Observation of Coulomb-Assisted Dipole-Forbidden Intraexciton Transitions in Semiconductors

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We use terahertz pulses to induce resonant transitions between the eigenstates of optically generated exciton populations in a high-quality semiconductor quantum-well sample. Monitoring the excitonic photoluminescence, we observe transient quenching of the 1s exciton emission, which we attribute to the terahertz-induced 1s-to-2p excitation. Simultaneously, a pronounced enhancement of the 2s-exciton emission is observed, despite the 1s-to-2s transition being dipole forbidden. A microscopic many-body theory explains the experimental observations as a Coulomb-scattering mixing of the 2s and 2p states, yielding an effective terahertz transition between the 1s and 2s populations.

Coulombically bound electron-hole pairs, i.e., excitons, often dominate the optical properties [1] of high-quality semiconductors at low temperatures. In direct-gap semiconductors, the intra-exciton energy spacing is typically in the terahertz (THz) frequency range, with the corresponding intra-excitonic transitions following the same selection rules as the dipole transitions in hydrogen atoms. Therefore, the presence of exciton populations can be unambiguously detected by monitoring the 1s-to-2p transition resonance [2–10]. One can also use strong THz pulses to induce nonlinear 1s-to-2p transitions [11–15], including Rabi flopping [16–19] and excitations in Λ systems [20].

An elegant way to study the influence of THz-induced exciton transitions is to monitor time-resolved photoluminescence (PL) spectra after combined optical and THz excitations [21]. Since only the optically active s states contribute to the PL, the THz-induced intra-excitonic 1s-to-2p population transfer can be observed as pronounced quenching of the 1s PL [22–24]. Similar PL quenching can be observed for quantum wells (QWs) when the THz energy is resonant with a dipole-allowed intersubband transition [25].

In this Letter, we apply this technique to monitor how a THz pulse changes the time-resolved PL long after a resonant optical pulse has generated exciton population in the system. Besides 1s-PL quenching, we observe an unexpected transient increase of the 2s-PL indicating a pronounced effective 1s-to-2s transition. Using a systematic many-body theory [26], we show that the THz-induced 1s-to-2p excitation is accompanied by an efficient 2p-to-2s transition which can be attributed to co-operative Coulomb scattering. This Coulomb-assisted THz-induced 1s-to-2s coupling is unique to interacting many-body semiconductor configurations and cannot be observed in atomic systems.

In our experiments, we study the intra-excitonic transitions in a 20× In_{0.06}Ga_{0.94}As QW structure at a lattice temperature of 10 K. The 8 nm wide QWs are separated by 130 nm GaAs barriers, rendering all nontrivial coupling effects negligible [27]. As such, the structure acts as a single QW whose PL magnitude is additively enhanced by the number of QWs [28]. The heavy-hole (hh) excitons are well separated from both the continuum and the light-hole states allowing us to address the resonances individually and efficiently yielding a two-band situation for the hh1 [29, 30]. For this particular In-Ga ratio, the energy difference between the 1s and 2s hh1 excitons is estimated as 6.7 meV and the 1s-to-2p transition energy is 6.9 meV.

We monitor the time-resolved PL after an excitation sequence where a THz pulse of a free-electron laser (FEL) follows an optical pulse of a Ti:Sapphire laser, as shown schematically in Fig. 1(a). The Ti:Sapphire laser emits pulses with a duration of 4 ps (FWHM) at a repetition rate of 78 MHz with a photon energy of 1.471 eV which is in resonance with the 1s hh state. The repetition rate is reduced to 13 MHz by an extra-cavity pulse picker equipped with an acousto-optic modulator to match the repetition rate of the FEL. The FEL emits 30 ps pulses with a wavelength of 191 μm which is in resonance with the 1s-to-2p hh exciton transition energy. The peak field strength of the FEL beam at the position of the sample is estimated to be 5 kV cm⁻¹.

