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Strain-Induced Darkening of Trapped Excitons in Coupled Quantum Wells at Low Temperature

N. W. Sinclair,^{*} J. K. Wuenschell, Z. Vörös,[†] B. Nelsen, and D.W. Snoke
Department of Physics and Astronomy, University of Pittsburgh, Pittsburgh, PA 15260, USA

M. H. Szymanska[‡]
University of Warwick, Coventry CV4 7AL, UK

A. Chin
Universität Ulm, 89069 Ulm, Germany

J. Keeling
University of St Andrews, St Andrews, Fife KY16 9AJ, Scotland

L. N. Pfeiffer and K.W. West
Department of Electrical Engineering, Princeton University, Princeton, New Jersey 08544, USA
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In GaAs/AlGaAs coupled quantum wells, strain-induced traps may be used to confine excitons in in-plane, harmonic traps. Using these traps, we have pursued Bose-Einstein condensation (BEC) of long-lived, spatially-indirect excitons. Here, we report a remarkable transition of the indirect exciton luminescence pattern with increasing strain, increasing exciton density, and decreasing temperature, to a spatial pattern exhibiting a large dark spot at trap center, where we expect the exciton density to be maximum. The mechanism of particle loss is ruled out as an explanation for this dark spot. While the onset criteria are approximately consistent with the conditions for BEC of a weakly interacting gas, the conspicuous proximity in energy of the indirect light-hole states suggests that an explanation employing the single-particle physics of light-hole/heavy-hole mixing may explain the phenomenon. The effect of the strain is modeled, and the resulting landscape of indirect exciton spin states is discussed. The relative oscillator strengths of these states are predicted by an exact numerical solution of the two-particle Schrodinger equation for electrons and holes in coupled quantum wells and an electric field. The contrast in oscillator strengths is sufficient to produce this luminescence pattern, but this analysis suggests a strongly diminished lifetime as stress is increased. The opposite lifetime dependence is observed experimentally. Additionally, the temperature dependence eludes explanation by this mechanism.

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I. INTRODUCTION

The remarkable realization of nearly ideal conditions in experiments on atomic Bose-Einstein condensates and degenerate gases comes at a cost in terms of apparatus complexity and overall price. In contrast, the pursuit of degenerate semiconductor systems, at more accessible temperatures and with less-elaborate trapping apparatus, represents an effort to bridge the experimental/theoretical divide with experiments that are more routine. Although these systems come with an added burden of untangling the desired many-body physics from features of the more complicated single-body states, the tunable single-particle physics of semiconductor heterostructures is potentially a desirable feature of the system, e.g. giving rise to novel and possibly more useful condensates.

While perhaps the most obvious example of uniquely solid-state degenerate behavior is the case of BCS superconductivity arising from phonon-mediated pairing, BEC-like behavior has been seen recently in exciton-

polaritons^{1,2}. Because excitons are electron-hole pairs, they may recombine to produce a photon. The strong coupling limit of quantum well excitons and photons in the microcavity structures produces a system of weakly-interacting, extremely low mass bosons called polaritons. The short polariton lifetime offers the ability to study the dynamics of non-equilibrium degenerate gases weakly coupled to a thermal phonon bath³. Along with displaying many typical weakly-interacting BEC features, e.g. a phonon-like Bogoliubov mode⁴, spatial and k-space narrowing², and long-range coherence⁵, the semiconductor microcavity polariton system also demonstrates less typical features, such as quantized half-vortices^{6,7} and spontaneous linear polarization of angular momentum^{2,8}.

In the opposite limit of photon coupling, the system of spatially-indirect excitons (IXs) in a pair of coupled quantum wells (CQWs) interacts only very weakly with photons⁹⁻¹¹. Indirect excitons, also often referred to as interwell excitons, are comprised of conduction-band electrons and valence-band holes forced by an electric field to reside in adjacent wells, separated spatially in

FIG. 1: (a) The stressor mechanism. Strain is produced by applying a localized force to the substrate side of the wafer. The optical side remains free to flex under the force of the pin, allowing expansion of the unit cell and, hence, lowering of the carrier energies. (b) A typical stressed spectrum of recombination PL after laser excitation. In order of decreasing energy, we see the heavy-hole DX, heavy-hole biexciton, and heavy-hole IX.

the wafer growth direction. Though separated, these electron-hole pairs remain strongly bound at temperatures below the energy of their Coulomb attraction. In the case of the 14 nm GaAs wells presented here, this binding energy is about 3 meV at a typical field strength of 1 meV/nm¹¹. While direct excitons (DXs), bound states of e-h pairs in a single well, quickly undergo electron-hole recombination with a lifetime of about 1 ns or lower, the small wavefunction overlap of the IX constituents strongly suppresses recombination, yielding particle lifetimes reaching 10 μ s or greater.

Because they are aligned dipoles, the excitons in this system have an intermediately strong interparticle interaction: much weaker than liquid helium but much stronger than polaritons^{12,13} and other particles with short-range interaction. The research toward excitonic condensation in this system cannot yet claim such unambiguous success as that of polariton condensation¹⁴. However, the exceptionally long lifetime of these particles continues to offer the promise of well-thermalized condensates of interacting particles.

In previous reports, it was demonstrated that applying localized force in the growth direction to a wafer creates a harmonic, in-plane trap for indirect excitons (IXs) in GaAs/AlGaAs coupled quantum wells¹⁵. In addition, it was shown that spatial equilibrium and thermal equilibrium are easily achieved within the lifetime of these particles^{17,18}.

Fig. 1(a) illustrates the stressor apparatus. A pin pressing against the GaAs substrate, under conditions in which the sample can flex on the opposing side, creates a predominantly hydrostatic expansion of the lattice on the opposite side, lowering carrier energies. At mod-

erate stress, this system has proven useful in probing the character of the IX-IX interaction¹⁹, but so far it has failed to yield any of the expected signals of quantum degeneracy²⁰.

Fig. 1(b) displays a typical moderate-stress spectral profile of the photoluminescence (PL) from electron-hole recombination. The pair of higher energy lines correspond to the heavy-hole DX and heavy-hole biexciton emission, while the lower line shows the energy profile of the trapped hh IXs. Biexcitons are not expected for IX states because of their repulsive interaction¹².

The achievable densities are limited by the strong repulsive interaction between these aligned-dipolar particles; a large spectral blue shift with density tends to flatten the trapping potential, requiring deeper traps to achieve higher densities. The flatter energy profile of the IX states in Fig. 1(b) shows this blue shift, and the lack of this effect in the DX emission indicates the difference in character of the DX-IX and IX-IX interactions. One approach to overcome the blue shift problem is to lower the temperature, allowing lower particle densities for quantum degenerate effects, but going to low temperature is problematic in this system because the excitons can become trapped in local minima created by disorder¹⁷. Instead, we have aimed to increase the applied stress, which increases both the trap depth and the spring constant of the trapping potential.

In this report, we demonstrate a remarkable transformation of the IX recombination PL profile upon increasing the stress beyond that of previous reports. In Section II, we outline the experimental apparatus and, in Section III, demonstrate an unexpected dimming of the PL at trap center, where we still expect the highest IX density to occur. We show that this PL dark spot pattern may occur either as strain is increased, as the temperature is decreased, or the exciton density is increased. Section IV discusses how strain causes mixing of the light-hole (lh) and heavy hole (hh) exciton states with stress and how this mixing produces a spatially-varying oscillator strength of the lowest IX energy level. We use the Pikus-Bir Hamiltonian, along with a piezoelectric modification of the local electric field, to interpret the effect of the strain predicted by a linear elastic, finite-element analysis. This analysis predicts both DX and IX energy levels which match well with the spectral profiles gathered experimentally. The oscillator strengths of the lh and hh IX states as a function of electric field are predicted by an exact numerical solution of the two-particle Schrodinger equation for Coulomb-interacting electrons and holes in an external potential of the CQW and applied electric field, within the anisotropic effective mass approximation.

