# First-order spatial coherence of excitons in planar nanostructures: A $k$-filtering effect

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We propose and analyze a $k_{\parallel}$-filtering effect which gives rise to the drastic difference between the actual spatial coherence length of quasi-two-dimensional excitons or microcavity polaritons in planar nanostructures and that inferred from far-field optical measurements. The effect originates from conservation of the in-plane wave-vector $k_{\parallel}$ in the optical decay of the particles in outgoing bulk photons. The $k_{\parallel}$-filtering effect explains the large coherence lengths recently observed for indirect excitons in coupled quantum wells but is less pronounced for microcavity polaritons at low temperatures, $T \lesssim 10$ K.

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Long-range spatial coherence is a fingerprint of well-developed Bose-Einstein (BE) statistics. Measurements of the first-order spatial coherence function $g^{(1)}$ and the coherence length $\xi$ have allowed visualization of the BE condensation transition in a trapped Bose gas of Rb atoms.[1] There are several recent reports on the observation of long-range spatial optical coherence in a low-temperature quasi-two-dimensional (quasi-2D) system of microcavity (MC) polaritons[2,3] and indirect excitons.[4–7] In this case, the resonant optical decay of MC polaritons or quantum well (QW) excitons in bulk photon modes allows mapping of the in-plane coherence function $g^{(1)}$ of the particles by measuring the optical coherence function $\tilde{g}^{(1)}$ of the emitted photons. It is commonly assumed that the coherence length of QW excitons (MC polaritons), $\xi_x(\xi_p)$, associated with $g^{(1)}$, is identical to that, $\xi_{\gamma}$, of the optical coherence function $\tilde{g}^{(1)}$.

In this Brief Report, we report a $k_{\parallel}$-*filtering effect*, which can strongly influence the optical coherence function $\tilde{g}^{(1)}$ measured from a planar nanostructure, and calculate $g^{(1)}$ and $\tilde{g}^{(1)}$ for QW excitons and MC polaritons. For QW excitons, the $k_{\parallel}$-filtering effect tremendously increases the optical coherence length $\xi_{\gamma}$, leading to $\xi_{\gamma} \gg \xi_x$, and can naturally explain the micron coherence lengths observed for indirect excitons and attributed to spontaneously developed coherence. The effect is less pronounced for MC polaritons, still with $\xi_{\gamma} \gtrsim \xi_p$.

The $k_{\parallel}$-filtering effect stems from the energy and in-plane momentum conservation in the resonant conversion "quasi-2D QW exciton (MC polariton) → outgoing bulk photon." For a QW structure surrounded by thick coplanar barrier layers, the case illustrated in Fig. 1, only low-energy optically active excitons from the radiative zone $k_{\parallel} \le k_0 = (\sqrt{\varepsilon_b}/c)\omega_0$, with $\varepsilon_b$ the dielectric constant of barrier layers and $\hbar\omega_0$ the exciton energy at $k_{\parallel}=0$, are bright, i.e., can emit far-field light.[8–11] In a far-field optical experiment with detection angle $2\alpha$ [see Fig. 1(b)], the fraction of QW excitons which contribute to the optical signal is drastically further reduced to the wave-vector band $\Delta k_{\parallel}$ given by $0 \le k_{\parallel} \le k_{\parallel}^{(\alpha)} = (k_0/\sqrt{\varepsilon_b})\sin\alpha \ll k_0$. The $\alpha$-dependent narrowing of the detected states results in an effective broadening of the first-order spatial coherence function $\tilde{g}^{(1)}$. In addition, the sharp cutoff of the detected states at $k_{\parallel}=k_{\parallel}^{(\alpha)}$ yields an unusual oscillatory behavior of $\tilde{g}^{(1)}$. The $k_{\parallel}$-filtering effect has no analogy in optics of bulk excitons or polaritons.