The two sources are synchronized electronically and their time delay is controlled by using a mechanical delay of the synchronizing pulses; for details, see [31]. The two beams are collinearly polarized and focussed directly onto
the QW. The spot sizes are chosen to be 300 and 50 µm for the FEL and the Ti:Sapphire laser, respectively. The PL is collected in a forward-scattering geometry through the transparent GaAs substrate. Only the center spot of about 20 µm in diameter is imaged within a small solid angle, carefully avoiding the transmitted laser beams as well as density-averaging effects. The PL is spectrally and temporally dispersed using a spectrometer attached to a streak camera with energy and time resolutions of 0.15 meV and 15 ps, respectively; the time is oversampled, collecting a spectrum every 4.4 ps.

The relatively weak optical pulse, that is \( \approx 10^{11} \text{ photons/(cm}^2\text{)} \) per pulse, couples to the 1s polarization which is converted into incoherent 1s excitons. This polarization-to-population conversion process occurs efficiently with a characteristic time scale of \(<10\) ps, as shown in Ref. [32]. Exemplary PL data without THz excitation are given in the left panel of Fig. 1(b) where the spectrally and temporally emitted PL intensities are plotted in false colors. The time \( t \) is defined with respect to the center of the THz pulse. Following the initial excitation, the PL decays exponentially on a time scale of 600 ps. The corresponding data for dual excitation are shown in the right panel. The THz pulse arrives 225 ps after the optical pulse, thus effectively exciting the incoherent 1s population into the 2p state and partially into the ionization continuum [19]. This population transfer quenches the 1s population and hence the related 1s resonance in PL as predicted in Refs. [22, 23]. Additionally, a clear spike in the 2s emission is visible when the FEL is incident, and the 1s PL recovers on a much longer time scale.

To better quantify the THz-induced changes, we plot emission spectra for three representative \( t \) with (solid line) and without (shaded area) THz excitation in Figs. 1(c)-1(e). For a short time delay of \( t = 50 \) ps just after the THz pulse [panel (c)], the 1s PL intensity decreases by 46%, verifying the usual THz-field-induced quenching scenario. However, in addition, the THz excitation induces a pronounced 2s resonance; this observation is unexpected because the direct 1s-to-2s transition is dipole forbidden. The 1s PL then recovers its intensity gradually while the 2s peak decays, as seen in (d) and (e). For \( t = 1200 \) ps, we observe another interesting feature: the 1s PL becomes larger with than without THz, indicating that more luminescing 1s excitons are present in the system long after the THz excitation. This is explained by THz-induced shelving of the overall exciton populations into optically dark states that cannot recombine radiatively. As the excitations relax back to the 1s state, we eventually observe excess 1s PL at later times because excitons experience a reduced overall radiative decay during the THz excitation-relaxation cycle.

Note that nonradiative recombination in these samples is negligible for the chosen excitation conditions as the overall time-integrated emission intensities with and without THz excitation match within the experimental error.

We follow the time evolution of these THz-induced phenomena by determining the differential photoluminescence \( \Delta PL(t) = PL_{THz} - PL_{ref} \) between the cases with \( (PL_{THz}) \) and without THz excitation \( (PL_{ref}) \). The measured \( \Delta PL \) spectra are shown in Fig. 2(a) as a contour plot. Again, the 1s quench (early times), 1s shelving (later times), and the 2s-excess PL are clearly visible. To monitor \( \Delta PL \) dynamics in more detail, Figs. 2(b) and 2(c) present \( \Delta PL_{1s} \) and \( \Delta PL_{2s} \), corresponding to energy-integrated spectra around the indicated 1s and 2s energies, respectively, as a function of time. The strong