We use the predicted oscillator strengths for the lh IX and hh IX, along with the lh/hh mixing predictions, to determine the spatially-varying oscillator strength for recombination in the lowest energy state. We find that this does predict a dark spot emerging in the center of the luminescence pattern. However, neither the

experimentally-observed increase in population lifetime with increasing stress nor the pattern's onset with decreasing temperature are predicted by this modeling. In addition, the contrast between the bright/dark regions is not as large in the lh/hh mixing simulations as in experimental images. These discrepancies suggest another factor must be at work in concert with the strain-induced mixing.

In Section V we show that particle loss in the center of the trap cannot explain the results, and we briefly examine the case for Bose-Einstein Condensation underlying this pattern formation. We find agreement in the onset criteria for BEC and for this pattern formation, yet an explanation relying solely on condensation of dark ($J=2$) excitons proves insufficient to cause the pattern formation.

Despite the remaining unknowns, the modeling presented here provides insight into the spin-structure of the strained states. We expect that the knowledge of the spin states gleaned from these experiments will motivate the direction of future work aimed at Bose-Einstein condensation. The landscape of spin-states resulting from stresses even higher than that required to produce the dark spot effect provides fertile ground for work on condensation phenomenon. Specifically, we note that the lowest-lying energy state at very high stress becomes the lh IX state and exhibits much sharper trap profiles and boasts an as-yet-unexplained increased population lifetime. In addition, we expect that, if degenerate phenomena are playing any substantial role in the deviations we show here from the lh/hh mixing analysis, this degeneracy will likely carry over to the high stress lh-dominated regime.

II. EXPERIMENTAL METHOD

The coupled quantum well structures are grown by MBE on a n-type GaAs substrate and contain a symmetric pair of GaAs wells separated by a 4 nm $\text{Al}_{0.45}\text{Ga}_{0.55}\text{As}$ barrier. The width of the GaAs wells varies from 8 nm to 14 nm in different samples, which affects the ground-state energy and lifetime of the indirect excitons. The wells are sandwiched between symmetric superlattice structures, which act to minimize the current through the wells. The conduction and valence bands of the structures in the presence of electric field normal to the wells are shown in Fig. 2.

A p-type GaAs layer caps the structure. Together with the n-type opposing face, this diode structure provides a built-in electric field in the wells. The total electric field derives from both this built-in field and from an applied voltage across the wafer. The spring-loaded conducting plates that serve to hold the sample also provide electrical contact to each of its sides.

The electric field in the wells tunes the indirect exciton energy below that of the direct exciton and dramatically increases the lifetime by pulling the constituent elec-

FIG. 2: Conduction and valence band in the evaporated structure. Solid lines show the energy of states with no voltage across the structure. Dotted lines show the effect of an applied electric field. The regions of the 125 μm -thick n-GaAs substrate and p-GaAs cap are not pictured.

trons and holes apart. The strength of this field varies from experiment to experiment, depending on our desired energy/lifetime, but it is typically in the range of 1 mV/nm (10 kV/cm). The time-averaged current through the samples, which have area of around 25 mm^2 , is typically less than 2 μA during optical excitation.

The sample is mounted in an optical cryostat and immersed in helium vapor or liquid. Measured temperatures are determined by a calibrated Si diode thermometer in close proximity to the sample. An 808 nm diode laser, typically pulsed on for 1-2 μs and subsequently switched off, serves to populate the wells. The excitation light is incident off-normal to the wafer to avoid collecting the laser's specular reflection. The excitation energy is tuned by varying the temperature of the laser, typically setting it either in resonance with the DX absorption at trap-bottom or several meV above, i.e. into excitation of the continuum of unbound e-h states. These choices avoid strongly stress-dependent absorption profiles that can be produced by tuning varying parts of the trap into DX resonances.

Light emitted from the sample is collected by an imaging lens external to the cryostat, collecting light up to about 15° from the normal to the wafer and focusing it onto the slit of a 500 mm imaging spectrometer with a 1200 g/mm grating. Images can be generated by opening the spectrometer slit and using the zero-order diffraction of the spectrometer. The diffracted or imaged light is collected by an intensified CCD (ICCD), allowing gate times down to a few nanoseconds. Time-resolved studies

FIG. 3: Conduction and valence band dispersion relations for a GaAs quantum well. The lh/hh splitting is exaggerated by showing the effect for 5 nm wells.

are achieved by triggering both the excitation laser and ICCD by a pulse generator with a programmable delay.

Several equilibration processes must occur to subsequently produce the indirect excitons. Excitons relax in energy, momentum, and angular momentum by interaction with the atomic lattice; they become indirect via one of the constituent particles tunneling through the barrier between the two wells. The rate of these processes in comparison with the radiative lifetime of each intermediate state determines the degree to which each of the higher energy states is visible. For example, the spectrum in Fig. 1(b) shows no sign of the indirect light-hole emission between the hh IX and hh biexciton lines. We understand this to be due to the long radiative lifetime of indirect states in comparison with the lh-hh relaxation time.

A simple estimate of the unstressed dispersion relations for carriers in a GaAs quantum well¹⁶ is displayed in Fig. 3. The split-off band is ignored in this paper due to its large separation in energy. The lh and hh, which are degenerate in bulk at zone center, are split by the quantum well confinement.

As discussed above, a pin pressing on the substrate side of the sample causes a spatially-varying strain profile, tuning the energy of both indirect and direct excitons, primarily through a localized hydrostatic expansion, but accompanied by off-diagonal strain components. The strain and energy profiles resulting from this mechanism are discussed in Section V. The applied force is determined by the reading on the micrometer driving the spring compression, subsequently loading the stressor pin. The wafer substrate is etched, post-fabrication,

FIG. 4: PL Images showing darkening with increasing applied stress, for symmetric coupled quantum wells with width 12 nm and temperature $T=2.2$ K. The laser excitation photon energy was 4 meV above the heavy-hole DX absorption at trap center; $P_{avg}=120 \mu W$. The trap depth changes from 4 meV to 5 meV from $F=2.5$ N to 2.8 N. An external electric potential difference of 0.50 V was applied to augment the built-in field. The horizontal axis (with weaker confinement) corresponds to the $[110]$ crystal axis, and the vertical axis corresponds to $[1\bar{1}0]$.

to 125 μm to facilitate this stress-trapping, providing sharper trap profiles with less force.

Images taken after a few hundred nanoseconds can safely be considered to show excitons in local thermal equilibrium, and images after 1-2 μs show a spatially-equilibrated population in the trapping potential¹⁷. Once the exciton gas is spatially equilibrated, the low-density intensity profiles are expected to have a Gaussian profile, with a width determined by the spring constant of the trapping potential and the exciton temperature. As density is increased, the trap potential is expected to be flattened by the exciton-exciton interaction, causing wider spatial distributions. Previous results at moderate stress followed this model very closely¹⁷.