The first-order spatial coherence function $g^{(1)}$ (Refs. 12 and 13) of quantum well excitons, at a fixed time, is given by $g^{(1)}(\mathbf{r}_{\parallel}',\mathbf{r}_{\parallel}'')=G^{(1)}(\mathbf{r}_{\parallel}',\mathbf{r}_{\parallel}'')/[G^{(1)}(\mathbf{r}_{\parallel}',\mathbf{r}_{\parallel}')G^{(1)}(\mathbf{r}_{\parallel}'',\mathbf{r}_{\parallel}'')]^{1/2}$, with $G^{(1)}(\mathbf{r}_{\parallel}',\mathbf{r}_{\parallel}'')=\langle\hat{\Psi}^{\dagger}(\mathbf{r}_{\parallel}')\hat{\Psi}(\mathbf{r}_{\parallel}'')\rangle$, where $\hat{\Psi}(\mathbf{r}_{\parallel}') = (1/\sqrt{S})\Sigma_{\mathbf{k}_{\parallel}} e^{i\mathbf{k}_{\parallel}\mathbf{r}_{\parallel}'} B_{\mathbf{k}_{\parallel}}$, $\mathbf{r}_{\parallel}'$ is the in-plane coordinate, $S$ is the area, and $B_{\mathbf{k}_{\parallel}}$ is the exciton operator. Thus for isotropically distributed QW excitons one receives

$$g^{(1)} = g^{(1)}(r_{\parallel}) = \frac{1}{2\pi n_{2d}}\int_0^{\infty} J_0(k_{\parallel}r_{\parallel})n_{k_{\parallel}}k_{\parallel}dk_{\parallel}, \tag{1}$$

where $r_{\parallel}=|\mathbf{r}_{\parallel}''-\mathbf{r}_{\parallel}'|$, $n_{2d}$ is the concentration of particles, $n_{\mathbf{k}_{\parallel}} = \langle B_{\mathbf{k}_{\parallel}}^{\dagger}B_{\mathbf{k}_{\parallel}}\rangle$ is the occupation number, and $J_0$ is the zeroth-order Bessel function of the first kind. For a classical gas of QW excitons at thermal equilibrium, Eq. (1), with $n_{\mathbf{k}_{\parallel}}$ given by the Maxwell-Boltzmann (MB) distribution function $n_{k_{\parallel}}^{\mathrm{MB}}$, yields the well-known result,[3,14]

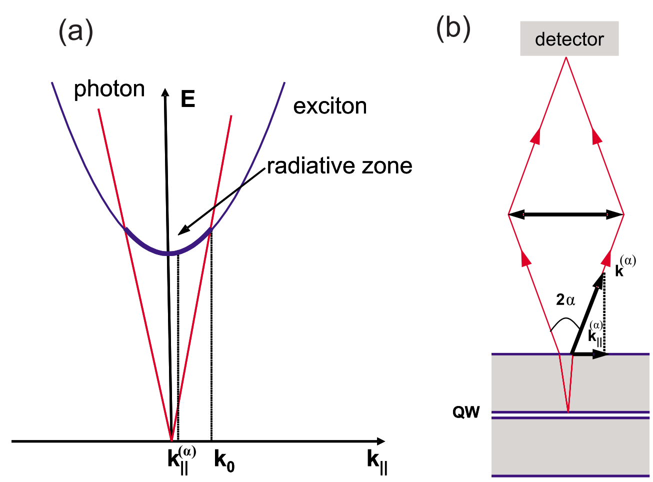


FIG. 1. (Color online) Schematic of the $k_{\parallel}$-filtering effect. (a) The exciton and photon dispersions. Only low-energy QW excitons from the radiative zone $k_{\parallel} \le k_0$ can emit outgoing bulk photons. (b) A far-field optical experiment with detection angle $2\alpha$: a small fraction of QW excitons with $|\mathbf{k}_{\parallel}| \le k_{\parallel}^{(\alpha)} = (k_0/\sqrt{\varepsilon_b})\sin\alpha$ contributes to the optical signal.

 

$$g^{(1)} = g^{(1)}_{\mathrm{cl}}(r_\parallel) = e^{-\pi r_\parallel^2/\lambda_{\mathrm{dB}}^2}, \tag{2}$$

where the thermal de Broglie wavelength is given by $\lambda_{\mathrm{dB}} = [(2\pi\hbar^2)/(M_x k_B T)]^{1/2}$, with $T$ the temperature and $M_x$ the exciton in-plane translational mass. For helium temperatures, one estimates from Eq. (2) the coherence length of MB-distributed indirect excitons in GaAs coupled QWs as $\xi_x \sim \lambda_{\mathrm{dB}} \sim 0.1$ $\mu$m.

