Localization-Delocalization Transition of Indirect Excitons in Lateral Electrostatic Lattices

M. Remeika, J.C. Graves, A.T. Hammack, A.D. Meyertholen, M.M. Fogler, and L.V. Butov
Department of Physics, University of California at San Diego, La Jolla, CA 92093-0319

M. Hanson and A.C. Gossard
Materials Department, University of California at Santa Barbara, Santa Barbara, California 93106-5050
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We study transport of indirect excitons in GaAs/AlGaAs coupled quantum wells in linear lattices created by laterally modulated gate voltage. The localization-delocalization transition (LDT) for transport across the lattice was observed with reducing lattice amplitude or increasing exciton density. The exciton interaction energy at the transition is close to the lattice amplitude. These results are consistent with the model, which attributes the LDT to the interaction-induced percolation of the exciton gas through the external potential. We also discuss applications of the lattice potentials for estimating the strength of disorder and exciton interaction.

Transport of particles in periodic potentials is a basic problem, which concerns a variety of systems extending from condensed matter systems with electrons in ionic lattices to engineered systems such as photons in photonic crystals and cold atoms in optical lattices. The particle localization and LDT are perhaps the most interesting transport phenomena. Particular cases of the latter—the metal-insulator and superfluid-insulator transitions—have been extensively studied for electrons, photons, and cold atoms in lattices [1, 2, 3, 4, 5, 6].

Tunability of system parameters has been essential in studies of cold atoms in optical lattices, allowing to probe localization of atoms with increasing lattice amplitude [3, 4, 5, 6]. Here, we study a condensed matter system with tunable parameters: excitons in electrostatic lattices created by a gate voltage. In this system parameters of both the lattice, e.g., the lattice amplitude, and the particles, e.g., the exciton density, can be controlled.

An indirect exciton in coupled quantum wells (CQW) is a bound state of an electron and a hole in separate wells (Fig. 1a). Lifetimes of indirect excitons exceed that of regular excitons by orders of magnitude and they can travel over large distances before recombination [7, 8, 9, 10, 11, 12]. Also, due to their long lifetime, these bosonic particles can cool to temperatures well below the quantum degeneracy temperature $T_{\text{DB}} = 2\pi \hbar^2 n/(mgk_B)$ [13]. (In the studied CQW, excitons have the mass $m = 0.22m_0$, spin degeneracy $g = 4$, and $T_{\text{DB}} \approx 3$ K for the density per spin $n/\hbar = 10^{10}$ cm$^{-2}$). Furthermore, indirect excitons in CQW have a dipole moment $ed$, where $d$ is close to the distance between the QW centers. This allows imposing external potentials $E(x, y) = edF_z(x, y) \propto V(x, y)$ for excitons using a laterally modulated gate voltage $V(x, y)$, which creates a transverse electric field $F_z(x, y)$.

The lattice potential for indirect excitons $E(x)$ was created by interdigitated gates. Base voltage $V_0 = 3$ V realized the indirect regime where indirect excitons are lower in energy than direct excitons, while voltage modulation $\Delta V$ controlled the lattice amplitude (Fig. 1b,c).

Note that in-plane electric field $F_z$ present near electrode edges can lead to exciton dissociation [14]. Therefore the CQW layers in our structure were positioned closer to the homogeneous bottom electrode. This design suppresses $F_z$, making the field-induced dissociation negligible [17]. An example of the calculated [17] (unscreened) $E(x)$ is shown in Fig. 1(c). Zero energy corresponds to zero voltage, the 4 meV energy shift due to binding energy of the indirect exciton is not shown. The potential...
modulation is nearly sinusoidal \( E(x) = E_0 \cos^2 (px/2) \). Its amplitude is \( E_0 = 3 \text{ meV} \) for \( \Delta V = 1 \text{ V} \) and scales linearly with \( \Delta V \). The lattice period \( 2\pi/q = 2 \mu m \) is determined by the electrode dimensions.