Following the procedure of our previous work, we gradually increased the depth of the trapping potential while observing the spatial and spectral intensity profiles. As stress was increased, we found a sharp transition to a luminescence profile with a dimmed center, with an intensity maximum tens to hundreds of μm away from the trap center. As discussed below, this luminescence darkening appears at slightly lower stresses with increasing density and/or decreasing temperature.

III. THE PL DARK SPOT

The 2D spatial profile of the PL from the indirect excitons with increasing applied stress is shown in Fig. 4. During this series, the laser was not pulsed, and it was focused on the center of the trap. The laser power and sample temperature were held fixed, at $120\ \mu\text{W}$ and $2.2\ \text{K}$, respectively. The anisotropic spatial profile displayed here, suggesting weaker confinement along one of the cleaving directions, is consistent with the effects of a piezoelectric energy shift, discussed in more depth in Section V. Although temporal gating was not employed in this series to cut out short-lifetime decays like the direct exciton luminescence, a bandpass filter, centered at $819\ \text{nm}$, with a $20\ \text{nm}$ width, was used to show only the indirect exciton luminescence. In this data series, the trap center PL is dimpled at $2.5\ \text{N}$, and, at $2.8\ \text{N}$, a dark spot is prominent in the PL image, with the intensity in a region up to $\sim 50\ \mu\text{m}$ from trap center dimmed by a factor of 3 relative to the intensity at the periphery of the trap.

In Figure 5, we demonstrate that the appearance of this feature depends on the IX density. This series of spectra was taken as we increased the average power of a focused excitation laser while holding all other parameters fixed. The PL profile cross-section displayed is along the weaker confinement axis of the stress trap. The spectrally broad, and spatially narrow, line through the center is substrate luminescence at the location of the excitation laser. The dark spot pattern is observed to appear as the singly-peaked low-density profile evolves to a high-density profile peaking at the perimeter regions of the trap. Both a blue shift in IX energy and a broadening of the linewidth are observed with increasing density, both of which are features of the IX-IX interaction¹⁹.

In Fig. 6, a series of 2D spatial images show the PL as a function of temperature, at a constant stress and constant average excitation power of $200\ \mu\text{W}$, demonstrating that under these conditions this effect appears as a function of temperature between $9\ \text{K}$ and $4\ \text{K}$. The intensity of the emission away from trap center brightens as the temperature is lowered, while, again, the center PL darkens.

To see the effect of the applied stress on the spin states, we measured the linear polarization of the IX luminescence at high stress. The upper half of Fig. 7 displays an image of the IX PL intensity overlaid with a polarization map, indicated by lines oriented parallel to the local linear polarization, with the degree of polarization reflected in each line's length. The CCD was gated to detect the PL $200\ \text{ns}$ after the excitation laser is switched off; hence the observed PL should reflect particles in local thermal equilibrium in IX states. The lattice temperature in this case was $1.4\ \text{K}$. The degree of linear polarization is defined as

$$\text{DOP} = \frac{I_{\text{Max}} - I_{\text{Min}}}{I_{\text{Max}} + I_{\text{Min}}}, \quad (1)$$

where I_{Max} and I_{Min} are the maximum and minimum

FIG. 5: The appearance of the dark spot is shown with increasing density: A series of spatially-resolved spectra with increasing laser excitation power, as stress, temperature, applied bias, and excitation laser duty-cycle remain fixed. The bright vertical line at the center of the trap is impurity luminescence from the substrate, demonstrating the position/profile of the pumping laser. The well width is $12\ \text{nm}$, $T=1.5\ \text{K}$. External bias was $0.6\ \text{V}$.

FIG. 6: The appearance of the dark spot with decreasing temperature. IX PL Images with decreasing temperature at high stress. Excitation $P_{\text{avg}}=50\ \mu\text{W}$. Bias= $0.5\ \text{V}$

values of the intensity at a particular location, measured as the orientation of the linear polarizer is rotated.

In addition to the PL darkening, with further increase in applied force past this PL darkening regime, we find that the center luminescence returns, within a larger region of still-dimmed PL. The intensity map underlying the polarization plot in Fig. 7 shows this pattern. The

FIG. 7: Measured and predicted linear polarization. Experimental linear polarizations at very high stress are displayed in the upper half, with an underlying image of the IX PL intensity. The line orientation and length indicate polarization angle and degree of linear polarization, respectively. The maximum degree of polarization, about 75%, is reached in the dimmed intensity region on the horizontal axis. The intensity image displays the reappearance of PL from trap center at very high stress. The lower half displays the theoretical prediction for linear polarization due to lh IX/hh IX mixing, as discussed in Section IV.

FIG. 8: IX Spectrum at a stress surpassing that of the PL dark spot, showing trap-center luminescence re-emerging.

high-stress IX PL spectrum in Fig. 8 also shows this dramatic reappearance of the PL at trap center. As discussed in Section V, we believe that this reappearance of the luminescence is connected to the crossing of the light-hole and heavy-hole IX bands.

In order to rule out any mechanism by which the pop-

ulation might preferentially occupy the higher energy perimeter states, we have sought to establish that the density remains peaked at trap center. Due to the recently measured monotonic increase of IX PL linewidth with increasing density¹⁹, we use the linewidth as a qualitative metric of the density. Fig. 9 shows the linewidth as a function of position along the trap, peaking in the center, in addition to the profile of the spectrally-integrated IX PL counts. The spectrum was recorded with a delay of 4600 ns after the end of the laser pulse. The contrast is apparent here between the centrally-peaked density profile and the radiation profile, dominated by emission from the perimeter. The significant noise in the linewidth measurement is not from low signal; it is an artifact of a periodic spatial pattern in the ICCD sensitivity, skewing the spectral profile and artificially widening or narrowing the profile. As a consequence, we can only trust this measure to within the margins of its variation, which we still deem sufficient to establish a centrally-peaked density profile.

The figures in this paper are from 12 nm and 14 nm samples, although well-widths of 8 nm and 10 nm were investigated. The same darkening behavior was observed in each. The noteworthy trend in this series of studies was that wider well-widths display the phenomenon at lower stress. This is a strong indication that the lh-hh crossing plays an important role, since the stress at which crossing occurs depends on the well width (see Section IV).

In addition, we have observed this effect, without notable change, in samples identical to the 14 nm well-width samples but lacking superlattice stacks. By this, we may rule out effects primarily owing to the presence of the superlattices.

Finally, we note that a similar dimming of the luminescence occurs in the lowest-energy DX emission from trap center at high stresses. This DX darkening displays no density dependence, and the phenomenon is easily understood to be due to the lh-DX pushing through the hh-DX to become the lowest energy state. The lh-DX is known to have a lower recombination probability than the hh-DX due to its increased penetration into the barrier regions¹¹. Also, the sharper trapping profile of the direct excitons at trap center makes the lh-DX identification simple, as the light-hole energy is more strongly influenced by stress.

This DX effect suggests an examination of the lh/hh IX crossing in order to explain the similar effect in the indirect exciton case. As mentioned in Section II, given the lh IX's position in energy between the hh DX and the hh IX, one might expect that DX excitation should cause some luminescence to be visible from lh IX states as excitons cascade to the lower hh IX states. However, the lh IX line is invisible in nearly all the data sets we have taken. We have, however, been able to observe the lh IX line under conditions that depend strongly on the excitation parameters, i.e. excitation angle and excitation energy. Figure 10 displays a spatially resolved spectrum from one of the experiments in which this line is observed.

FIG. 9: Profiles of IX PL linewidth and spectrally-integrated counts across the weak-confinement axis of the trap. The linewidth is used here as a qualitative measure of density. $T = 2.0$ K, $P_{\text{avg}} = 2$ mW, Bias = 1 V.