Compared with Eq. (1), the spatial coherence function $\tilde{g}^{(1)}$ of photons emitted by QW excitons is given by

$$\tilde{g}^{(1)}(r_\parallel) = \frac{\int_0^\infty G_f(k_\parallel) J_0(k_\parallel r_\parallel) n_{k_\parallel} k_\parallel dk_\parallel}{\int_0^\infty G_f(k_\parallel) n_{k_\parallel} k_\parallel dk_\parallel}, \tag{3}$$

where $G_f = \Theta(k_\parallel^{(\alpha)} - k_\parallel)\Gamma_{x-\gamma}(k_\parallel)$ is the $k_\parallel$-filtering function, with $\Theta(x)$ the step function and $\Gamma_{x-\gamma}(k_\parallel)$ the efficiency of the resonant conversion of a QW exciton in an outgoing bulk photon. The function $G_f$ reduces the integration limits on the right-hand side (rhs) of Eq. (3) to the narrow band $\Delta k_\parallel = [0, k_\parallel^{(\alpha)}]$ and describes the $k_\parallel$-filtering effect in high-quality planar nanostructures. If both the function $\Gamma_{x-\gamma}(k_\parallel)$ and the occupation number $n_{k_\parallel}$ do not change significantly in the narrow band $\Delta k_\parallel$, Eq. (3) yields

$$\tilde{g}^{(1)} = \tilde{g}^{(1)}_f(r_\parallel) = 2J_1(k_\parallel^{(\alpha)} r_\parallel)/(k_\parallel^{(\alpha)} r_\parallel), \tag{4}$$

where $J_1$ is the first-order Bessel function of the first kind. From Eq. (4) one concludes that the optical coherence length $\xi_\gamma$, evaluated as the half width at half maximum of $\tilde{g}^{(1)} = \tilde{g}^{(1)}_f(r_\parallel)$, is given by

$$4J_1(k_\parallel^{(\alpha)}\xi_\gamma) = k_\parallel^{(\alpha)}\xi_\gamma \rightarrow k_\parallel^{(\alpha)}\xi_\gamma \simeq 2.215. \tag{5}$$

Equations (4) and (5) illustrate the net $k_\parallel$-filtering effect in the absence of instrumental aberrations: $\xi_\gamma \propto 1/k_\parallel^{(\alpha)} \propto 1/\sin\alpha$ strongly increases with decreasing aperture angle $2\alpha$. Below we analyze in more detail the exciton function $g^{(1)}$ against the optical $\tilde{g}^{(1)}$, assuming no phase transition to a collective (superfluid) state.

*First-order spatial coherence of noninteracting quasi-2D bosons (excitons) in equilibrium*. In this case, the chemical potential $\mu_{2d}$ is given by $\mu_{2d}^{(0)} = k_B T \ln(1 - e^{-T_0/T})$, with $k_B T_0 = (2\pi/g)(\hbar^2/M_x) n_{2d}$ the quantum degeneracy temperature where $g$ is the spin degeneracy factor of bosons ($g=4$ for indirect excitons). By substituting $n_{k_\parallel} = n_{k_\parallel}^{\mathrm{BE}}$ into Eq. (1), where $n_{k_\parallel}^{\mathrm{BE}}$ is the Bose-Einstein occupation number, one receives

$$g^{(1)} = g^{(1)}_{\mathrm{nint}}(r_\parallel) = \frac{T}{T_0} g_1(1 - e^{T_0/T}, e^{-\pi r_\parallel^2/\lambda_{\mathrm{dB}}^2}) = \frac{T}{T_0}\sum_{n=1}^{\infty} \frac{(1 - e^{-T_0/T})^n}{n} e^{-\pi r_\parallel^2/n\lambda_{\mathrm{dB}}^2}. \tag{6}$$

Here, the generalized Bose function[14] $g_\nu(x,y)$ with $\nu=1$ is defined as $g_\nu(x,y) = \Sigma_{k=1}^{\infty}(x^k y^{1/k})/k^\nu$.

