Radiative Renormalizations for Excitonic Molecules

A. L. Ivanov and H. Haug

Institut für Theoretische Physik, J.W. Goethe Universität Frankfurt, Robert-Mayer-Strasse 8, D-60054 Frankfurt, Germany

(Received 26 August 1994)

The influence of the polariton effects on the internal structure of an excitonic molecule (xx) is examined by the solution of the new xx Schrödinger equation which includes the polariton effects self-consistently. In 3D (CuCl) the radiative renormalizations yield a nonparabolic xx dispersion with infinite effective xx mass at zero momentum; in quasi-2D (GaAs) they give rise to a "Mexican hat" structure in the xx dispersion at small momenta together with a considerable increase of both the xx binding energy ϵ^{xx} and the inverse xx radiative lifetime Γ^{xx}. In quasi-1D the influence of the polariton effects is so strong that the xx exists only as a broad resonance ($\epsilon^{xx} \sim \Gamma^{xx}$).

PACS numbers: 71.35.+z, 71.36.+c

Radiative corrections are well known in atomic physics due to the Lamb shift [1]. This shift of a transition frequency of an atom in the photon vacuum is due to spontaneous emission. It can be calculated within perturbation theory. We report here radiative renormalizations for an excitonic molecule in direct-gap semiconductors. These renormalizations originate from the polariton effect and cannot be included perturbationally. They give rise to the significant observable modifications of the xx dispersion for all dimensions: $D = 3$ (bulk semiconductors), $D = 2$ [quasi-2D quantum wells (QWL)] and $D = 1$ [quasi-1D quantum wires (QWR)].

Novel high-precision techniques with a spectral resolution ≤ 10 μeV have been applied to reinvestigate an xx in bulk CuCl [2]. Recently, the xx has been observed also in QWL [3,4]. The xx ground state has been calculated variationally within the underlying electron-hole (e-h) picture for 3D [5], 2D [6], and 1D [7] cases, respectively. The exact interparticle Coulomb potentials have been included, while the polariton effects have been neglected completely. The xx radiative decay is then treated perturbationally as a direct optical conversion $xx \to$ exciton (x) + photon (γ) with a giant oscillator strength [8]. The strong interaction with the electromagnetic field will, however, also modify the xx *dispersion* and give rise to an xx Lamb shift Δ^{xx}.

The e-h picture can be reduced to the x representation (see, e.g., [9,10]), if the x binding energy ϵ^x is much larger than that of molecule ϵ^{xx}. This approximation holds, e.g., for bulk CuCl ($\epsilon^x \simeq 190$ meV and $\epsilon^{xx} \simeq 34$ meV) and for GaAs QWL ($\epsilon^x \simeq 10$ meV and $\epsilon^{xx} \simeq 1-2$ meV). If one neglects the interaction with the light field, the xx eigenfunction $\Psi_0(p,K)$ and ground state energy Ω_K^{xx} are obtained from its Schrödinger equation in momentum representation (with $\hbar = 1$):

$$\sum_{\mathbf{p}'}\left[\sum_{\sigma=\pm} \omega^x(\sigma\mathbf{p} + \mathbf{K}/2)\delta_{\mathbf{p},\mathbf{p}'} + W_{xx}(\mathbf{p} - \mathbf{p}')\right]\Psi_0(p',K) = \Omega_K^{xx}\Psi_0(p,K). \quad (1)$$

Here, $\mathbf{p}$ and $\mathbf{K}$ are the momenta of the relative and center-of-mass motion of the xx, respectively; $\omega^x(p) = \omega_t + p^2/2M_x$ is the x kinetic energy, ω_t is the transverse x frequency, M_x is the x translational mass; $W_{xx}(p)$ is the attractive potential between two singlet x's with opposite $e(h)$ spin orientation. The attractive potential which follows from the $e(h)$-$e(h)$ Coulomb interactions includes both the direct and the exchange contributions. Although W_{xx} is known explicitly [9,10], it is often replaced by simpler model potentials such as the deuteron or the harmonic oscillator potential [2,9]. Because of the quadratic x dispersion the center-of-mass motion splits off in Eq. (1), i.e., $\Psi_0(p,K') = \delta(K - K')\Psi_0(p)$. The wave function of the relative motion $\Psi_0(p)$ and the corresponding xx binding energy $\epsilon_0^{xx} = 2\omega_t + K^2/2M_{xx} - \Omega_K^{xx}$ are independent of K. Here $M_{xx} = 2M_x$ is the xx translational mass.