FIG. 10: Low-stress, spatially resolved spectrum in the trap, showing, in order of decreasing energy, lh-DX, hh-DX, lh IX, and hh IX. Excitation incidence angle was about 55° from normal. Excitation $\lambda = 797$ nm.

The energy shift of the lh IX state with applied voltage follows that of the hh IX very closely. Note that the lh lines are affected to a greater degree by the applied

stress, in accord with the theory presented in Section IV, providing a sharper trapping profile for those excitons.

IV. A LIGHT/HEAVY HOLE MIXING MECHANISM

Several trends in this behavior suggest an interaction between the light-hole and heavy-hole states to be the underlying cause of this luminescence pattern. First of all, the linear polarization demands a mixing of pure angular momentum states. The spatial pattern of this polarization, rotating about the stressor location, is suggestive that the local strain may determine the extent of this mixing. Next, the trend of decreased stress required to manifest this pattern in wider wells is consistent with this mixing scenario, as the lh IX/hh IX crossing occurs at lower stress in wider wells, where the confinement energy splitting between light and heavy holes is lowered. Furthermore, the reappearance of the trap-center PL at higher stress, along with the analogous dimming in the DX state as the lh-DX is pushed through the hh-DX, indicates that the lh IX state may be pushed to lower energy than the hh IX at very high stress.

In order to make predictions regarding the effect of this interaction, we model the effect of strain on the exciton states and its subsequent effect on the radiative decay rates of these excitons. We seek to predict the lh/hh IX mixing and the density, stress, and temperature dependencies by examining the physics of a single indirect exciton in the strained structure. We begin this analysis by understanding the effect of strain on the energies and interactions of these states.

The trap profile can be determined by the equilibrium strain needed to balance the force of the pin. We calculate the equilibrium strain in the plane of the quantum well using ANSYS, a finite element modelling program, which has been used successfully to fit the stress effects of microcavity polaritons²⁵. The electron energy is roughly given by

$$E_{\text{electron}} = \frac{\hbar^2}{2m_e} \left(\frac{\pi}{a} \right)^2 + a_c(\epsilon_{xx} + \epsilon_{yy} + \epsilon_{zz}). \quad (2)$$

The hole energy is somewhat more complicated and is approximated by the Pikus-Bir Hamiltonian (evaluated at $k_x = k_y = 0, k_z \approx \pi/a$). It is written in the basis of the four spin states of the $j=3/2$ valence band, $m_z=3/2, 1/2, -1/2$ and $-3/2$ (where the $m_z=\pm 3/2$ and $m_z=\pm 1/2$ states correspond to the heavy and light hole states, respectively)¹⁶.

$$E_{\text{hole}} = \begin{bmatrix} \frac{\hbar^2}{2m_{hh}} \left(\frac{\pi}{a} \right)^2 + P_\epsilon + Q_\epsilon & 0 & R_\epsilon & 0 \\ 0 & \frac{\hbar^2}{2m_{lh}} \left(\frac{\pi}{a} \right)^2 + P_\epsilon - Q_\epsilon & 0 & R_\epsilon \\ R_\epsilon^\dagger & 0 & \frac{\hbar^2}{2m_{lh}} \left(\frac{\pi}{a} \right)^2 + P_\epsilon - Q_\epsilon & 0 \\ 0 & R_\epsilon^\dagger & 0 & \frac{\hbar^2}{2m_{hh}} \left(\frac{\pi}{a} \right)^2 + P_\epsilon + Q_\epsilon \end{bmatrix}, \quad (3)$$

where

$$P_\epsilon = -a_v(\epsilon_{xx} + \epsilon_{yy} + \epsilon_{zz}), \quad (4)$$

$$Q_\epsilon = -b(\epsilon_{xx} + \epsilon_{yy} - 2\epsilon_{zz}), \quad (5)$$

and

$$R_\epsilon = \frac{\sqrt{3}b}{2}(\epsilon_{xx} - \epsilon_{yy}) - id\epsilon_{xy}. \quad (6)$$

The parameters in Eqns. (3)–(6) can be found in Chuang¹⁶. The close proximity of the wells to the surface of the sample implies that ϵ_{xz} and ϵ_{yz} are negligible (an approximation that is verified by our strain calculations). We are primarily concerned with the lowest-energy hole state (the highest energy valence state). At zero stress, this will be the doubly-degenerate heavy hole; when ϵ_{xy} is non-zero or $\epsilon_{xx} \neq \epsilon_{yy}$, it will be a mixed heavy- and light-hole state.

The hole energy can be found by diagonalizing the above Hamiltonian, but to find the exciton energy rigorously it is necessary to solve the electron and hole wave functions consistently, taking the Coulomb interaction into account. We performed an exact numerical solution of the two-particle Schrodinger equation for Coulomb-interacting electrons and holes in an external potential of the CQW and an applied electric field, within the anisotropic effective mass approximation. We included the single-particle potential and the Coulomb interaction between the electron and hole on an equal footing, and performed exact diagonalization of the two-particle problem within a finite basis set of one-particle wave functions. The one-particle (electron and hole) wave functions in an external potential were calculated using the conjugate-gradient minimization technique with preconditioning of the steepest descent vector.

Because the calculation is computationally intensive, it was done only at zero stress for a range of electric field values to determine the dipole moment and the oscillator strength (proportional to the square of the overlap integral of the electron and hole wavefunctions). The lh and hh exciton energies are displayed in Fig. 11(a) as a function of electric field in the wells. A linear fit to the calculated Stark shifts was performed to extrapolate between the electric field points lying in our experimental range. Similarly, an exponential dependence was fit to the numerical results for oscillator strength.

Extrapolating to non-zero stress, we approximate the energy of the direct exciton wave function as the sum of the electron and hole energies, with a constant binding energy:

$$E_{exciton} = E_{gap} + E_{electron} + E_{hole} - E_{bind} \quad (7)$$

This energy is fit to a representative experimental low-stress trap profile for the direct excitons, using the applied force, an overall ϵ_{xy} component to the strain, and a constant energy shift (the sum of the gap energy, the

binding energy and the confinement energy) as the fitting parameters. The overall shear strain is a relatively weak effect, as the energy profile is dominated by the shift in the conduction band, but needs to be included to account for the asymmetry of the degree of polarization along the strong and weak axes. The result for the DX exciton energy is shown in Fig. 12. For the direct excitons, the shift due to the DC electric field is not included; direct excitons have a relatively small electric dipole (around 1-2 nm), so the electric field will only introduce a small, constant shift in the direct line (which should be accommodated by the constant energy shift).

The effect of strain on the IX lh and hh states is shown in Fig. 13. At low stress the higher energy state is primarily lh and the lower energy state is primarily hh; at high stress they switch. Near cross-over there is a small anti-crossing due to the built-in shear strain.

For 14 nm quantum wells separated by a 4 nm barrier, the IX effective dipole length should be about 22 nm, according to the numerical results shown in Fig. 11(a). In the limit of high electric field, this will lower the energy of the IX line with respect to the DX line according to the simple approximation

$$\Delta E_{IX} = -e(\mu_{IX} - \mu_{DX})F_{Well}, \quad (8)$$

where μ_{IX} is the IX dipole moment, μ_{DX} is the DX dipole moment, and F_{Well} is the electric field in the wells.