For distances $r_\parallel \gtrsim r_\parallel^{(q)} = \lambda_{\mathrm{dB}}[-(2/\pi)\ln(1 - e^{-T_0/T})]^{1/2}$ Eq. (6) reduces to the quantum limit when the sum on the rhs cannot be approximated by the first term,

$$g^{(1)}(r_\parallel \gtrsim r_\parallel^{(q)}) \simeq 2\frac{T}{T_0} K_0\left(\frac{r_\parallel}{r_0}\right), \tag{7}$$

where $K_0$ is the modified Bessel function of the second kind and $r_0 = \lambda_{\mathrm{dB}}/[-4\pi \ln(1 - e^{-T_0/T})]^{1/2}$. For $r_\parallel \gtrsim r_0 \gtrsim r_\parallel^{(q)}$, Eq. (7) reduces further to

$$g^{(1)} = g_q^{(1)}(r_\parallel \gtrsim r_0) = \sqrt{2\pi}\frac{T}{T_0}\sqrt{\frac{r_0}{r_\parallel}} e^{-r_\parallel/r_0}. \tag{8}$$

For temperatures $T \gg T_0$, the spatial coherence function is well approximated by Eq. (2), and the quantum corrections given by Eqs. (7) and (8) refer to large $r_\parallel \gtrsim r_\parallel^{(q)} \simeq \lambda_{\mathrm{dB}}\sqrt{(2/\pi)\ln(T/T_0)} \gg \lambda_{\mathrm{dB}}$ and, therefore, to very small values of $g^{(1)}$. For $T \lesssim T_0$, when BE statistics is well developed, Eqs. (7) and (8) are valid for distances larger than $r_\parallel^{(q)} \simeq \lambda_{\mathrm{dB}}\sqrt{(2/\pi)} e^{-T_0/2T} \ll \lambda_{\mathrm{dB}}$, so that $g^{(1)}$ is well approximated by $g_q^{(1)}$ for any $r_\parallel$. The quantum statistical effects, which are included in Eq. (7) through $T_0 \propto \hbar^2$, considerably increase the correlation length $\xi_x$, giving rise to $\xi_x \simeq [\lambda_{\mathrm{dB}}/(2\sqrt{\pi})] e^{T_0/2T}$ for $T \lesssim T_0$ (see Fig. 2).

*The coherence function $g^{(1)}$ of weakly interacting thermal QW excitons*. For circularly polarized excitons in single QWs, the case relevant to MC polaritons, the repulsive interaction between the particles is well approximated by a contact potential $U_{\mathrm{SQW}} = (u_0/2)\delta(\mathbf{r}_\parallel)$, with $u_0 = u_0^{\mathrm{SQW}} > 0$. In this case, the mean-field Hartree-Fock (HF) interaction only shifts the chemical potential, $\mu_{2d} = \tilde{\mu}_{2d}^{(0)} = \mu_{2d}^{(0)} + u_0 n_{2d}$, leaving unchanged Eqs. (6)–(8).

For indirect excitons in coupled QWs, the mid-range dipole-dipole repulsive interaction $U_{\mathrm{CQW}}$ of the particles cannot be generally approximated by a contact potential. Following Ref. 15, we use the two-parametric model potential $U_{\mathrm{CQW}}(r_\parallel) = [(\sqrt{\pi}u_0 w)/r_\parallel^3](1 - e^{-r_\parallel^2/w^2})$, with parameters $u_0 = u_0^{\mathrm{CQW}} \simeq 4\pi(e^2/\varepsilon_b)d_z$ (Refs. 16 and 17) and $w \simeq a_x^{(2d)}$, where $d_z$ is the distance between coupled quantum wells and $a_x^{(2d)}$ is the radius of an indirect exciton. The model potential reproduces $1/r_\parallel^3$ behavior at $r_\parallel \gtrsim a_x^{(2d)}$ and $1/r_\parallel$ Coulomb repulsive potential at $r_\parallel \lesssim a_x^{(2d)}$. The self-consistent HF analysis[18] of the Hamiltonian $H_x = \Sigma_{\mathbf{p}_\parallel}[p_\parallel^2/(2M_x)]B_{\mathbf{p}_\parallel}^\dagger B_{\mathbf{p}_\parallel} + 1/(2S)\Sigma_{\mathbf{p}_\parallel,\mathbf{l}_\parallel,\mathbf{q}_\parallel} U_{\mathrm{CQW}}(\mathbf{q}_\parallel) B_{\mathbf{p}_\parallel}^\dagger B_{\mathbf{l}_\parallel}^\dagger B_{\mathbf{l}_\parallel+\mathbf{q}_\parallel} B_{\mathbf{p}_\parallel-\mathbf{q}_\parallel}$ yields the $n_{2d}$- and $T$-dependent change in the in-plane translational mass $M_x$. In this case, $\mu_{2d}$ is