![FIG. 1. (color online). (a) Schematic experimental setup. (b) False-color representation of the temporally and spectrally resolved PL without (left) and with (right) THz excitation. (c)-(e) Measured reference PL spectra (shaded area) vs. THz-induced PL without (left) and with (right) THz excitation in Figs. 1(c)-1(e). For a short time delay of \( t = 50 \) ps just after the THz pulse [panel (c)], the 1s PL intensity decreases by 46%, verifying the usual THz-field-induced quenching scenario. However, in addition, the THz excitation induces a pronounced 2s resonance; this observation is unexpected because the direct 1s-to-2s transition is dipole forbidden. The 1s PL then recovers its intensity gradually while the 2s peak decays, as seen in (d) and (e). For \( t = 1200 \) ps, we observe another interesting feature: the 1s PL becomes larger with than without THz, indicating that more luminescing 1s excitons are present in the system long after the THz excitation. This is explained by THz-induced shelving of the overall exciton populations into optically dark states that cannot recombine radiatively. As the excitations relax back to the 1s state, we eventually observe excess 1s PL at later times because excitons experience a reduced overall radiative decay during the THz excitation-relaxation cycle. Note that nonradiative recombination in these samples is negligible for the chosen excitation conditions as the overall time-integrated emission intensities with and without THz excitation match within the experimental error. We follow the time evolution of these THz-induced phenomena by determining the differential photoluminescence \( \Delta PL(t) = PL_{THz} - PL_{ref} \) between the cases with \( (PL_{THz}) \) and without THz excitation \( (PL_{ref}) \). The measured \( \Delta PL \) spectra are shown in Fig. 2(a) as a contour plot. Again, the 1s quench (early times), 1s shelving (later times), and the 2s-excess PL are clearly visible. To monitor \( \Delta PL \) dynamics in more detail, Figs. 2(b) and 2(c) present \( \Delta PL_{1s} \) and \( \Delta PL_{2s} \), corresponding to energy-integrated spectra around the indicated 1s and 2s energies, respectively, as a function of time. The strong...](image-url)
energy of the excitons are defined by the two-particle correlations within one of the QWs. The microscopic properties for this purpose, we concentrate on the excitation dynamics THz-induced exciton transitions and the related PL. For this purpose, we concentrate on the excitation dynamics THz-induced exciton transitions and the related PL. For this purpose, we concentrate on the excitation dynamics THz-induced exciton transitions and the related PL. For this purpose, we concentrate on the excitation dynamics THz-induced exciton transitions and the related PL.

\[ E_{1s} = \langle c^e_{\mathbf{k}} | \hat{h}^* | c^{h}_{\mathbf{k}} \rangle. \]

The corresponding electron (hole) distribution is \( f^e_n = \langle c^e_{\mathbf{k}} | \hat{h}^* | c^e_{\mathbf{k}} \rangle \). In the so-called main-sum approximation [1, 26], the PL dynamics is given by

\[
\hbar \frac{\partial}{\partial t} c^q_{\mathbf{k'}, \mathbf{q'}} = E^e_{\mathbf{k'}, \mathbf{q'}} c^q_{\mathbf{k'}, \mathbf{q'}} - A_{\text{THz}}(t) \cdot j_{\mathbf{k'} + \mathbf{q} - \mathbf{k}} c^q_{\mathbf{k'}} c^q_{\mathbf{k'}}
+ (1 - f^e_{\mathbf{k'}} - f^h_{\mathbf{k'} - \mathbf{q}}) \sum_{l} V_{l - k} c^q_{\mathbf{k'}, \mathbf{l}}
- (1 - f^e_{\mathbf{k'} + \mathbf{q}} - f^h_{\mathbf{k'}}) \sum_{l} V_{l - k'} c^q_{\mathbf{k'}, \mathbf{l}} + T q_{\mathbf{k'}, \mathbf{k}}, \tag{1}
\]

where \( E^e_{\mathbf{k'}, \mathbf{q'}} \) contains the renormalized energy of an electron-hole pair, \( V_{\mathbf{k}} \) is the Coulomb-matrix element, \( A_{\text{THz}}(t) \) is the vector potential of the THz pulse, and \( j_{\mathbf{k}} = -e|\mathbf{b}|/\mu \) is the current matrix element with the reduced electron–hole mass \( \mu \). The three-particle correlations are symbolically denoted by \( T q_{\mathbf{k'}, \mathbf{k}} \), see the Supplemental Material (SM) for more details.