Additionally, there will be a small, stress-dependent electric field caused by the piezoelectric effect. In layered media, the electric field in a particular layer is calculated by requiring that the component of the flux density normal to the interface is the same in each layer and fixing the total voltage drop (V_{app}) across the sample¹⁶. In the well and barrier regions, the piezoelectric field will end up being proportional to the difference in the piezoelectric constants of GaAs and $Al_{0.45}Ga_{0.55}As$:

$$F_W = \frac{\epsilon_B}{L_W\epsilon_B + L_B\epsilon_W}V_{App} + \frac{L_B}{L_W\epsilon_B + L_B\epsilon_W}(P_B - P_W), \quad (9)$$

$$F_B = \frac{\epsilon_W}{L_B\epsilon_W + L_W\epsilon_B}V_{App} + \frac{L_W}{L_B\epsilon_W + L_W\epsilon_B}(P_W - P_B), \quad (10)$$

where L_B and L_W are the total thickness of $Al_{0.45}Ga_{0.55}As$ and GaAs in the insulating region of the sample, respectively (they will be almost equal for our structure), and for GaAs, the relevant piezoelectric polarization can be written as

$$P_B - P_W = (e_{14}^B - e_{14}^W)\epsilon_{xy}. \quad (11)$$

We assume that the strain is the same in the wells and barrier, justified by their very similar elastic constants. The piezoelectric field will also shift the IX energy according to Eq. (8). Rigorously, the electric field felt by the IX is some combination of the fields in the barrier and the wells - for thick wells, it is safe to approximate that the field in the well regions dominates. Assuming

FIG. 11: Results of the numerical solution to two-particle Schrodinger equation in coupled quantum wells. Points marked with a circle or triangle are the results of numerical simulation. (a) The magnitude of the confined Stark shift of the hh and lh IX energies as a function of electric field, with a linear fit in our region of interest to extract the electric dipole moment. (b) The oscillator strength of the hh and lh excitons as a function of applied electric field. A rough fit in the region around our built-in field suggest near exponential dependence ($I_0 e^{-\alpha E}$). The lh value is roughly ten times greater than that of the heavy hole, a ratio which is relatively independent of field in the region of interest.

that the shift of the direct excitons with electric field is negligible, we can find the piezoelectric shift at a given applied force by subtracting the DX and IX lines, and subtracting off the offset at the center of the trap.

The prediction, measurement, and fit of the IX piezoelectric shift is demonstrated in Fig. 14(a-c). The data

displays an asymmetry in the piezoelectric shift, which we attribute to the trap's position near the edge of the window in the clamping plates. The overshoot at the edges of the fit may be due to the density-dependent blueshift in the indirect line. In addition, we can measure a value for the piezoelectric constant from our fitting, obtaining a value of 4.70 eV. We may predict the value of this constant from material parameters from the literature²⁶, which is given by

$$\frac{\Delta E_{IX}}{\epsilon_{xy}} = -e\mu_{IX}(e_{14}^B - e_{14}^W) \frac{L_B}{L_W\epsilon_B + L_B\epsilon_W}, \quad (12)$$

yielding a piezoelectric constant of 2.25 eV. There are a few plausible reasons for this discrepancy. The expression above is strongly dependent on the rather small difference in the piezoelectric constant in the two materials. Due to unavailability of some parameters at low temperature, the calculation was done using room temperature values. If the piezoelectric constants of the two materials have a significant temperature dependence, it could have a drastic impact on the value of this constant. Additionally, variation of the strain in the insulating region of the sample could alter this result. Nonetheless, the close fit is consistent with the prediction that the difference between the IX and DX lines are proportional to ϵ_{xy} , suggesting that the potential can be closely modeled with our strain calculations with the addition of a piezoelectric shift.

It is easy to see from the form of the Pikus-Bir Hamiltonian that ϵ_{xy} induces mixing between the heavy- and light-hole states. If we consider the spin 3/2 heavy hole state, the applied stress will result in mixing with the -1/2 light hole state,

$$|3/2'\rangle = \eta_H |3/2\rangle + \eta_L |-1/2\rangle. \quad (13)$$

FIG. 12: The solid line shows a fit to the measured trap-profile of the direct hh exciton at low stress (higher energy points) using the discussed model. Also shown is the measured trap profile for the corresponding biexciton line (lower energy points).

FIG. 13: The hh and lh indirect exciton energy at trap center for 14 nm (solid lines) and 12 nm (dashed lines) coupled quantum wells as a function of applied force. The higher energy state of each pair at low stress corresponds to the lh IX, while the lower energy state is the hh IX.

A plot of the mixing fraction, $|\eta_L|^2$, as a function of stress is shown in Fig. 15. At reasonable values of stress, this fraction appears to be quite low. If we believe that the suppression of luminescence in the center is caused by a drop in the relative lifetime on the sides of the trap, the lh IX lifetime must be significantly shorter than the hh IX lifetime.

This hypothesis is supported by computational estimates of the indirect heavy- and light-hole exciton lifetimes from the numerical solution. The relative oscillator strengths from the calculation are shown, in units normalized to the zero-field hh DX oscillator strength, in Fig. 11(b). The light hole lifetime is roughly a factor of ten times shorter than that of the heavy-hole for the entire range of these electric fields. This can be understood from the different masses. The smaller out-of-plane mass of the light hole improves its ability to tunnel through the barrier, increasing both the tunneling rate and the electron-hole overlap. The opposite is true of the direct light and heavy hole excitons - the lighter mass of the light hole in the confinement direction allows it to penetrate more deeply into the walls of the well, decreasing the electron-hole overlap within the well. In addition, the larger in-plane mass of the light hole yields a more tightly bound exciton, increasing the e-h overlap for both direct and indirect excitons. However, the strength of this in-plane effect must be small, as it loses out to the effect of tunneling in the case of the hh vs. lh DX oscillator strengths.

The lifetime of the mixed state as a function of the mixing fraction is given by

$$\tau_{Mix} = \frac{1}{f_{hh}\tau_{hh}^{-1} + f_{lh}\tau_{lh}^{-1}}, \quad (14)$$

where f_{lh} is the fraction of lh IX mixed into the lowest energy state, as plotted in Fig. 15, and f_{hh} is the remaining hh IX fraction. Even for a small light hole fraction, the radiative rate can become significantly faster. For a light hole fraction of 0.09 seen in Fig. 18, the oscillator strength ratio of 10 from the theory prediction shown in Fig. 14(b) implies that the radiative rate will be a factor of two greater at the edge of the trap than in the center, where the excitons are purely heavy hole states.

Another symptom of light-heavy hole mixing is the appearance of linear polarization in the luminescence. The position of shear-dependent terms in the Pikus-Bir Hamiltonian matrix induces mixing between the $J = +1$ heavy hole states and the $J = -1$ light hole states. Without mixing, these would produce right- and left-handed circularly polarized luminescence, respectively.