$$\mu_{2d} = \tilde{\mu}_{2d}^{(0)} + \frac{u_0}{2(\lambda_{\mathrm{dB}}^*)^2}\left\{\frac{T_0^*}{T} + \sqrt{\pi}\frac{w}{\lambda_{\mathrm{dB}}^*}\left[\frac{\sqrt{\pi}}{2}\frac{w}{\lambda_{\mathrm{dB}}^*}\mathrm{Li}_2(F) - \mathrm{Li}_{3/2}(F)\right]\right\}, \tag{9}$$

where, together with Eq. (6), both the de Broglie wavelength $\lambda_{\mathrm{dB}}^*$ and the degeneracy temperature $T_0^*$ are changed according to $M_x \rightarrow M_x^*$, $F = 1 - e^{-T_0^*/T}$, and $\mathrm{Li}_\nu(x) = \Sigma_{k=1}^{\infty} x^k/k^\nu$ is the polylogarithm. The particle mass $M_x^*$ renormalized by the dipole-dipole interaction is given as a single solution of the transcendental equation,

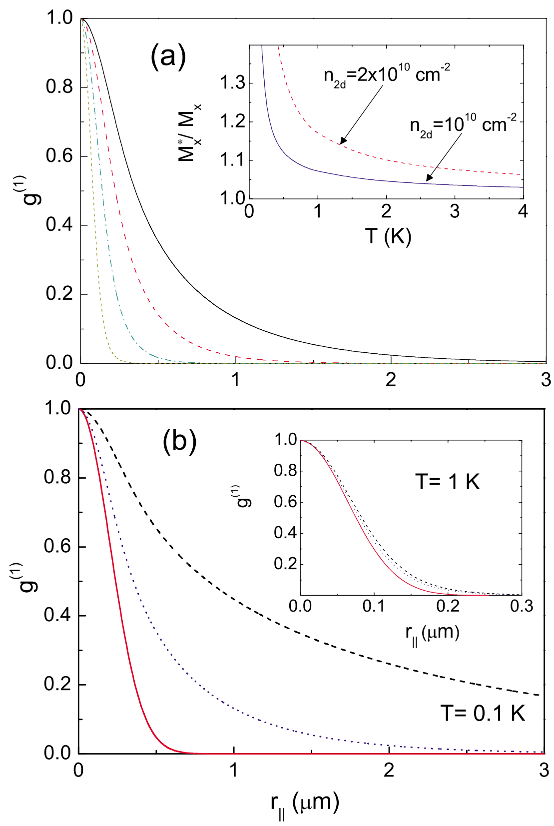


FIG. 2. (Color online) (a) The first-order spatial coherence function $g^{(1)}=g^{(1)}_{\mathrm{ind}}(r_{\|})$ of indirect excitons in a GaAs coupled QW structure with $d_z=11.5$ nm and $w=15$ nm: $n_{2d}=10^{10}$ cm$^{-2}$ and $T=1$ (dotted line), 0.4 (dash-dotted line), 0.2 (dashed line), and 0.1 K (solid line). Inset: the renormalized mass $M^*_x$ against temperature $T$, calculated with Eq. (10). (b) $g^{(1)}=g^{(1)}_{\mathrm{cl}}(r_{\|})$ (solid line), $g^{(1)}=g^{(1)}_{\mathrm{nint}}(r_{\|})$ (dashed line), and $g^{(1)}=g^{(1)}_{\mathrm{ind}}(r_{\|})$ (dotted line): $n_{2d}=10^{10}$ cm$^{-2}$ and $T=0.1$ K. Inset: the same functions evaluated for $n_{2d}=10^{10}$ cm$^{-2}$ and $T=1$ K.

