi - \mathcal{U}\phi$,