As shown in Ref. [26], \( T \) is dominantly built up via the Boltzmann-type Coulomb scattering where exciton correlations exchange momentum with the plasma and the other two-particle correlations. Consequently, \( T \) becomes a complicated functional of exciton populations due to the quantum kinetics involved. However, one does not need to determine the full quantum kinetics explicitly to explain the consequence of the \( T \)-related Coulomb scattering on THz transitions. Instead, we only need to consider that the incoming excitons scatter into new momentum states with a constraint that the number of incoming and outgoing exciton correlations remains constant.

One can explain the consequences of the diffusive Coulomb scattering on THz transitions using an ansatz

\[ T^q_{\text{diff}} = -i \hbar \gamma \left[ c^q_{\mathbf{k'}, \mathbf{k}} - \frac{1}{2\pi} \int_0^{2\pi} d\theta_{\mathbf{k}} c^q_{\mathbf{k'} + \mathbf{K}, \mathbf{k} + \mathbf{K}} \right]. \tag{2} \]

where \( \theta_{\mathbf{k}} \) is the direction of the scattering \( \mathbf{K} \) that is assumed to have a constant magnitude and \( \gamma \) defines the overall scattering strength. The introduced \( T^q_{\text{diff}} \) is a generalization of the diffusive model [32] that explains the principal effects of excitation-induced dephasing [26] beyond the constant-dephasing approximation.

To determine the effect of diffusive Coulomb scattering on the THz-generated 2p populations, we use the exciton transformation \( \Delta N^\lambda_{\mathbf{q}} = \sum_{\mathbf{k'}, \mathbf{k}} \phi_{\lambda}(\mathbf{k}) \phi_{\lambda}^*(\mathbf{k'}) c^q_{\mathbf{k'}, \mathbf{k}} \) where \( \phi_{\lambda}(\mathbf{k}) \) is the exciton wave function and \( c^q_{\mathbf{k'}, \mathbf{k}} \) is the center-of-mass representation of \( c_{\mathbf{k}} \) [26]. The specific exciton populations are given by the diagonal elements \( \Delta N^\lambda_{\mathbf{q}, \mathbf{q}} \) while the off-diagonal elements \( \Delta N^\lambda_{\mathbf{q}, \mathbf{q'}}, \mathbf{q}\neq\mathbf{q'} \) define

\[ \Delta N^\lambda_{\mathbf{q}, \mathbf{q'}} = \sum_{\mathbf{k}, \mathbf{k'}} \phi_{\lambda}(\mathbf{k}) \phi_{\lambda}^*(\mathbf{k'}) c^q_{\mathbf{k'}, \mathbf{k}} \]

\[ \Delta N^\lambda_{\mathbf{q}, \mathbf{q'}} = \sum_{\mathbf{k}, \mathbf{k'}} \phi_{\lambda}(\mathbf{k}) \phi_{\lambda}^*(\mathbf{k'}) c^q_{\mathbf{k'}, \mathbf{k}} \]

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\[ \Delta N^\lambda_{\mathbf{q}, \mathbf{q'}} = \sum_{\mathbf{k}, \mathbf{k'}} \phi_{\lambda}(\mathbf{k}) \phi_{\lambda}^*(\mathbf{k'}) c^q_{\mathbf{k'}, \mathbf{k}} \]
the exciton transitions. The diffusive scattering from 2p to 2s can be deduced by projecting Eqs. (1)–(2) with the exciton transformation and following the 2p contributions, yielding \[ \frac{\Delta N_{2s}^{2s}}{\Delta N_{2p}^{2p}} \bigg|_{2p} = \Delta N_{2p,2p}^{2p}/\tau_{\text{conv}} \] (see SM) where we have defined a scattering time

\[ \frac{2\pi}{\tau_{\text{conv}}} = \gamma \int_0^{2\pi} d\theta K \left| \sum_k \phi_{2p}(k) \phi_{2s}(k + K) \right|^2. \] (3)

We see that the 2p-to-2s coupling is present as long as \( K = 0 \).