$$\frac{1}{M^*_x}=\frac{1}{M_x}+\frac{u_0 w}{8\sqrt{\pi}\hbar^2\lambda^*_{\mathrm{dB}}}\left[\sqrt{\pi}\frac{w}{\lambda^*_{\mathrm{dB}}}\frac{T^*_0}{T}-\mathrm{Li}_{1/2}(F)\right]. \tag{10}$$

In Fig. 2(a) we plot $g^{(1)}=g^{(1)}_{\mathrm{ind}}(r_{\|})$ evaluated numerically by using Eqs. (6), (9), and (10) for indirect excitons in a GaAs coupled QW structure. In Fig. 2(b), the coherence function $g^{(1)}_{\mathrm{ind}}$ is compared with $g^{(1)}_{\mathrm{cl}}$ evaluated with Eq. (2) and $g^{(1)}_{\mathrm{nint}}$ calculated with Eq. (6) for noninteracting excitons. The main result is that the dipole-dipole repulsive interaction induces an increase in the translational mass [see the inset of Fig. 2(a)] and, therefore, decreases the coherence length $\xi_x$ compared to that of noninteracting particles [see also Fig. 3(a)]. The effect, however, becomes pronounced only at temperatures well below 1 K. For $T=1$ K all three correlation functions, $g^{(1)}_{\mathrm{ind}}$, $g^{(1)}_{\mathrm{cl}}$, and $g^{(1)}_{\mathrm{nint}}$, nearly coincide, as is clearly seen in the inset of Fig. 2(b). In other words, for $n_{2d}=10^{10}$ cm$^{-2}$ and $T=1.5$ K, which are relevant to the experiments,[4–7] the quantum limit, i.e., $g^{(1)}=g^{(1)}_q$ given by Eq. (8), cannot build up: One estimates $T_0\simeq T^*_0\approx 0.65$ K and $n^{\mathrm{BE}}_{k_{\|}=0}\approx 0.54<1$, so that BE statistics is rather weakly developed to influence the coherence length $\xi_x$.

The given description of $g^{(1)}$ refers to temperatures above

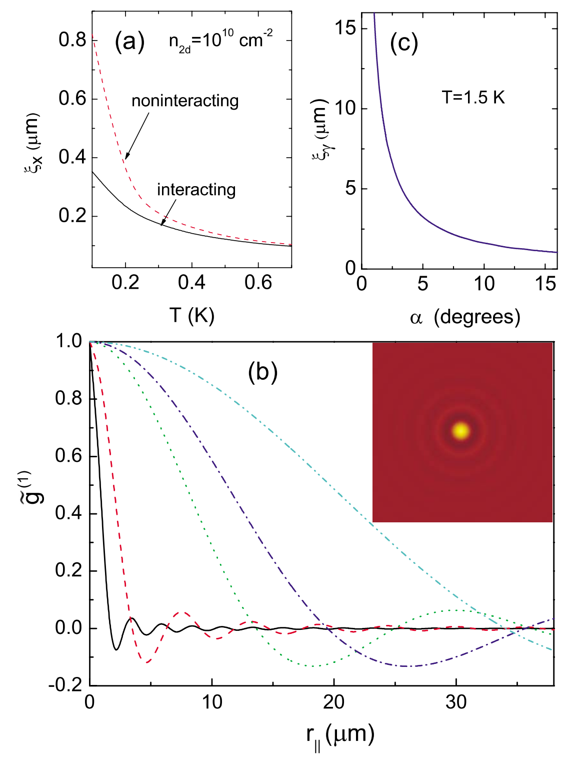


FIG. 3. (Color online) (a) The dependence of the correlation length $\xi_x$ against temperature $T$, calculated for noninteracting (dashed line) and dipole-dipole interacting (solid line) indirect excitons. (b) The $k_{\|}$-filtering effect: $\tilde{g}^{(1)}=\tilde{g}^{(1)}(r_{\|})$ evaluated for $\alpha=18.9°$ (solid line), 8.3° (dashed line), 2.1° (dotted line), 1.4° (dash-dotted line), and 0.8° (dash-double-dotted line). Inset: the real-space 2D image of $\tilde{g}^{(1)}$. (c) The coherence length $\xi_\gamma$ against the aperture angle $2\alpha$.

