Suppression of space broadening of exciton polariton transport by Bloch oscillation effect

Xudong Duan, Bingsuo Zou and Yongyou Zhang

Beijing Key Lab of Nanophotonics & Ultrafine Optoelectronic Systems and School of Physics, Beijing Institute of Technology, Beijing 100081, People’s Republic of China

E-mail: yyzhang@bit.edu.cn

Received 2 July 2015, revised 24 August 2015
Accepted for publication 3 September 2015
Published 9 October 2015

Abstract

We theoretically study the transport of exciton polaritons under different applied photon potentials. The relation between the photon potentials and the thickness of the cavity layer is first calculated by finite-element simulation. The theoretical analysis and numerical calculation indicate that the cavity photon potential is proportional to the thickness of the cavity layer with the coefficient being about 1.8 meV nm⁻¹. Further, the periodic and linear photon potentials are considered to control the transport of the exciton polaritons in weak- and strong-field pump situations. In both situations the periodic potential cannot by itself effectively suppress the scatterings of the disorder potentials of the cavity photons and excitons and the nonlinear exciton–exciton interaction. When the linear potential is added to the cavity photons, the exciton polariton transport exhibits the Bloch oscillation behavior. Importantly, the polariton Bloch oscillation can strongly suppress the space broadening of the exciton polariton transport due to the disorder potentials and nonlinear exciton–exciton interaction, which is beneficial for designing the polariton circuits.

Keywords: exciton polariton, microcavity, Bloch oscillation

1. Introduction

Studies on polaritons, especially on microcavity polaritons (MPs), have developed rapidly in recent decades, though they were first predicted by Hopfield and Agranovich in the late 1950s [1, 2]. Pictorially, MPs are a kind of photon-dressed excitation: a quantum state with part of the cavity photon and part of the quantum well (QW) exciton. The strong coupling between the QW excitons and cavity photons takes overall responsibility for the form of the MPs. For the MPs, there is a sharp dip at the center of the dispersion which implies that the effective mass of the MPs is very light, and thus it is easy to achieve the Bose–Einstein condensation [3, 4], superfluid [5], and vortex [6, 7], for the MPs. In addition, the exciton part in the MPs leads to the nonlinear polariton–polariton interaction [8] and the numerous dynamics of the MPs [9–12]. With the help of the dip dispersion and nonlinear interaction, people have achieved or designed optical amplifiers [13] and switches [14] based on the microcavities. This research has prompted a great deal of work toward the development of practical devices.

With today’s micro-/nano-fabricating technologies, physical experimentalists have been able to manufacture many kinds of semiconductor microcavities (SMCs), including photonic dots [15], microcavity arrays [16–18], and one-dimensional microcavities [19]. In these microcavities, scalar potentials can be generated for the QW excitons or and the cavity photons. For example, the polariton trap used in the Bose–Einstein condensation is induced by applying a mechanical stress [20]. In addition, one can turn to the dependence of the cavity mode energy on the microcavity structure. With this idea, position-dependent potentials for the cavity photons can be designed into arbitrary geometries, such as to deposit surface metallic strips on top of the SMC [21]. These developments are the foundation of the accurate design of the microcavity devices, such as optical diodes [22] and quantum circuits [23–25]. Generally speaking, it is easier to design the photonic potentials than to design the excitonic...
ones. When integrating the polariton devices together, the communication among them also needs an accurate design on the microcavities to suppress the MP scatterings due to the disorder potentials and the nonlinear exciton–exciton interaction.

The foundation of the communication among the polariton devices integrated together is to achieve the transport of the MPs without any broadening. This inspires us to study the influence of the applied potentials to the transport of the MPs. Because the photonic potentials are more convenient to achieve than the excitonic ones, we only consider the effects of the photonic potentials here. The photonic potentials originate from the defects of the microcavities or the geometry of the distributed Bragg reflectors (DBRs); see figure 1. In the present work, we focus on three kinds of DBR geometries, namely, planar (in figure 1(a)), linearly tilted (in figure 1(b)), and periodic (in figure 1(c)). The method used in figure 1 can introduce larger photonic potentials than the one in which optical structures are deposited on top of the SMC [21]. For the planar SMC in figure 1(a), there is no applied photonic potential except a random photonic potential due to the random defects in the microcavity. It is known that different geometries of the DBRs give different photonic potentials. Though the quantitative relation between them is very important in order for people to design the exciton polariton devices, it has not been calculated or studied in detail so far, and thus we will first determine this relation. In the present work, we will prove that the Bloch oscillation of the polariton, being similar to that of the electromagnetic wave [26, 27], can effectively suppress the space broadening of the polariton transport.