To find the polarization of the luminescence as a function of position in the stress-induced trap, we need to find the recombination rate as a function of polarization angle. To simplify the calculation, we assume the emission angle is perpendicular to the plane of the quantum wells. The recombination rate, R , is proportional to the square of the matrix element²⁷,

$$R \propto |\langle u_c | \hat{\epsilon}(\theta) \cdot \vec{p} | u_v \rangle|^2, \quad (15)$$

where $|u_c\rangle$ and $|u_v\rangle$ are the conduction and valence band states comprising the IX, $\vec{p}$ is the momentum operator, and $\hat{\epsilon}(\theta)$ is the polarization unit vector of the emitted photon. The intensity observed normal to the plane of the quantum well is proportional to the rate of photons emitted in that direction due to recombination. For normal emission, the polarization angle can be expressed in terms of x and y components in the plane of the quantum wells,

$$\hat{\epsilon}(\theta) = \begin{pmatrix} \cos(\theta) \\ \sin(\theta) \end{pmatrix}. \quad (16)$$

Excitons formed from the unmixed heavy or light hole states will emit circularly polarized light, will have no linear polarization dependence and will exhibit zero degree of polarization. The mixing coefficients η_H and η_L will be complex numbers whose relative phase determines the polarization and whose magnitudes determine the relative fraction of heavy and light hole in the ground state. To find the polarization dependence of the emission, we decompose the hole states into spherical harmonics and calculate the matrix elements:

$$I(\theta) \propto \frac{|\eta_H|^2 I_H^2}{4} + \frac{|\eta_L|^2 I_L^2}{12} + \frac{|\eta_H||\eta_L| I_H I_L}{2\sqrt{3}} \cos(2\theta + \theta_R), \quad (17)$$

FIG. 14: The spatial dependence of the ϵ_{xy} component of the shear strain as predicted by calculations (a), is compared to the energy difference between the indirect and direct lines (minus an overall constant) in (b). In (c), the calculated value of the strain (solid line) is fit to the measured difference between the indirect and direct lines, using a multiplicative constant as a fitting parameter.

where I_H and I_L represent the overlap integrals of the electron and hole states in the z-direction for the heavy- and light-hole excitons respectively. The angle θ_R is the complex angle of the R-term in the Pikus-Bir Hamiltonian (which arises from shear strain). From the theory prediction shown in Fig. 14(b), we approximate $I_L = \sqrt{10}I_H$. Note that some level of mixing is required for any degree of linear polarization. From this equation, we can determine the polarization angle of maximum intensity and the degree of polarization; for a representative stress, a plot of these as a function of position is

FIG. 15: (a) A numerical model of the expected spatial dependence of the lh-fraction in the lower energy IX state as a function of applied stress.

shown in the lower half of Fig. 7, represented by the angle and length of the lines respectively.

Qualitatively, this compares favorably with the experimental data, displayed above those predicted results. Both data sets display polarization angles favoring the cleaving axes (horizontal and vertical), rather than displaying purely radial polarizations. In order to replicate the strong/weak-confinement axis asymmetry appearing in the measurements, it is necessary to include a small, constant component of ϵ_{xy} that enhances the amount of mixing along the weak-confinement axis and suppresses the mixing along the strong-confinement axis. This strain could result from the piezoelectric effect due to the built-in electric field. This asymmetry of the $[110]$ and $[\bar{1}\bar{1}0]$ axes can be easily seen in both data sets as the polarization along the vertical axis changes non-monotonically with distance from the trap center, in contrast to the steady rise along the horizontal axis.

With some certainty regarding the nature of the lh/hh mixing provided by the preceding polarization analysis, along with the simulation results predicting the considerably larger lh IX oscillator strength, we can make a qualitative argument predicting the manifestation of the dark spot in the luminescence pattern. The appearance of the dark spot with increasing density can be understood as coming from the exciton-exciton repulsion which causes the exciton cloud to spread out spatially as density is increased. This causes the excitons at the periphery of the cloud to lie in regions of higher radiative recombination rate.

The equilibrium density can be calculated using a Maxwell-Boltzmann distribution and assuming that the exciton-exciton interaction energy is proportional to the local density,

$$n(r) = n_0 e^{-\frac{V(r) + \gamma n(r)}{kT}}, \quad (18)$$

where the potential energy, $V(r)$, is determined by the

FIG. 16: Predicted IX recombination PL intensity profiles with increasing force applied to the sample, using the theory of a spatially-varying oscillator strength.

FIG. 17: Predicted IX recombination PL intensity profiles with increasing exciton density.

strain profile and γ is the proportionality constant relating the blue shift with density (for reasonable values, see discussion in Appendix). This equation is solved self-consistently at each point in the trap profile using an iterative method to determine the density. The particle number is fixed by integrating the density over the trap and re-scaling n_0 after each iteration. This density $n(r)$ is multiplied by the radiative rate from Eq. (14) to get the total intensity at each point r . The result is shown in Fig. 16 for different applied forces, as the total particle number and temperature are held constant. In Fig. 17, the applied force and temperature are fixed, and the total particle number is increased over several orders of magnitude. The strong exciton-exciton repulsion pushes particles away from the high-density central region and into the shorter lifetime regions on the side of the trap.

This analysis of the spatial profile indicates that the appearance of the dark spot can be at least partially understood in terms of the mixing of the light hole and heavy hole states in the exciton single-particle spectrum. However, the same analysis implies predictions for the overall population lifetime which do not appear to be consistent with the experiments.

If we approximate a constant shape to the density profile, varying only by a constant factor with time, then we can define a population lifetime τ_{pop} for a given exciton density profile. We write

$$\frac{\partial N}{\partial t} = - \int d^2r \frac{n(r)}{\tau(r)} = - \frac{N}{\tau_{\text{pop}}}, \quad (19)$$

i.e., since each local region has a decay rate which is linear with density, the lifetime of the total population will be a constant that depends on the recombination probability profile, which is a function of the light-hole/heavy-hole

mixing, and the exciton density profile $n(r)$.

The local recombination rate, given by the inverse of Eq. (14), suggests that the lifetime of states on the periphery of the trapped exciton cloud should fall drastically with respect to that of trap-center states when the lh-hh mixing begins. This implies that the overall population lifetime will decrease dramatically as light-hole/heavy-hole mixing begins. This is not observed in the experiments; we see in Fig. 18(b) that the population lifetime slowly increases with increased force, throughout the onset of the dark hole phenomenon and continuing after the bright lh IX luminescence appears in the center of the trap at high stress.

However, the mechanism suggested here to produce the dark spot pattern does not require the total population lifetime to decrease in order to display a relative dimming of the luminescence. It allows for the possibility that the total lifetime increases, as long as the lifetime of a particle at trap center increases faster than that of the total population, so that states located on the periphery remain the dominant recombination channel. Thus some other effect could be giving an overall longer population lifetime, compensating the decrease due to light-hole mixing.

We do expect a gradual overall lifetime increase due to the piezoelectric electric field generated by the shear strain. This field creates regions of higher electric field along the weaker confinement axis, and this augmented electric field causes both higher occupation in this region as well as an extended lifetime. However, our calculations indicate that this effect should produce an increase in the lifetime of only a couple percent, while the population lifetime from mixing should vary smoothly from the hh IX lifetime at low stress to near the lh IX lifetime (a factor of ten shorter) once the lh IX state is pushed below the

FIG. 18: (a) IX Luminescence as a function of excitation laser P_{avg} . No significant deviations from the linear fit emerge as the new luminescence pattern appears. (b) IX lifetime vs. trap depth. Trap depth is indicative of total applied force and the magnitude of strain. Lifetimes plotted are decay constants of single-exponential decays fit to the time-resolved IX PL. The inset shows this fitting for one point in the lifetime plot.

hh IX state in energy. Therefore, in the stress regime where the dark center is present, we expect a population lifetime quite a bit shorter than the hh IX lifetime.

Another shortcoming of an explanation entirely in terms of the lh-hh mixing is that it fails to explain the temperature dependence. The decay rate profile predicted here should be nearly independent of temperature; gaining only an overall increase in decay rate with decreasing temperature.