In Ref. [10] we have shown that the polariton effects can be included self-consistently in the description of the xx in direct-gap semiconductors by the following wave equation:

$$\sum_{\mathbf{p}'}\left[\sum_{\sigma=\pm} \omega^{\mathrm{pol}}(\sigma\mathbf{p} + \mathbf{K}/2)\delta_{\mathbf{p},\mathbf{p}'} + \tilde{W}_{xx}(\mathbf{p},\mathbf{p}',\mathbf{K})\right]\tilde{\Psi}(\mathbf{p}',\mathbf{K}) = \tilde{\Omega}_K^{xx}\tilde{\Psi}(\mathbf{p},\mathbf{K}), \quad (2)$$

with $\omega^{\mathrm{pol}}(\mathbf{p}) = \omega^-(\mathbf{p})$, where $\omega^{\pm}(\mathbf{p})$ is the dispersion of the upper (+) and lower (−) polariton branch, respectively, which are given by the roots of $(\omega_{\mathbf{p}}^{\gamma})^2 = c^2p^2/\varepsilon_0 = \omega^2 + \Omega_c^2\omega^2/(\omega_t^2 + \omega_t p^2/M_x - \omega^2)$. Here, $\omega_{\mathbf{p}}^{\gamma}$ is the γ dispersion, ε_0 is the background optical dielectric constant, $\Omega_c = \sqrt{2\omega_{\mathrm{lt}}\omega_t}$, where ω_{lt} is the longitudinal-transverse splitting. The effective potential $\tilde{W}_{xx}$ is determined by the x components of the two interacting polaritons: $\tilde{W}_{xx}(\mathbf{p},\mathbf{p}',\mathbf{K}) = f(\mathbf{p},\mathbf{K})W_{xx}(\mathbf{p} - \mathbf{p}')$, where $f(\mathbf{p},\mathbf{K}) = \prod_{\sigma=\pm} \varphi^-(\sigma\mathbf{p} + \mathbf{K}/2, \omega^-(\sigma\mathbf{p} + \mathbf{K}/2))$, and $\varphi^{\pm}(p,\omega) = \omega^{\pm}(p)(\omega - \omega_p^{\gamma})/\omega[\omega^{\pm}(p) - \omega^{\mp}(p)]$. For the description of the "ground state" of the decaying xx, only the lower polariton dispersion branch has been taken

0031-9007/95/74(3)/438(4)$06.00

into account in Eq. (2). This equation contains simultaneously the x-x attraction as well as the x-γ polariton coupling. The x-weight functions $\varphi^{\pm}(p,\omega)$ satisfy the conditions $\varphi^{+}(p,\omega^{+}(p)) \geq 0$, $\varphi^{-}(p,\omega^{-}(p)) \geq 0$, and $\varphi^{+}(p,\omega^{+}(p)) + \varphi^{-}(p,\omega^{-}(p)) = 1$.

Equation (2) shows that the name "bipolariton" would be more appropriate for an xx than "biexciton." The most important consequence of Eq. (2) is that the xx is no longer a stable state as in Eq. (1). The *linear* integral Eq. (2) describes both the radiative decay of a xx state and the renormalization of the xx energy due to polariton effects. Because of $f(\mathbf{p},\mathbf{K})$ Eq. (2) cannot be transformed easily to real space. Equation (2) reduces to the xx Schrödinger equation (1), if both momenta p and p' are outside of the optical range ($p, p' \gg k_{\rm opt} = \omega_t\sqrt{\varepsilon_0}/c$). The xx radiative decay is described here as a continuous evolution of the xx internal state rather than a discrete optical conversion act to $x + \gamma$. When an x in the xx acquires a small momentum within the optical range ($p \leq k_{\rm opt}$), the xx undergoes a radiative annihilation. For the unperturbed x of Eq. (1), the optical range in momentum space corresponds to the x *turning points* in real space where the x group velocity vanishes: $v_g^x|_{p\to 0} = \partial\omega^x/\partial p \to 0$. In the polariton representation the x relative motion in real space is rather different. With increasing x distance and decreasing the momentum p, the group velocity increases $v_g^{\rm pol}|_{p\to 0} = \partial\omega^{\rm pol}/\partial p \to c/\sqrt{\varepsilon_0}$. Instead of returning, the x's start to leave each other; the radiative decay of the xx sets in.