This work is organized as follows. In section 2, we first find the relation between the photonic mode and the DBR geometry, and then determine the corresponding parameters of the photonic mode. Based on these, we use the Gross–Pitaevskii equations to describe the dynamics of the MPs. In section 3, we discuss the influence of the three geometries of the DBRs on the transport of the MPs. The Bloch oscillation could be observed for certain DBR geometry and can suppress the space broadening of the exciton polariton transport. Finally, we present a brief conclusion in section 4.

2. Theoretical model

When the SMC is not planar, as shown in figures 1(b) to (c), the photon energy will depend on the transversal coordinates, which will influence the transport behavior of the MPs strongly. Without loss of generality, we assume the geometry of the DBR only depends on the coordinate \( y \); see figure 1. The starting point of the present section is to find the dependence of the Hamiltonian of the cavity photons on the geometry of the DBR.

2.1. Photonic Hamiltonian

For the planar SMC shown in figure 1(a), the Hamiltonian of the cavity photons is well known, namely,

\[
H_{ph} = -\frac{\hbar^2}{2m_c} \nabla^2 + E_0^c
\]

(1)

where \( \nabla^2 = \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} \). The values of \( E_0^c \) and \( m_c \) are the fundamental energy and effective mass of the cavity photons, respectively. They can be roughly estimated by \( E_0^c = \hbar \omega_0 c/n_c \) and \( m_c = n_c/m_0 = E_0^c/\frac{\hbar^2 k_0^2}{2c^2} \), where \( \hbar \), \( c \) and \( n_c \) are the Planck constant, vacuum light speed, and refractive index of the cavity layer, respectively, and \( k_0 = 2\pi n_c/L_c \) being the effective thickness of the cavity layer. The precondition of equation (1) is that the in-plane wave vector of the cavity photon \( k \) must be far less than \( k_0 \) which, for general microcavities, is satisfied.

When the effective thickness of the cavity is added by a small value \( \Delta L_c \), equation (1) can be cast into

\[
E^c(k) \approx (E_0^c + \Delta E_0^c) + \left(1 - \frac{\Delta E_0^c}{E_0^c}\right) \frac{\hbar^2 k^2}{2m_c}
\]

(2)

where \( \Delta E_0^c \) represents the increase value of the fundamental photon energy due to \( \Delta L_c \). For the GaAs-based SMCs, \( E_0^c \) is about 1.5 eV, and the kinetic energy of the cavity photon is about several or tens of meV. Hence, compared with the fundamental energy, the variation of the kinetic energy of the cavity photon due to the cavity thickness could be neglected. Therefore, we can rewrite the Hamiltonian of the cavity photon as follows:

\[
H_{ph} = -\frac{\hbar^2}{2m_c} \nabla^2 + V_c(x) + E_0^c
\]

(3)

where \( E_0^c = \frac{\hbar^2 k_0^2}{2m_c} \) and \( L_0 \) being a reference cavity thickness. The value of \( V_c(x) \) represents the photonic potential which could be rewritten as

\[
V_c(x) = E_0^c \left[ \frac{L_0}{L_0 + d(x)} - 1 \right] \approx -E_0^c \frac{d(x)}{L_0}
\]

(4)

where \( d(x) \) is the thickness deviation of the cavity layer from \( L_0 \) and is commonly far less than \( L_0 \). The minus sign of the last term in equation (4) indicates that the eigenfrequency of the cavity photon decreases with increasing the cavity thickness. The effectiveness of the Hamiltonian in equation (3) can be confirmed by comparing its results with the numerical simulation of the finite-element method.
We take the periodic GaAs-based SMC as an example; see figure 1(c). The refraction indexes of GaAs and AlAs are \( n_a = 3.6 \) and \( n_b = 3.0 \), respectively. The SMC is designed with the fundamental photon energy \( E_0 = 1478.8 \) meV. In addition, the potential \( V_c(x) \) of the cavity photons is taken as the following form,