Furthermore, two discrepancies are evident in comparison with the actual intensity profiles. First, the degree of the darkening of the center luminescence with increasing stress seems to be too low in our prediction. The largest observable drop in the luminescence predicted in the model is around 40-50%, while, in the experimental data, a drop of 80% is not unusual. Using a light hole oscillator strength of ten times that of the heavy hole, this is achievable in the mixing model with 50% mixing, which is possible at extremely high shear strain. However, at these strain values the model predicts the light hole IX energy significantly below the heavy hole IX at trap center, suggesting this should not occur before the lh IX luminescence appears at trap center. The other incompatibility of our predictions with the measured profiles is that in general the high intensity lobes are closer to the center in the model. The location of the dimmed region is also puzzling in our polarization/intensity plot in Fig. 7, where the polarization along the horizontal axis is substantial even in the dimmed intensity region. In contrast, we expect both the linear polarization and the local decay rate to be proportional to the lh/hh mixing. However, the polarization profile matches well between

the theory and experiment, suggesting that our mixing analysis is correct, but that additional physics, such as many-body effects, are needed to explain all the data.

V. OTHER EXPLANATIONS AND PERSPECTIVES

The most obvious possible cause of this PL darkening is that we could simply be losing excitons; either by another exciton decay mechanism, which manifests at non-zero stress, or by a change in the emission character with stress that hinders light collection from the center of the trap. The simplest checks for this are to demonstrate a decrease in the measured lifetime or in the amount of total collected light as this effect manifests.

Fig. 18(a) shows a measure of the total collected light from the IX PL as the excitation intensity is increased; a linear increase in the collected PL suggests that we do not begin to lose more emitted photons as the effect turns on. We choose to measure the collected light as a function of density, rather than stress, because a change in the energy profile of the trap could easily affect the laser absorption. For this measurement, we opened the spectrometer slit to allow luminescence from the entire emitting area to be collected by the CCD. This allowed us to sum the total indirect exciton luminescence by collecting all first-order diffracted light in the spectral region of IX luminescence. The data is well fit by a linear increase in the collected light with average power, showing no obvious deviations at the density of dark spot onset. This suggests that, while there are excitons that are not

FIG. 19: Relative particle number required for dark spot onset vs. temperature for two different trap depths. The $N^{1/2}$ line suggested by the non-interacting BEC theory is shown for comparison. Data shown for 12 nm wells.

collected by our detection system, the proportion missed is fixed throughout the investigated range; we do not see a sublinear increase that would suggest a sudden increase in lost excitons – from a new decay mechanism or from a change in the collecting efficiency with power.

In addition, a series of lifetime measurements was performed at decreasing stress values, beginning in a high-stress regime where the dark spot pattern is prominent and ending with a moderate-stress profile peaked at trap center. The indirect exciton lifetime measurement was taken by collecting the total intensity in the same manner as the total intensity vs. excitation power experiment, by opening the slit to integrate the diffracted PL. The lifetimes in Fig. 18(b) are plotted against trap depth, which increases with applied force. The relatively large error bars are mostly due to the relatively long lifetimes in comparison to the repetition period. Longer repetition periods would be ideal to measure the decay, allowing more of the time evolution to be observed, but they severely diminish the signal/noise due to the reduction in experiments performed.

Another possible scenario explored is that of condensation of a dark ($J=2$) population. In this scenario, bright states attain a higher energy by exchange interactions that require e-h recombination as an intermediate process. With a lower energy dark state, the stage would be set for a dark condensate, one not dipole-coupled to photons. This possibility was predicted recently for any system of indirect excitons²¹. The idea that condensation may be playing a role in the appearance of the dark spot in our experiments is supported by the fact that the onset criteria seem to correspond to increased degeneracy.

The critical temperature for BEC of a trapped, 2D, weakly-interacting, and dilute Bose gas is given in terms of the trap's spring constant, ω , and the total particle

number, N , as the following²²:

$$T_{crit} = \frac{\sqrt{6}\hbar\omega}{k_B\pi} N^{1/2}. \quad (20)$$

A plot of the relative particle number required for an off-center intensity peak vs. the sample temperature is plotted in Fig. 19, along with the $N^{1/2}$ dependence suggested by the criteria for weakly-interacting BEC. The relative number is plotted here, rather than total particle number, due to the difficulty of measuring the absolute number of excitons. The particle number after pumping for a time long compared to the lifetime is expected to be simply the exciton generation rate multiplied by the IX lifetime. Hence, the excitation laser power multiplied by the lifetime at each temperature should be proportional to density at each temperature. This is the quantity plotted as relative particle number in the figure.

On the other hand, we may make some estimates of the absolute particle number in our experiments, either using the measured excitation power and the expected efficiency of exciton creation or by the observed blue shift. These estimates are discussed in the Appendix to this paper. We estimate that the value labeled '1' on the particle number axis in Fig. 19 corresponds to a particle number of roughly 4×10^5 IXs. If we put this number of particles within the spatial distribution before the PL dark spot appears, we obtain a density of $8 \times 10^9 \text{ cm}^{-2}$.

We can make a naive calculation of the weakly-interacting BEC temperature for this number of particles. Using an exciton mass of $0.2m_e$ and a spring constant of 30 eV/cm^2 to calculate ω , this particle number suggests a critical temperature of 2 K. If we consider the dark excitons to be doubly-degenerate from the electron spin, we would reduce this critical temperature to 1.4K.

The saturation of the critical temperature at high density is not surprising given the strong exciton-exciton repulsion, which can cause divergence from the weakly-interacting theory. At very high density we expect that a plot of the critical temperature vs. density would turn over, requiring lower temperatures for increasing density²³.

Despite the suggestive correspondence of the onset criteria to that of BEC, no other condensation-like behavior is evident thus far. In addition, we find that a simple local density approximation solution, where the density is allowed to be a non-monotonic function of the effective chemical potential, $(\mu - V(r))$, due to a condensed dark population, cannot by itself explain the PL profile's density dependence. For example, while we do see an actual dimming of the center luminescence with increased strain, we never see a dimming with increasing density, as that analysis would demand. Rather, the brightness of the periphery of the cloud increases faster as density increases. Thus, while we do not rule out many-body effects, we determine that they are not the sole cause of this luminescence pattern.

VI. DISCUSSION AND CONCLUSIONS

We have demonstrated a marked dimming of the luminescence from indirect excitons in the high-density region at trap center. This effect occurs with increasing stress, decreased temperature, and increasing density. Our collection efficiency of the PL from these excitons remains constant through the onset of this PL pattern, and the lifetime of the particles is shown to steadily increase, suggesting that no drastic changes are occurring in the decay process.

We have presented a theoretical analysis predicting a stress-dependent, spatially-varying fraction of the lh IX state mixed into the lowest energy indirect state. The exact numerical solutions of the two-particle Schrodinger equation for Coulomb-interacting electrons and holes in structures used in the experiment suggest a factor of ten difference in the lh IX oscillator strength over the hh oscillator strength, in contrast to the direct case, where the hh DX state is known to have a higher oscillator strength than the lh DX state. Taken together, these results suggest a relative darkening of the IX luminescence from the center of the trap near the stress of light-hole/heavy-hole indirect exciton crossing, and reasonable assumptions indicate that it may occur only at higher density.

There are three large discrepancies with this analysis and the experiment. First, the predicted intensity profiles are significantly less dimmed and have the brighter lobes closer to the center. Second, we are unable to account for the temperature dependence. Last, the predicted lifetime decrease is not observed.