$T_0$, i.e., when classical or weakly developed BE statistics are realized, and to a quantum gas of indirect excitons at $T\le T_0$, but still above the phase transition temperature. In all these cases the correlation function for a quasi-2D system of weakly interacting excitons is universally given by Eq. (6).

*The optical spatial coherence function* $\tilde{g}^{(1)}$ *of indirect excitons*. In order to explain the experiments,[4–7] which demonstrate a coherence length $\xi_\gamma$ much larger than $\xi_x\sim 0.1$ μm, we implement the concept of $k_{\|}$-filtering. In this case, $\tilde{g}^{(1)}=\tilde{g}^{(1)}_{\mathrm{ind}}(r_{\|})$ is given by Eq. (3) with the efficiency of the "indirect exciton → bulk photon" conversion $\Gamma_{x-\gamma}=(2k_0^2-k_{\|}^2)/[k_0(k_0^2-k_{\|}^2)^{1/2}]$.[8–10,19] In Fig. 3(b), we plot $\tilde{g}^{(1)}_{\mathrm{ind}}$ calculated for various aperture angles, $2°\lesssim 2\alpha\lesssim 40°$. The dependence $\tilde{g}^{(1)}=\tilde{g}^{(1)}_{\mathrm{ind}}(r_{\|})$ is well approximated by Eq. (4). The above approximation of $\tilde{g}^{(1)}$ by the "device function" $\tilde{g}^{(1)}_f$ is valid when $n_{k_{\|}}=n^{\mathrm{BE}}_{E=\hbar^2k_{\|}^2/2M_x}$ is nearly constant in the rather narrow energy interval $0\le E\le E^{(\alpha)}$, i.e., when $E^{(\alpha)}=(\hbar k^{(\alpha)}_{\|})^2/2M_x\ll k_BTe^{-T_0/T}$. For indirect excitons, this inequality with $T_0$ replaced by $T^*_0$ is definitely held for $n_{2d}\sim 10^{10}$ cm$^{-2}$ and $T\sim 1$ K (e.g., for $\alpha=20°$ the cutoff energy $E^{(\alpha)}$ is only 1.2 μeV). Thus the $k_{\|}$-filtering effect yields the correlation length $\xi_\gamma\simeq 2.215\sqrt{\varepsilon_b}/(k_0\sin\alpha)$, with $k_0\simeq 2.8\times 10^5$ cm$^{-1}$, according to Eq. (5). As a result, $\xi_\gamma$ is intrinsically scaled by the photon wavelength, i.e., is in the micron length scale [see Fig. 3(c), where $\xi_\gamma$ is plotted against the angle $\alpha$].

Compared to standard interference patterns in Young's double-slit experiment, with visibility contrast determined by $\tilde{g}^{(1)}$, the oscillatory behavior of the optical coherence function $\tilde{g}^{(1)}=\tilde{g}^{(1)}(r_{\parallel})$ is rather unusual [see Eq. (4) and Fig. 3(b)]. This is a signature of the $k_{\parallel}$-filtering effect: The $k_{\parallel}$-filtering function $G_f \propto \Theta(k_{\parallel}^{(\alpha)}-k_{\parallel})$ gives a sharp cutoff at $k_{\parallel}=k_{\parallel}^{(\alpha)}$ in the integrals of Eq. (3) that results in oscillations of $\tilde{g}^{(1)}(r_{\parallel})$. In some aspects, the effect is similar to Friedel oscillations in a Fermi liquid, with $\hbar k_{\parallel}^{(\alpha)}$ akin to the Fermi momentum.