CQW structure was grown by MBE. \( n^+ \)-GaAs layer with \( n_{33} = 10^{10} \text{ cm}^{-3} \) serves as a homogeneous bottom electrode. Semitransparent top electrodes were fabricated by magnetron sputtering a 90 nm indium tin oxide layer. CQW with 8 nm GaAs QWs separated by a 4 nm \( \text{Al}_{0.33}\text{Ga}_{0.67}\text{As} \) barrier were positioned 100 nm above the \( n^+ \)-GaAs layer within an undoped 1 \( \mu \text{m} \) thick \( \text{Al}_{0.33}\text{Ga}_{0.67}\text{As} \) layer. Excitons were photogenerated by a 633 nm HeNe or 786 nm Ti:Sapphire laser focused to a spot \( \sim 10 \mu \text{m} \) in diameter in the center of the \( 150 \times 150 \mu \text{m} \) lattice. Exciton density was controlled by the excitation power. Photoluminescence (PL) images of the exciton cloud were captured by a CCD with a filter 800 \( \pm 5 \) nm covering the spectral range of the indirect excitons. The diffraction limited spatial resolution was 1.5 \( \mu \text{m} \) (N.A. \( = 0.28 \)). The spectra were measured using a spectrometer with resolution 0.18 meV.

Figure 1 shows images of the exciton cloud at zero \( \Delta V = E_0 = 0 \) and finite \( \Delta V = 1.2 \text{ V} \), \( E_0 = 3.7 \text{ meV} \) lattice amplitude for different excitation powers \( P \). At low \( P \) (Fig. 1d,g), the cloud profile essentially coincides with the laser excitation spot. This indicates that excitons do not travel beyond the excitation spot, i.e., they are localized. On the contrary, at high \( P \) (Fig. 1e,f,h,i), the excitons spread beyond the excitation spot indicating that they are delocalized. The LDT occurs both with and without the lattice. While the cloud is practically symmetric at \( \Delta V = 0 \) (Fig. 1h,i), at finite \( \Delta V \) it is compressed in the lattice direction (Fig. 1e,f).

The lattice potential also causes periodic modulations of PL characteristics. Figure 2 presents a PL image in energy–\( x \) coordinates in the delocalized regime. Both the integrated PL intensity \( I(x) \) and the average PL energy \( \hbar \omega(x) \) show small modulation at the lattice period superimposed on a smoothly varying profile (Fig. 2b,c). To demonstrate the modulations more clearly, we subtracted the smooth component and plot the remainder on a magnified scale in Fig. 2b,c. Minima in energy correspond to the maxima in intensity. We define the amplitude of energy modulation as the difference between adjacent maxima and minima \( \Delta \omega = \omega_{\text{max}} - \omega_{\text{min}} \). Figure 2 shows that \( \hbar \Delta \omega \) is much smaller than the lattice amplitude \( E_0 = 3.7 \text{ meV} \) that is discussed below. No intensity or energy modulation was observed at \( \Delta V = 0 \). This indicates that the modulations in question are not due to the partial light absorption in the top electrodes.

The PL intensity has a maximum along a ring in the regime of delocalized excitons (Figs. 1e,f,h,i, 2c). This so-called inner ring was previously observed in PL patterns of indirect excitons without lattices [9]. It was explained in terms of exciton transport and cooling [3,11]. The inner ring effect persists in the lattice (Figs. 1e,f, 2c).

The full width at half maximum of the exciton cloud in the \( x \)-direction is plotted in Fig. 3a. Initially, it is practically independent of \( P \) but then starts to grow as \( P \) increases. We define the excitation power at the transition \( P_{\text{LDT}} \) as the point where the extrapolation of this growth to small \( P \) becomes equal to the low-\( P \) constant. At the LDT, the exciton cloud starts to spread beyond the excitation spot and the cloud extension changes from constant to increasing with \( P \). Figure 3a shows that the transition is smooth. This yields \( P_{\text{LDT}} \approx 2 \text{ \mu W} \) for low lattice amplitudes \( E_0 \leq 1 \text{ meV} \). At higher \( E_0 \) the LDT for the \( x \)-direction shifts to higher excitation powers with increasing lattice amplitude (Fig. 3a,c). The LDT was also observed with reducing lattice amplitude, see the data for \( P = 5 \text{ \mu W} \) (Fig. 3b). Note that excitons remain localized at lower \( P = 0.9 \text{ \mu W} \) and delocalized at higher \( P = 21 \text{ \mu W} \) for all \( E_0 \) in Fig. 3b.

The exciton transport along the \( y \)-direction is only weakly effected by the lattice (Fig. 1d-i). The LDT for this direction shifts slightly to lower excitation powers with increasing \( E_0 \) (Fig. 3c).

The smooth component of the \( \hbar \omega(x) \) also exhibits an interesting behavior. It increases with increasing exciton density, both with increasing \( P \) or reducing \( |x| \). Let
the middle curve corresponds to the percolation point. It also explains why $P_{LDT}$ increases as $E_0$ goes up, see Fig. 3d.