\[
V_c(x) = -\alpha d(x) = -\alpha d_0 \cos \left( \frac{2\pi}{L_p} x \right)
\]

where \( d_0 \) and \( L_p \) denote the amplitude and period of the variation of the cavity thickness, and \( \alpha \) is a coefficient that describes the photon potential owing to the variation of the cavity thickness. With \( d(x) \), one can numerically calculate the photonic bands of the SMC by the finite-element method. It is natural that the band gaps at \( \pm \pi/L_p \) increase with increasing \( d_0 \), just as shown in figures 2(a) to 2(c). In figure 2, all the data denoted by the dots are calculated from the finite-element simulation, while all the solid lines are calculated from the Hamiltonian given in equation (3). In order to make the theoretical results meet the simulation data, we fit the dispersion in figure 2(a) by a parabolic dispersion (whose band gap is zero due to \( d_0 = 0 \)) and find the effective mass of the cavity photon \( m_c = 3.1 \times 10^{-3} m_e \) where \( m_e \) is the mass of a free electron. Using this effective mass of the cavity photon, we find the optimum value of \( \alpha \) is about \( \alpha = 1.8 \) meV nm\(^{-1}\) for the photon bands shown in figure 2(b). By figure 2(c), we further confirm the effectiveness of these two parameters. In practice, we also verify the availability of these two values of \( \alpha \) to the cases with other \( d_0 \). The calculations indicate that the parameter of \( \alpha = 1.8 \) meV nm\(^{-1}\) is effective enough when \( d_0 \lesssim 12 \) nm. As a result, we can expect that the Hamiltonian in equation (3) is effective enough when the deviation of the cavity thickness is small. In the present work, this assumption will always hold. Besides, we neglect the polarization difference between the transverse electrical mode and transverse magnetic mode for convenience. This means that the effective masses of the transverse electric mode and transverse magnetic mode are taken as equal to each other.

At last, we point out that the damping rate of the photon mode in the present photonic crystal, strictly speaking, depends on the form of the cavity potential \( V_c(x) \), while the corresponding dependence is very weak, due to the fact that the \( d_0 \) we consider is far less than the thickness of the cavity layer. For convenience, we will not consider the weak dependences of the photon damping rate on the cavity potential, and thus will take it as a constant in the following calculations. Our calculations from the finite-element simulation also support this simplification.

2.2. Polaritonic crystal

With the help of equation (3) the polariton Hamiltonian can be written as

\[
H_{pl} = \int d^3r \hat{\psi}^\dagger(r) \hat{\psi}(r) + \frac{1}{2} g \int d^3r \hat{\psi}^\dagger_\sigma(r) \hat{\psi}^\dagger_\tau(r) \hat{\psi}_\tau(r) \hat{\psi}_\sigma(r)
\]

where \( \hat{\psi} = (\hat{\psi}_e, \hat{\psi}_x)^T \). \( \hat{\psi}_e \) and \( \hat{\psi}_x \) are the photonic and excitonic field operators, respectively, and \( g \) represents the nonlinear exciton–exciton coupling. The Hamiltonian \( \hat{H}_0 \) is

\[
\hat{H}_0 = \begin{pmatrix}
-\frac{\hbar^2}{2m_e} \nabla^2 + V_c(x) + E_0 & \Omega_g \\
\Omega_g & -\frac{\hbar^2}{2m_e} \nabla^2 + E_0
\end{pmatrix}
\]
where $m_x$ and $E_x^0$ are the exciton mass and fundamental energy, respectively, and $\Omega_R$ is the Rabi coupling between the cavity photon and QW exciton. Note that all energy bands shown in (a) to (c) belong to the lower polariton branch. (d) Variation of the first band gap $E_g$ with the photon-exciton detuning. The SMC structure could refer to figure 1(c) and other parameters used in all panels are as follows: the upper and lower DBRs hold 24 periodic cells; the cavity thickness of the SMC is described by $d(x) = d_0 \cos(2\pi x/L_p)$ with $L_p = 1.3 \mu m$ and $d_0 = 4$ nm; and $E_x^0 = 1478.8$ meV, $m_x = 0.22 m_e$, and $\Omega_R = 2.65$ meV.