Besides the 2p-to-2s coupling, the Coulomb interaction also relaxes 2s populations toward the quasi-equilibrium on a time scale \( \tau_{\text{rel}} \). The combined effect of \( \tau_{\text{conv}} \) and \( \tau_{\text{rel}} \) creates new Coulomb-mediated eigenstates where 2s and 2p state become mixed, see SM. In particular, the Coulomb-induced state mixing induces an effective THz transition between the original 1s and 2s states. Note that the dipole-allowed 1s-to-2p and 2p-to-2s transitions cannot generate efficient 1s-to-2s population conversion without Coulomb scattering because the THz pulse is off-resonant with the 2p-to-2s transition.

Due to the scattering nature of 2s-2p mixing, the created 2s population decays with rate \( \tau_{2s}^{-1} = \tau_{\text{rel}}^{-1} - \tau_{\text{conv}}^{-1} \) which defines also the fast decay of the excess 2s PL, as shown in the SM. For the late times, the 2s population reaches a quasi-equilibrium, yielding a slower decay of the 2s PL on the time scale of the remaining phonon relaxation \( \tau_{\text{phon}} \).

We numerically solve THz dynamics (1) including all the relevant exciton states for the optically bright and dark excitons, and the diffusive Coulomb scattering (2). The center-of-mass momentum of bright and dark excitons is fully taken into account. We also include the radiative decay of bright excitons as well as the relaxation of excitons toward the thermodynamic equilibrium on a \( \tau_{\text{phon}} = 900 \text{ ps} \) time scale, agreeing well with independent microscopic computations [6, 26]. The diffusive Coulomb scattering is chosen to give \( \tau_{2s} = 120 \text{ ps} \) (\( \tau_{\text{conv}} = 56.0 \text{ ps} \), \( \tau_{\text{rel}} = 38.2 \text{ ps} \)) that is substantially faster than \( \tau_{\text{phon}} \). The quasi-stationary PL spectra are computed via the PL-Elliott formula [1], as discussed in the SM.

Figure 3 shows the PL changes induced by THz excitation when the diffusive Coulomb scattering is included under the same excitation conditions as in Fig. 2. The computations not only explain the qualitative behavior of 1s quench, exciton shelving, and excess 2s PL, but they quantitatively determine the 1s quench and 2s excess levels. They also explain the double-decay of the 2s PL as switching from the fast Coulomb-equilibration \( \tau_{2s} \) to the slow phonon-relaxation \( \tau_{\text{phon}} \). We also have performed a computation where we reduce \( 1/\tau_{\text{conv}} \) by a factor of 20 in the \( c_\lambda \) dynamics. The dashed line in Fig. 3(c) compares this \( \Delta \text{PL}_{2s} \) with the full result (solid line). We see that the full and \( 1/\tau_{\text{conv}} \)-reduced computations decay similarly at the late times; note that some excess 2s PL remains due to thermal relaxation from the ionized excitons toward the 1s and 2s states. However, only the full computation produces a strong \( \Delta \text{PL}_{2s} \) transient that decays fast. Hence, the \( \Delta \text{PL}_{2s} \) transient does not originate from relaxation but follows from the Coulomb-induced population transfer and the subsequent equilibration. The \( \Delta \text{PL}_{2s} \) peak also emerges on a time scale similar to the THz excitation, determined by the THz-pulse duration (here 30 ps). Thus, the Coulomb interaction cooperates with the THz excitation to open a new 1s-to-2s transition that is much faster than the relaxation processes.