Adding disorder does not modify this picture greatly as long as $E_0$ remains larger than the characteristic amplitude of $E_{\text{rand}}$. Otherwise, the percolation is determined by the random potential [21, 22], so that the dependence of $P_{LDT}$ and $\Delta\omega_{LDT}$ on $E_0$ saturates. The saturation point gives the estimate of $E_{\text{rand}}$. From Fig. 3d, we find $E_{\text{rand}} \sim 0.8$ meV. This number is comparable to the PL linewidth at low densities, suggesting that the disorder is responsible for both of these energy scales.

To further develop this model, we make the following simplifying assumptions: 1) $E_{\text{rand}} \equiv 0$, while $E(x)$ can be considered slowly varying, 2) excitons reach a quasi-equilibrium state with local chemical potential $\mu(x)$ and temperature $T(x)$, which are also slowly varying, 3) exciton interaction is local (dipolar tails, see below, are neglected). Under these assumption, we obtain

\begin{equation}
\tilde{\zeta}(x) \equiv \mu(x) + E(x) = \text{const},
\end{equation}

\begin{equation}
n(x) = \int_{0}^{\infty} \frac{e^{\nu_1 de}}{\exp[(e + \Re \Sigma(e, x) - \mu(x)) / k_B T(x)]} - 1,
\end{equation}

where $\zeta$ is the electrochemical potential, $\nu_1 = m/(2\pi \hbar^2)$ is the density of states per spin species, and $\Sigma(e)$ is the self-energy (in the uniform state of the same $n$).

To find the equation for the PL energy shift $\hbar \Delta \omega$ we take advantage of the smallness of $Q$, the range of in-
interaction. At large r the interaction is known to be dominated by dipole repulsion $\gamma r^2/(4\pi \varepsilon_0 r^2)$. When excitons approach each other, the interaction potential becomes complicated. What appears to be certain is that for $d = 12 \, \text{nm}$ in our experiment the exciton potential remains strictly repulsive [23, 24], and so Re $\Sigma(0)$ increases with density: Re $\Sigma(0) = n$. The growth of Re $\Sigma(0)$ with $n$ implies an increase of the PL energy $\hbar \omega$, which is observed experimentally. The calculation of Re $\Sigma(e) \approx \text{Re} \Sigma(0)$, which is reasonable for short-range interactions. Substituting $\Sigma = tn$ into Eq. (3), after some algebra we get

$$\exp \left( -\frac{n}{k_B T N_{\text{v}}(1)} \right) + \exp \left( \frac{\mu n}{k_B T v_1} \right) = 1, \quad \mu = \zeta - E(x),$$

where the dimensionless parameter $\gamma = \tau v_1$ characterizes the strength of the interaction. Given $\zeta, T$, and $\gamma$, Eq. (4) can be solved numerically for each $E(x)$. The results for $n(x)$ and $\Delta \omega(x)$ are shown in Fig. 4a,b.

The exciton interaction results in screening of the lattice potential at points where local density is not small. Because of this screening, the amplitude of the PL energy modulation $\hbar \Delta \omega$ is much smaller than $E_0$. Consider points $|x| \approx 10 \, \text{nm}$, which are midway between the center and the edge of the exciton cloud in Fig. 2, on the aforementioned inner ring. (Here the exciton temperature is close to the sample temperature $T = 1.6 \, \text{K}$ [11].) From Fig. 2b we see that $\hbar \Delta \omega(x) \approx 0.07 \, \text{meV}$, more than an order of magnitude smaller than $E_0 = 3.7 \, \text{meV}$.

Using the above equations we calculated $\Delta \omega$ as a function of the adjustable parameter $\gamma$, cf. the inset in Fig. 4. The experimental value of $\hbar \Delta \omega$ gives a rough estimate $\gamma = 2.3$. In comparison, the mean-field Hartree approximation [25, 26, 27, 28] yields the so-called “plate capacitor” formula $\gamma_{\text{cap}} = (2d/a_0)(m_{\text{e}}/m_{\text{n}}) \approx 7$, where $a_0 = 10 \, \text{nm}$ is the electron Bohr radius, see Fig. 4. The reduction of the interaction constant compared to the $\gamma_{\text{cap}}$ can be due to correlation effects [23, 28]. A systematic analysis of $\gamma$ remains a problem for future research.