There are a number of possible ways the discrepancy in intensity profiles may be interpreted. First, it is possible that the light hole is even brighter than predicted by our calculations. On the other hand, the oscillator strength and mixing fraction both factor into the polarization model, which seems to match the experimental results. Second, it would be possible to achieve higher shear strain in the lobes prior to the hh-lh crossover point if the zero-stress energy splitting between the heavy and light hole states was larger than expected.

Another interesting option is that the splitting between the light and heavy hole is density dependent. The asymmetry of the polarization and luminescence profiles lead us to believe there is some built-in shear strain, which should imply non-zero mixing at the center of the trap. Because that is the only component of the shear strain at trap-center, the hh-lh mixing at that point will be entirely governed by the splitting. In other words, if the splitting increases with density (if the lh line blue shift in the presence of hh excitons is larger than the hh line blue shift), then high density at the center of the trap would suppress mixing there, further darkening it. Exploratory calculations including this effect fail to significantly improve the contrast, but they succeed in moving the bright PL lobes further from the center of the trap.

The temperature dependence remains unexplained because, in order to create the observed spatial pattern, this

mechanism requires only an increased oscillator strength on the perimeter and an exciton cloud that extends into this enhanced recombination region. Hence, the role of temperature to control this transition remains puzzling. The exciton cloud size should only increase with increased temperature, adding to the density in the enhanced region. Higher temperature does increase the occupation of higher in-plane momentum states, lengthening the overall lifetime and, most likely, lengthening spin relaxation lifetimes by decreasing the mobility. It is conceivable that these changes in time scales are significant. Again, there still remains the possibility of Bose statistics playing an important role in occupation of the spin states, which would clearly show threshold-like behavior, but the mechanism by which Bose statistics might cause the darkening is unclear.

As discussed above, an explanation solely in terms of the lh-hh crossing suggests that the recombination lifetime should shorten drastically as we observe this pattern form, while we experimentally measure the opposite. If the theory of stress-induced lh/hh mixing is correct, then an as-yet-unexplained mechanism exists that counteracts the predicted drastic lifetime decrease. Even supposing an additional mechanism to lengthen the lifetime with increasing stress that allows the lh IX lifetime to increase more slowly than the hh IX lifetime (hence, still allowing the pattern formation mechanism suggested here), a decrease in lifetime of the IX population should be observed as the stress is increased to the point where the lh IX becomes the lowest energy state. The lack of this feature in the lifetime data suggests that there may be an intermediate decay process (e.g. spin relaxation) responsible for buffering the overall recombination lifetime.

One idea is that we have a large population of dark excitons with a very long spin flip time, so that the IX lifetime is controlled by the time required to flip from dark to bright. The dark spot effect would still be produced by the bright-exciton decay rate profile, but the decay rate would be limited, at least after the initially populated bright states have decayed. This line of thinking is motivated by the recently observed long spin flip time in similar GaAs CQW systems²⁴, but the lower limits quoted in that work are many orders of magnitude from the 20 μ s relaxation time required for this explanation, since the spin flip time in this case would have to be comparable to the observed population lifetime.

Apart from the assumption of IX-IX repulsion which gives an expanding cloud with increased exciton density, all of the analysis presented here has been single-particle physics, so perhaps the cause of these divergences from our predictions lies in many-body physics. The puzzle of these discrepancies remains an open question for further investigation.

The fact remains that there must be mixing of spin states occurring as a function of stress, e.g. this is demanded by the linear polarization of the luminescence. We have achieved qualitative agreement between these measurements and predicted linear polarizations from

lh/hh mixing.

We have observed the lh IX as it becomes the lowest energy state, becoming a bright spot at trap center, and we have also observed PL from this state at lower stresses, when the excitation conditions and branching ratios in the relaxation process were favorable for recombination in the higher energy lh IX state. The proximity in energy of this lh IX state to the lower energy hh IX state suggests that the mixing is indeed between these states, as predicted by our analysis.

Preliminary studies of the lh IX luminescence appearing at high stress in the center of the trap suggest that these traps have spring constants about 4 times that of the hh IX stress traps. The increased confinement offered by these states, along with the mysteriously extended population lifetimes at high stress, promise increased degeneracy among these lh IXs. Further studies will aim to find degenerate many-body phenomena, particularly condensation, within this system of trapped light-hole, indirect excitons.

Appendix A: Density Estimate

There are two methods to ascertain a density range for the data presented here. First, we may exploit the linear dependence of the PL energy with density to relate a measured blue shift with a density estimate. Second, we can try to account for the pumping efficiency to relate a measured excitation power to a exciton creation rate.

To begin the density estimate based on blue shift, we must first realize that the proportionality factor is highly disputed. The simplest estimate for the linear dependence of blue shift on density derives from the ‘capacitor formula’, depending only on the density, n , the well separation, d , and the permittivity in the barrier, which is

$$\Delta_{cap} = \frac{de^2}{\epsilon} n. \quad (A1)$$

For 14 nm wells, this suggests a factor of 1.9×10^{-10} meV cm². Hence, for a routinely observed blue shift of 2 meV, we would predict a density of 10^{10} cm⁻².

One may go one step beyond this to consider the effect of lowered temperature to cause an anti-correlation of the excitons; the excitons rarely have enough kinetic energy to approach each other against the repulsive interaction^{28,29}. Hence, these high energy regions barely figure into the exciton energy and the blue shift is greatly reduced at low temperature, with the correction factor multiplying the capacitor estimate reaching 0.1 at about 5 K. Thus, we would increase that estimated density by factor of 10 to 10^{11} cm⁻² for a blue shift of 2 meV.

Experimentally, the proportionality factor is difficult to measure owing to the lack of a decent absolute exciton density measurement. Yet, it can be measured with a large degree of uncertainty by employing an approximated absolute density based on the expected light absorption¹⁷. That measurement yielded 0.9 ×

10^{-10} meV cm² in 12 nm wells at 5K (where the capacitor formula predicts 1.6×10^{-10} meV cm²). Thus, in that case, the proportionality constant fell between the capacitor formula prediction and the corrected value, though a bit nearer the capacitor formula.

Using a similar method to that in Ref. 16, we may make an estimate of the number of excitons in our traps by accounting for our photon losses and conversion efficiency. By measuring our optical transmission, accounting for losses in the barrier between the surface and the wells, and estimating our absorption in the quantum well, using a DX absorption length³⁰ of $1.65 \mu\text{m}^{-1}$, we obtain a pumping rate of direct excitons. We measure the DX to IX conversion efficiency by comparing the brightness of the DX and IX spectral lines during CW excitation, and achieve an IX pumping rate. Taking into account the IXs lost during pumping by recombination, we can make a guess at the total number of IXs around when the laser is turned off.

Combining the efficiency factors, we achieve 0.0015 indirect excitons per incident photon, yielding 4.8×10^{11} excitons per second for a peak laser power of 80 μW at 800 nm. Thus, we estimate in Fig. 19 that the position labeled ‘1’ on the axis marking relative critical particle number corresponds to a particle number of roughly 4×10^5 IXs. If we put this number of particles within the spatial distribution before the dark spot appears in the PL image, we obtain a peak density of 8×10^9 cm⁻².

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- * nws5@pitt.edu
- † Current address: Universität Innsbruck, Christoph-Probst-Platz, Innrain 52, A - 6020 Innsbruck, Austria
- ‡ M.H.Szymanska@warwick.ac.uk
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