In conclusion, our experiment-theory analysis shows the existence of Coulomb-assisted THz transitions converting 1s into 2s excitons, i.e., a process that extends the dipole-selection rules as a direct consequence of many-body interactions. This effect survives even when an appreciable amount of disorder is present, as shown in the SM. The related 2p-to-2s transfer is significantly faster than other relaxation processes, making the Coulomb-induced scattering an active partner in the THz transitions. This work not only highlights a pronounced difference between excitons and atoms, but also opens up a

![FIG. 3. (color online). THz-induced effect on PL. (a) Computed differential PL. Integrated differential \( \Delta \text{PL}_\lambda(t) \) for (b) \( \lambda = 1s \) and (c) \( \lambda = 2s \). The conditions are same as in Fig. 2.](attachment:image)
new mechanism to manipulate excitons through combining many-body and THz effects.

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[1] M. Kira and S. W. Koch, Semiconductor Quantum Optics, 1st ed. (Cambridge Univ. Press, 2011)
[2] T. Timusk et al., Solid State Comm. 25, 217 (1978)
[3] D. Labrie et al., Phys. Rev. Lett. 61, 1882 (1988)
[4] C. C. Hodge et al., Phys. Rev. B 41, 12319 (1990)
[5] R. H. M. Groeneveld and D. Grischkowsky, J. Opt. Soc. Am. B 11, 2502 (1994)
[6] M. Kira et al., Phys. Rev. Lett. 87, 176401 (2001)
[7] R. A. Kaindl et al., Nature 423, 734 (2003)
[8] M. Kira, W. Hoyer, and S. W. Koch, Solid State Commun. 129, 733 (2004)
[9] I. Galbraith et al., Phys. Rev. B 71, 073302 (2005)
[10] T. Suzuki and R. Shimano, Phys. Rev. Lett. 103, 057401 (2009)
[11] K. B. Nordstrom et al., Phys. Rev. Lett. 81, 457 (1998)
[12] M. Kubouchi et al., Phys. Rev. Lett. 94, 016403 (2005)
[13] R. Huber et al., Phys. Rev. Lett. 96, 017402 (2006)
[14] J. R. Danielson et al., Phys. Rev. Lett. 99, 237401 (2007)
[15] S. Leinß et al., Phys. Rev. Lett. 101, 246401 (2008)
[16] C. W. Luo et al., Phys. Rev. Lett. 92, 047402 (2004)
[17] S. G. Carter et al., Science 310, 651 (2005)
[18] M. Wagner et al., Phys. Rev. Lett. 105, 167401 (2010)
[19] B. Ewers et al., Phys. Rev. B 85, 075307 (2012)
[20] J. L. Tomaino et al., Phys. Rev. Lett. 108, 267402 (2012)
[21] R. Ulbricht et al., Rev. Mod. Phys. 83, 543 (2011)
[22] J. Cerne et al., Phys. Rev. Lett. 77, 1131 (1996)
[23] M. S. Salib et al., Phys. Rev. Lett. 77, 1135 (1996)
[24] J. Kono et al., Phys. Rev. Lett. 79, 1758 (1997)
[25] S. Zybell et al., Appl. Phys. Lett. 99, 041103 (2011)
[26] M. Kira and S. W. Koch, Prog. Quantum Electron. 30, 155 (2006)
[27] M. Hubner et al., Phys. Rev. Lett. 83, 2841 (1999)
[28] M. Schafer et al., Phys. Rev. B 74, 155315 (2006)
[29] S. Chatterjee et al., Phys. Rev. Lett. 92, 067402 (2004)
[30] T. Grunwald et al., Phys. Status Solidi C 6, 500 (2009)
[31] J. Bhattacharyya et al., Rev. Sci. Instrum. 82, 103107 (2011)
[32] M. Kira and S. W. Koch, Phys. Rev. Lett. 93, 076402 (2004)