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[1] P.A. Lee, T.V. Ramakrishnan, Rev. Mod. Phys. 57, 287 (1985); D. Belitz, T. R. Kirkpatrick, Rev. Mod. Phys. 66, 261 (1994).
[2] D.S. Wiersma, P. Bartolini, A. Lagendijk, R. Righini, Nature 390, 671 (1997).
[3] M. Greiner, O. Mandel, T. Esslinger, T.W. Hansch, I. Bloch, Nature 415, 39 (2002).
[4] C.D. Fertig, K.M. OHara, J.H. Huckans, S.L. Rolston, W.D. Phillips, J.V. Porto, Phys. Rev. Lett. 94, 120403 (2005).
[5] J.K. Chin, D.E. Miller, Y. Liu, C. Stan, W. Setiawan, C. Sanner, K. Xu, W. Ketterle, Nature 443, 961 (2006).
[6] G. Roati, C. D’Errico, L. Fallani, M. Fattori, C. Fort, M. Zaccanti, G. Modugno, M. Modugno, M. Inguscio, Nature 453, 895 (2008).
[7] M. Hagn, A. Zrenner, G. Böhm, G. Weimann, Appl. Phys. Lett. 67, 232 (1995).
[8] A. V. Larionov, V.B. Timofeev, J. Hvam, K. Soerensen, Sov. Phys. JETP 90, 1093 (2000).
[9] L.V. Butov, A.C. Gossard, D.S. Chemla, Nature 418, 751 (2002).
[10] Z. Vörös, R. Balili, D.W. Snoke, L. Pfeiffer, K. West, Phys. Rev. Lett. 94, 226401 (2005).
[11] A.L. Ivanov, L.E. Smallwood, A.T. Hammack, Sen Yang, L.V. Butov, A.C. Gossard, Europhys. Lett. 73, 920 (2006).
[12] A. Gartner, A.W. Holleitner, J.P. Kotthaus, D. Schuh, Appl. Phys. Lett. 89, 052108 (2006).
[13] L.V. Butov, A.L. Ivanov, A. Imamoglu, P.B. Littlewood, A.A. Shashkin, V.T. Dolgopolov, K.L. Campman, and A.C. Gossard, Phys. Rev. Lett. 86, 5608 (2001).
[14] S. Zimmermann, A.O. Gorov, W. Hansen, J.P. Kotthaus, M. Bichler, W. Wegscheider, Phys. Rev. B 56, 13414 (1997).
[15] T. Huber, A. Zrenner, W. Wegscheider, M. Bichler, Phys. Stat. Sol. (a) 166, R5 (1998).
[16] J. Krauß, J.P. Kotthaus, A. Wixforth, M. Hanson, D.C. Driscoll, A.C. Gossard, D. Schuh, M. Bichler, Appl. Phys. Lett. 85, 5830 (2004).
[17] A.T. Hammack, N.A. Gippius, Sen Yang, G.O. Andreev, L.V. Butov, M. Hanson, A.C. Gossard, J. Appl. Phys. 99, 066104 (2006).
[18] G. Chen, R. Rapaport, L.N. Pfieifer, K. West, P.M. Platzman, S. Simon, Z. Vörös, D. Snoke, Phys. Rev. B 74, 045309 (2006).
[19] A.A. High, A.T. Hammack, L.V. Butov, M. Hanson, A.C. Gossard, Opt. Lett. 32, 2466 (2007).
[20] A.A. High, E.E. Novitskaya, L.V. Butov, M. Hanson, A.C. Gossard, Science 321, 229 (2008).
[21] P. Lukan, D. Clément, P. Bouyer, A. Aspect, M. Lewenstein, L. Sanchez-Palencia, Phys. Rev. Lett. 98, 170403 (2007); T. Nattermann, V. L. Pokrovsky, Phys. Rev. Lett. 100, 060402 (2008); B. I. Shklovskii, Semicond.(St. Petersburg) 42, 909 (2008).
[22] Y.P. Chen, J. Hitchcock, D. Dries, M. Junker, C. Welford, R.G. Hulet, Phys. Rev. A 77, 033632 (2008).
[23] C. Schindler, R. Zimmermann, Phys. Rev. B 78, 045313 (2008).
[24] A. D. Meyertholen, M. M. Fogler, Phys. Rev. B 78, 235307 (2008).
[25] D. Yoshioka, A.H. MacDonald, J. Phys. Soc. Jpn. 59, 4211 (1990).
[26] X. Zhu, P.B. Littlewood, M. Hybertsen, T. Rice, Phys. Rev. Lett. 74, 1633 (1995).
[27] Yu.E. Lozovik, O.L. Berman, JETP Lett. 64, 573 (1996).
[28] S. Ben-Tabou de-Leon, B. Laikhtman, Phys. Rev. B 63, 125306 (2001).
[29] A.L. Ivanov, Europhys. Lett. 59, 586 (2002).