Figure 3 shows the band structure of the polaritonic crystal under several different photon-exciton detunings $\delta$. The photon-exciton detuning is defined as $\delta = E_x^0 - E_\delta^x$. In addition, the photonic potential $V_c(x)$ takes the form given in equation (5). Note that only the lower-branch polariton bands are plotted and the lowest three bands are shown as the black solid, red dashed, and green dotted curves in figures 3(a)–(c). Because of $m_x \gg m_e$, the exciton level is almost non-dispersive, and therefore, the polariton bands near the exciton level become denser and denser, displayed as the top dense black curves in figures 3(a)–(c). For easy observation, the first forbidden band is denoted as the light gray region in which the value of the band gap $E_g$ is shown. From figures 3(a)–(c), the band gap decreases as the photon-exciton detuning increases; see figure 3(d). This is due to the fact that the exciton component in the lower polariton branch increases with increasing the detuning. Seen from equation (7), the exciton potential is not introduced (which is not easy to achieve in the SMCs), thus the band gap of the polariton crystal shown in figure 3 originates from the photon component of the polaritons. When $\delta$ is much less than zero, the photon component dominates in the lowest several polariton bands, and so the polariton band gap approaches the photon component of 7.0 meV; see figure 2(c). When $\delta$ is much larger than zero, the dominant component in the lowest several polariton bands turns to be the excitons, and thus the polariton band gap turns to zero. Between these two
situations, there is a region where the polariton band gap changes fast with the detuning; see figure 3(d).

2.3. Gross–Pitaevskii equations

In order to study the transport of the exciton polariton in the SMCs, for example a polaritonic crystal, we follow a standard mean-field treatment on equation (6), so that the evolution of the field functions can be described by the coupled Gross–Pitaevskii equations [28–31]

$$i\hbar\frac{\partial}{\partial t}\psi_e = \left[\frac{-\hbar^2}{2m_e} \nabla^2 + V^D_e + V_e(x) + \delta - \frac{\gamma_e}{2}\right]\psi_e + \Omega R \psi_x + f_p(x,y)e^{-i\Delta/\hbar},$$

(8a)

$$i\hbar\frac{\partial}{\partial t}\psi_x = \left(-\frac{\hbar^2}{2m_x} \nabla^2 + V^D_x - i\frac{\gamma_x}{2}\right)\psi_x + \Omega R \psi_e + g\left|\psi_e\right|^2\psi_x$$

(8b)

where $\hbar$ is the Planck constant which, for convenience, will be set as 1 hereafter. The exciton fundamental energy $E_0^e$ is chosen as the energy reference point. After the mean-field treatment, the photon and exciton fields, namely, $\psi_e$ and $\psi_x$, have been taken as complex numbers. The values $\gamma_e$ and $\gamma_x$ represent the photon and exciton losses, respectively; $\delta$ and $\Delta$ are the photon-exciton detuning and pump energy; $f_p$ measures the pump field whose space distribution and time dependence can be controlled by the applied pump lasers; $V^D_e$ and $V^D_x$ are taken as the static disordered potentials (DPs) for the cavity photons and excitons, respectively; and $V_e(x)$ represents the photon potential determined by the thickness of the cavity layer. With the help of the Gross–Pitaevskii equations, one can study the transport of the polaritons conveniently.

For clarity, the parameters used throughout the present work are listed as follows. In the following calculations and discussions, we adopt the parameters of the typical GaAs-AlGaAs SMCs: $E_0^e = 14788.8$ meV, $\delta = -2.0$ meV, $\gamma_e = 0.02$ meV, $m_e = 3.1 \times 10^{-5}m_e$, $m_x = 0.22m_e$, $\Omega_R = 2.65$ meV, $g = 0.015$ meV $\mu$m$^2$ [15], $\alpha = 1.8$ meV nm$^{-1}$. The static disordered potentials $V^D_e$ and $V^D_x$ are randomly generated, and their random-oscillation amplitudes are 0.4 and 0.1 meV, respectively.