The quasistationary solution $\tilde{\Psi}(\mathbf{p},\mathbf{K})$ of Eq. (2) represents an outgoing spherical wave due to the radiative decay of the xx state with $\tilde{\Omega}_K^{xx} = \Omega_K^{xx} + \Delta^{xx}(K) - i\Gamma^{xx}(K)/2$. Here, $\Omega_K^{xx} = \Omega_{K=0}^{xx} + K^2/2M_{xx}$ is the unperturbed xx energy. For the small relative momenta of the optical range, $\tilde{\Psi}(\mathbf{p},\mathbf{K})$ is so strongly modified that it cannot be calculated by perturbation theory.

Instead we decompose the xx wave function into

$$\tilde{\Psi}(\mathbf{p},\mathbf{K}) = \Psi_0(p) + \delta\Psi(\mathbf{p},\mathbf{K}), \tag{3}$$

where $\Psi_0(p)$ is the known solution of the xx Eq. (1), while $\delta\Psi(\mathbf{p},\mathbf{K})$ describes the outgoing part. The latter term is large only in the optical range, i.e., for $p \leq k_{\rm opt} \ll a_{xx}^{-1}$, where a_{xx} is the xx radius. For these small momenta one can approximate $\sum_{\mathbf{p}'} \tilde{W}_{xx}(\mathbf{p},\mathbf{p}',\mathbf{K})\delta\Psi(\mathbf{p}',\mathbf{K})$ by $\tilde{W}_{xx}(\mathbf{p},\mathbf{0},\mathbf{K})\sum_{\mathbf{p}'}\delta\Psi(\mathbf{p}',\mathbf{K})$, which yields

$$\delta\tilde{\Omega}(\mathbf{p},\mathbf{K})\delta\Psi(\mathbf{p},\mathbf{K}) - W_{xx}(p)f(\mathbf{p},\mathbf{K})\int\frac{d^Dp'}{(2\pi)^D}\delta\Psi(\mathbf{p}',\mathbf{K}) = -E(\mathbf{p},\mathbf{K})\Psi_0(p), \tag{4}$$

where

$$\delta\tilde{\Omega}(\mathbf{p},\mathbf{K}) = \tilde{\Omega}_K^{xx} - \sum_{\sigma=\pm}\omega^{-}(\sigma\mathbf{p} + \mathbf{K}/2) \tag{5}$$

is the energy difference between the bipolariton and the two separate polaritons. Furthermore,

$$E(\mathbf{p},\mathbf{K}) = \delta\tilde{\Omega}(\mathbf{p},\mathbf{K}) + (\epsilon_0^{xx} + p^2/M_x)f(\mathbf{p},\mathbf{K}). \tag{6}$$

Equation (4) is a Fredholm integral equation with separable kernel. The corresponding solution is given by

$$\delta\Psi(\mathbf{p},\mathbf{K}) = -\frac{E(\mathbf{p},\mathbf{K})\Psi_0(p) + CW_{xx}(p)f(\mathbf{p},\mathbf{K})}{\delta\tilde{\Omega}(\mathbf{p},\mathbf{K})}, \tag{7}$$