*The coherence function* $\tilde{g}^{(1)}$ *of MC polaritons*. In this case, the "MC polariton $\rightarrow$ bulk photon" conversion function in Eq. (3) is $\Gamma_{x-\gamma}=\Psi(k_{\parallel})/\tau_{\gamma}(k_{\parallel})$, with $\Psi$ $(0\leq\Psi\leq1)$ the photon component along a MC polariton branch and $\tau_{\gamma}$ the radiative (escape) lifetime of a MC photon. In Fig. 4, $g^{(1)}=g_{\mathrm{MC}}^{(1)}(r_{\parallel})$ calculated with Eq. (6) for circularly polarized MC polaritons is compared with $\tilde{g}^{(1)}=\tilde{g}_{\mathrm{MC}}^{(1)}(r_{\parallel})$ evaluated with Eq. (3). According to the experiments,[2,3] we assume the BE distribution of MC polaritons along the lower polariton branch which is taken in the parabolic approximation with an effective in-plane mass $M_{\mathrm{MC}}^{lb}$. Compared to the case of QW excitons, the difference between $g_{\mathrm{MC}}^{(1)}$ and $\tilde{g}_{\mathrm{MC}}^{(1)}$ is much smaller, still giving $\xi_{\gamma}>\xi_{p}$. This is because the cutoff energy $E^{(\alpha)}$ in the $k_{\parallel}$-filtering effect is much larger than that relevant to QW excitons, due to $M_{\mathrm{MC}}^{lb}\ll M_x$. If $k_B T\ll E^{(\alpha)}\sim1$ meV, $g_{\mathrm{MC}}^{(1)}$ and $\tilde{g}_{\mathrm{MC}}^{(1)}$ nearly coincide (see Fig. 4).

We qualitatively explain a sharp increase in the coherence length with decreasing temperature, found in the experiments with coupled QWs,[4–7] by combining the $k_{\parallel}$-filtering effect with screening of disorder by dipole-dipole interacting indirect excitons.[17] In high-quality GaAs coupled QWs the screening process effectively develops at $T\lesssim5$ K, giving rise to a well-defined single-particle momentum $\hbar\mathbf{k}_{\parallel}$, as has been observed, e.g., in the experiments.[20,21] Thus the large correlation length $\xi=\xi_{\gamma}\sim1$ $\mu$m can naturally be explained by the $k_{\parallel}$-filtering effect and cannot be interpreted as a signature of BE condensation in a system of indirect excitons.

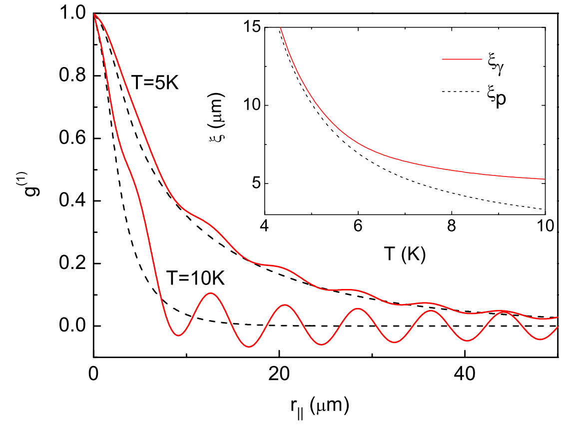


FIG. 4. (Color online) The MC polariton coherence function $g^{(1)}=g_{\mathrm{MC}}^{(1)}(r_{\parallel})$ (dashed lines) against that of emitted photons, $\tilde{g}^{(1)}=\tilde{g}_{\mathrm{MC}}^{(1)}(r_{\parallel})$ (solid lines). Inset: the coherence lengths $\xi_p$ and $\xi_{\gamma}$ versus temperature $T$. The calculations, which model the experiments (Ref. 3) refer to a GaAs microcavity with positive detuning $\delta=7$ meV and Rabi splitting $\Omega_{\mathrm{MC}}=4$ meV. The density of MC polaritons $n_{2d}=10^8$ cm$^{-2}$ and the aperture half-angle $\alpha=16.7°$, so that $T_0=27.6$ K and $E^{(\alpha)}=0.96$ meV.

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