3. Transport of exciton polaritons

In this section, we study the transport of exciton polaritons in several different applied potentials. Besides the DPs of the QW excitons and cavity photons, we consider the two following photonic potentials,

$$V^P_e(x) = V_p \times \cos \left(\frac{2\pi x}{L_p}\right),$$

(9a)

$$V^P_x(x) = V_L \times x.$$

(9b)

Here, $V^P_e(x)$ and $V^P_x(x)$ are the periodic potential (PP) and linear potential (LP) of the cavity photons, respectively, and their strengths are set to be $V_p = 7.2$ meV and $V_L = 0.5$ meV $\mu$m$^{-1}$. The period of $V^P_e(x)$ is set to be $L_p = 1.3$ $\mu$m. These two photonic potentials correspond to $d(x) = -4 \cos \left(\frac{2\pi x}{L_{pp}}\right)$ nm and $d(x) = -0.28 x$ ($x$ takes the unit of $\mu$m), respectively. The tilted angle of the DBR due to the linear potential is only about $2.8 \times 10^{-4}$ rad. Hence, the dependence of the photon damping rate on the tilted potential is very weak, and could be taken as a constant for simplicity. In addition, the pump field takes the following form

$$f_p(x,y) = F_p e^{-\frac{x^2+y^2}{w^2}} e^{ik_y y}$$

(10)

where $F_p$, $w$, and $k_p$ represent the amplitude, width, and wave vector along the $y$-direction, respectively. In this section, $w$, and $k_p$ are taken to be $w = 6 \mu$m, and $k_p = 2 \mu$m$^{-1}$. The pump frequency is set to be the energy of the lower-branch exciton polariton when the wave vector takes $k_p = 2 \mu$m$^{-1}$, namely, $\Delta = -1.55$ meV. Thus, the pump scheme adopted in the present work is a resonant excitation.

For clearly displaying the transport of exciton polaritons, this section will be divided into two subsections according to the pump-field strength: (A) weak-field pump and (B) strong-field pump. For the former case, the nonlinear exciton–exciton interaction could be neglected, while for the latter one, it must be considered.

3.1. Weak-field pump

Figure 4 shows the transport of the exciton polaritons under different applied photonic potentials which are denoted on the lower-left corner of each panel. When only the disorder potentials $V^D_e$ and $V^D_x$ are introduced (this case will be denoted as DPs hereafter), see figures 4(a) and (d), the exciton polaritons are scattered away from the transport direction, namely, the direction of the horizontal. Therefore, the transport beam is broadened in its transport path. When a periodic potential for the cavity photons given in equation (9a) is added (this case will be denoted as DPs + PP hereafter), this behavior can also be seen, see figures 4(b) and (e), but is weakly suppressed. The beam broadening of the exciton polaritons can be further suppressed by adding the linear photon potential given in equation (9b) (see figures 4(c) and (f)) (this case will be denoted as DPs + PP + LP hereafter). Obviously, the transport of the exciton polaritons in figures 4(c) and (f) exhibits an oscillation behavior whose period is about 48.5 $\mu$m. This oscillation behavior is the well-known Bloch oscillation (BO). In order to confirm that the oscillation in figures 4(c) and (f) is the BO, we first neglect the optical pump field and disorder potentials in equation (8), and then under slow wave approximation transform [26] then
into the following form

$$i\lambda \frac{\partial \varphi_\alpha}{\partial y} = \left( -\frac{\lambda^2}{2n} \frac{\partial^2}{\partial x^2} + \mathcal{V}_\alpha(x) + \varepsilon_\alpha \right) \varphi_\alpha + \Omega_\alpha \varphi_\alpha,$$  \hspace{1cm} (11a)

$$i\lambda \frac{\partial \varphi_\alpha}{\partial y} = \left( \frac{\lambda^2}{2n} \frac{\partial^2}{\partial x^2} + \varepsilon_\alpha + g \left| \varphi_\alpha \right|^2 \right) \varphi_\alpha + \Omega_\alpha \varphi_\alpha,$$  \hspace{1cm} (11b)

where $\varphi_\alpha(x) = \psi(x)e^{-\imath k_p y + \imath \Delta t}$. Other parameters are defined as follows: $\lambda = \hbar k_p/m_\alpha$, $n = \lambda k_p$, $\Omega_\alpha = \Omega_p$, $\varepsilon_\alpha = \hbar^2 k_p^2/2m_\alpha - \Delta$, $\varepsilon_\alpha = \frac{m_\alpha}{m_\alpha} \left( \delta + \hbar^2 k_p^2/2m_\alpha - \Delta \right)$, and $\mathcal{V}_\alpha(x) = \frac{m_\alpha}{m_\alpha} V_\alpha(x)$. From equation (11), the BO period $L_B$ can be found to be $L_B = 2\pi n_\alpha/m_\alpha V_\alpha L_p = 48.3 \mu m$ [32–34], which is consistent with figures 4(c) and (f). Thus, the oscillations in figures 4(c) and (f) correspond to the Bloch oscillation.

We can further confirm this by comparing the distributions of the exciton and photon fields calculated from equations (11) and (8), respectively (see figure 5). Obviously, their field distributions hold almost the same transport patterns, except for the values and widths of the field beams. The differences of the field values and beam widths between them are due to their different excitations.

In calculating figures 5(a) and (b) through equation (11), the excitation is given by the initial fields of the photons and excitons, namely, $\varphi_\alpha(x)|_{t=0} = F_p e^{-\imath k_p y^2}$ and $\varphi_\alpha(x)|_{t=0} = 0$ with $F_p = 0.1 \text{ meV } \mu m^{-1}$. In calculating figures 5(c) and (d) through equation (8), the excitation is given by equation (10) with the same parameters as in figures 4(c) and (f). This different excitation does not influence the conclusion that the oscillations in figures 4(c) and (f), as
well as in figures 5(c) and (d), are the Bloch oscillation. The excitons and photons strongly couple together to form the exciton polaritons by the Rabi terms; see equation (8). Hence, the transport patterns of the excitons and photons are similar. When the weightings of the exciton and photon components are different in the exciton polaritons, the absolute values of the exciton and photon fields appear different, as shown in figures 4 and 5.

The BO plays an important role in achieving the long-distance polariton transport with a narrow beam width. Comparing figure 4(c) with figures 4(a) and (b), we can find that the field beam width in figures 4(c) almost does not change with increasing the transport distance. In order to describe this effect, we introduce the effective width of the beam, $W_{\text{eff}}$, as follows:

$$W_{\text{eff}} = \left[ \frac{\int (x - \bar{x})^2 |\psi|_x^2 \, dx}{\int |\psi|_x^2 \, dx} \right]^{1/2}$$

where $\bar{x} = \left( \int x |\psi|_x^2 \, dx \right) / \left( \int |\psi|_x^2 \, dx \right)$, being the center position of the field beam along the direction of $x$. We plot the effective width $W_{\text{eff}}$ as the function of the transport distance in figure 6. Due to the scattering of the disorder potentials of the cavity photons and excitons, the photon field beams become wider and wider with increasing the transport distance for the cases of DPs and DPs + PP; see the upper two curves. For them, $W_{\text{eff}}$ approximately linearly increases with increasing the transport distance, which is mainly determined by the strength of the DPs. However,
when the linear potential for the cavity photons is applied, the width of the photon field beam almost does not increase after a long transport distance and is about 2.6 μm which is less than the width of the pump field \( w = 6 \) μm. We attribute this phenomenon to the Bloch oscillation which also leads to the periodical oscillation of \( W_{\text{eff}} \), as shown by the lowest curve in figure 6. The oscillation period of \( W_{\text{eff}} \) is the same as the period of the BO, i.e. about 48.5 μm. If only the LP is applied for the cavity photons, the photon and exciton field patterns look like the trace of the free-body fall until they meet the cavity boundary (which can reflect them with solid boundary conditions). This implies that when only the LP is applied, the transport of the polaritons is not along one main direction, which is why we do not consider them in the present work. According to the above discussions, we could get that the BO effect of the polariton can be used to achieve the polariton transport without space broadening for the weak-field pump situation.

### 3.2. Strong-field pump

In this subsection, we study the transport of the polaritons with the strong-field pump. For the situations of the weak-field pump, it is known that the nonlinear exciton–exciton interaction could be neglected, while this nonlinear interaction plays important roles in the situation of the strong-field pump. It can enhance the scattering of DPs during the polariton transport, as shown in figure 7. Compared with figures 4, the field patterns become wider, especially for the cases of DPs and DPs + PP. Because of the strong nonlinear interaction, the BO patterns shown in figures 7(c) and (f) also become indistinct. However, the BO still limits the space broadening of the polariton field, which can also be confirmed by figure 8.

We also plot the effective width of the field beam, \( W_{\text{eff}} \), for the strong-field pump case in figure 8. The BO makes the polariton fields move along one main direction without space broadening, as shown by the lowest curve in figure 8. The lowest curve in figure 8 indicates that the average of \( W_{\text{eff}} \) is about 4.0 μm. This value of \( W_{\text{eff}} \) is less than the width of the pump field \( w = 6 \) μm, and does not increase with increasing with the transport distance, though it is a little larger than the situation of the weak-field pump. On the other hand, \( W_{\text{eff}} \) increases with increasing the transport distance for the cases of DPs and DPs + PP, which is similar to the weak-field pump situations. Therefore, the BO not only can shield the scattering due to the DPs but also can shield the scattering due to the nonlinear exciton–exciton interaction. According to this, the BO can be used to achieve the polariton transport without space broadening.

In fact, the exciton transports in the above two pump cases are also calculated for the photon damping rates being 0.01 meV and 0.03 meV. Though the large damping rate can decrease the transport distance of the polariton, it does not change the conclusion that the BO effect of the polariton can suppress the space broadening.

### 4. Conclusion

We have theoretically investigated the transport of exciton polaritons under different applied photon potentials in the semiconductor microcavities. First, we proved that the cavity photon potential is proportional to the thickness deviation of the cavity layer from the reference value and the corresponding ratio coefficient is about 1.8 meV nm\(^{-1}\) fitted from the results of the finite-element method. Then, we studied the polariton transport under DPs of the QW excitons and cavity photons. When the periodic and linear potentials are applied to the cavity photons, the Bloch oscillation of the polaritons can be observed. In the situation of the weak-pump field, the Bloch oscillation of the polaritons is further confirmed by the method of slow wave approximation. Subsequently, we found that the Bloch oscillation can effectively suppress the scattering of the disorder potentials and the nonlinear exciton–exciton interaction. The corresponding behavior is that the effective width of the photon and exciton fields does not increase with increasing the transport distance, either in the situations of the weak- or strong-field pumps. Thus, we conclude that the Bloch oscillation effect of the polaritons can be used to achieve the polariton transport without space broadening, which is beneficial for designing the polariton circuits.

### Acknowledgments

This work is supported by NSFC (grant no. 11304015), Beijing Higher Education Young Elite Teacher Project (grant no. YETP1228), BIT Foundation for Basic Research (grant nos. 20121842005, 20131842002).

### References

[1] Hopfield J J 1958 Theory of the contribution of excitons to the complex dielectric constant of crystals Phys. Rev. 112 1555–67
[2] Agranovich V M 1959 Zh. Eksper. Teoret. Fiz. 37 1555
[3] Kasprzak J et al 2006 Bose–Einstein condensation of exciton polaritons Nature 443 409–14
[4] Byrnes T, Kim N Y and Yamamoto Y 2014 Exciton-polariton condensates Nature Phys. 10 803–13
[5] Amo A, Lefrère J, Pigeon S, Adrados C, Ciuti C, Carusotto I, Houdré R, Giacobino E and Bramati A 2009 Superfluidity of polaritons in semiconductor microcavities Nature Phys. 5 805–10
[6] Rabo Y G 2007 Half vortices in exciton polariton condensates Phys. Rev. Lett. 99 106401
[7] Lagoudakis K G, Ostatnický T, Kavokin A V, Rabo Y G, André R and Deveaud-Plédran B 2009 Observation of half-quantum vortices in an exciton-polariton condensate Science 326 974–6
[8] Zoubi H and la Rocca G C 2007 Microscopic theory of nonlinear polariton interactions in strongly coupled hybrid organic-inorganic microcavities Phys. Rev. B 76 035325
[9] Winkler K, Schneider C, Fischer J, Rahimi-Iman A, Amthor M, Forchel A, Reitzenstein S, Höfling S and Kamp M 2013 Electroluminescence from spatially confined exciton polaritons in a textured microcavity Appl. Phys. Lett. 102 041101
[10] Nakayama M, Murakami K, Furukawa Y and Kim D G 2014 Photoluminescence characteristics of polariton condensation in a CuBr microcavity Appl. Phys. Lett. 105 021903
[11] Yulin A V, Egorov O A, Lederer F and Skryabin D V 2008 Dark polariton solitons in semiconductor microcavities Phys. Rev. A 78 061801
[12] Karr J P, Baas A and Giacobino E 2004 Twin polaritons in semiconductor microcavities Phys. Rev. A 69 063807
[13] Huyhn A, Tignon J, Larsson O, Roussignol P, Delalande C, André R, Romestain R and Dang L S 2003 Polariton parametric amplifier pump dynamics in the coherent regime Phys. Rev. Lett. 90 106401
[14] Zhang Y, Jin G and Ma Y 2007 Phase effects on the exciton polariton amplifier Appl. Phys. Lett. 91 191112
[15] Verger A, Ciuti C and Carusotto I 2006 Polariton quantum blockade in a photonic dot Phys. Rev. B 73 193306
[16] Park S-J, Chen K-F, Ostrom N P and Eden J G 2005 40000 pixel arrays of ac-excited silicon microcavity plasma devices Appl. Phys. Lett. 86 111501
[17] Masumoto N, Kim N Y, Byrnes T, Kusudo K, Löfler A, Höfling S, Forchel A and Yamamoto Y 2012 Exciton-polariton condensates with flat bands in a two-dimensional kagome lattice New J. Phys. 14 065002
[18] Kim N Y, Kusudo K, Löfler A, Höfling S, Forchel A and Yamamoto Y 2013 Exciton-polariton condensates near the dirac point in a triangular lattice New J. Phys. 15 035032
[19] Gamouras A, Britton M, Khairy M M, Mathew R, Dalacu D, Poole P, Poitras D, Williams R L and Hall K C 2013 Energy-selective optical excitation and detection in InAs/InP quantum dot ensembles using a one-dimensional optical microcavity Appl. Phys. Lett. 103 253109
[20] Balili R, Hartwell V, Snoke D, Pfeiffer L and West K 2007 Bose–Einstein condensation of microcarrier polaritons in a trap Science 316 1007–10
[21] Lai C W et al 2007 Coherent zero-state and π-state in an exciton-polariton condensate array Nature 450 529–32
[22] Espinosa-Ortega T, Liew T C H and Shelykh I A 2013 Optical diode based on exciton-polaritons Appl. Phys. Lett. 103 191110
[23] Probst S, Kukharchyk N, Rotzinger H, Tkalčec A, Wünsch S, Wieck A D, Siegel M, Ustinov A V and Bushev P A 2014 Hybrid quantum circuit with implanted erbium ions Appl. Phys. Lett. 105 162404
[24] Liew T C H, Kavokin A V and Shelykh I A 2008 Optical circuits based on polariton neurons in semiconductor microcavities Phys. Rev. Lett. 101 016402
[25] Espinosa-Ortega T and Liew T C H 2013 Complete architecture of integrated photonic circuits based on and not logic gates of exciton polaritons in semiconductor microcavities Phys. Rev. B 87 195305
[26] Longhi S 2007 Topological optical Bloch oscillations in a deformed slab waveguide Opt. Lett. 32 2647–9
[27] Longhi S 2012 Bloch–Zener oscillations of strongly correlated electrons Phys. Rev. B 86 075144
[28] Zhang Y, Jin G and Ma Y-Q 2009 Boundary effects on the dynamics of exciton polaritons in semiconductor microcavities J. Appl. Phys. 105 033105
[29] Leyder C, Liew T C H, Kavokin A V, Shelykh I A, Romanelli M, Karr J P, Giacobino E and Bramati A 2007 Interference of coherent polariton beams in microcavities: polarization-controlled optical gates Phys. Rev. Lett. 99 196402
[30] Shelykh I A, Rubo Y G, Malpuech G, Solnyshkov D D and Kavokin A 2006 Polarization and propagation of polariton condensates Phys. Rev. Lett. 97 066402
[31] Flayac H and Savenko I G 2013 An exciton-polariton mediated all-optical router Appl. Phys. Lett. 103 201105
[32] Dreisow F, Szameit A, Heinrich M, Pertsch T, Nolte S, Tünnemann A and Longhi S 2009 Bloch–Zener oscillations in binary superlattices Phys. Rev. Lett. 102 076802
[33] Longhi S 2006 Optical Zener–Bloch oscillations in binary waveguide arrays Europhys. Lett. 76 416
[34] Zhang H R and Sun C P 2010 Bloch oscillations of polaritons of an atomic ensemble in magnetic fields Phys. Rev. A 81 063427