where

$$C = C(\tilde{\Omega}_K^{xx}) = A/(2\pi)^D(1 + B), \tag{8}$$

where $A = \sum_{\mathbf{p}}[E(\mathbf{p},\mathbf{K})/\delta\tilde{\Omega}(\mathbf{p},\mathbf{K})]\Psi_0(p)$, and $B = -\sum_{\mathbf{p}}[W_{xx}(p)/\delta\tilde{\Omega}(\mathbf{p},\mathbf{K})]f(\mathbf{p},\mathbf{K})$. According to Eq. (7), the roots $\mathbf{p}_0 = \mathbf{p}_0(\mathbf{K})$ of $\delta\tilde{\Omega}(\mathbf{p},\mathbf{K}) = 0$ give rise to singularities in $\delta\Psi(\mathbf{p},\mathbf{K})$ and $\tilde{\Psi}(\mathbf{p},\mathbf{K})$. These singularities describe the two outgoing polariton waves with $\mathbf{p}_0 + \mathbf{K}/2$ and $-\mathbf{p}_0 + \mathbf{K}/2$. However, the functions $\Psi_0(p)$, $\delta\Psi(\mathbf{p},\mathbf{K})$, and $\tilde{\Psi}(\mathbf{p},\mathbf{K})$ as well as the initial Eq. (2) are symmetric with respect to the substitution $\mathbf{p} \to -\mathbf{p}$. As a result, the functions $\delta\Psi(\mathbf{p},\mathbf{K})$ and $\tilde{\Psi}(\mathbf{p},\mathbf{K})$ can be decomposed into outgoing and incoming parts with equal contributions, e.g., $\tilde{\Psi} = \tilde{\Psi}^{\rm out}_{\Omega-i\Gamma/2} + \tilde{\Psi}^{\rm in}_{\Omega-i\Gamma/2}$. Thus the wave function includes not only the true outgoing part, but also the "unphysical" incoming component. Moreover, the residues of the roots of $\delta\tilde{\Omega}^{xx}(\mathbf{p},\mathbf{K}) = 0$ which correspond to the incoming waves and lie on the unphysical sheet of the complex p plane contribute to the integrals of Eq. (8). In order to avoid this difficulty, one can work only on the physical sheet [11] by treating the complex conjugate of Eq. (2). In this case $\tilde{\Omega}_K^{xx} \to [\tilde{\Omega}_K^{xx}]^*$, $\tilde{\Psi}(\mathbf{p},\mathbf{K}) \to \tilde{\Psi}(\mathbf{p},\mathbf{K})^* = \tilde{\Psi}^{\rm out}_{\Omega+i\Gamma/2} + \tilde{\Psi}^{\rm in}_{\Omega+i\Gamma/2}$, where $\tilde{\Psi}^{\rm out}_{\Omega-i\Gamma/2} = [\tilde{\Psi}^{\rm in}_{\Omega+i\Gamma/2}]^*$. In this way we treat the conjugate problem of resonant xx creation out of two polaritons instead of the xx optical decay. This procedure also avoids the task to normalize the outgoing wave function $\tilde{\Psi}^{\rm out}_{\Omega-i\Gamma/2}$.

In order to satisfy the boundary conditions, the polariton correction $\delta\Psi(\mathbf{p},\mathbf{K})$ has to vanish outside of the optical range, i.e., for $p \gg k_{\rm opt}$. From this requirement the complex energy eigenvalue $\tilde{\Omega}_K^{xx}$ will be calculated. With a model potential which satisfies the relation $W_{xx}(p) = \beta\Psi_0(p)$, where β is a negative constant, $\delta\Psi$ vanishes indeed outside the optical range, provided that

$$(\tilde{\Omega}_K^{xx})^* - \Omega_K^{xx} = \Delta^{xx} + i\Gamma^{xx}/2 = -\beta C^*(\tilde{\Omega}_K^{xx}). \tag{9}$$

The transcendent Eqs. (8) and (9) determine self-consistently both the radiative width Γ^{xx} and the xx Lamb shift Δ^{xx}. Because the polariton dispersion does not allow to split off the center-off-mass motion in Eq. (2), the total xx momentum K influences the relative motion of the x's. Thus $\delta\Psi$, Δ^{xx}, and Γ^{xx} are K dependent.

For the numerical evaluations of Eqs. (8) and (9) the deuteron potential $W_{xx}(p) = W_0^{(D)}/(p^2 + 1)^{(D+1)/2}$ has been used. As usual, the deuteron wave function for Eq. (1) is calculated variationally with the ansatz

$\Psi_0^{(D)}(p) = S^{(D)}/(p^2 + \alpha^2)^{(D+1)/2}$, where $S^{(D)}$ is the normalization constant and α the variational parameter. We use dimensionless units scaled with the 3D xx Rydberg and the 3D xx radius. The energy minimum is reached with $\alpha = f(W_0^{(D)})$. In order to satisfy the assumed proportionality between $W_{xx}(p)$ and $\Psi_0(p)$ we invert the result of the variation calculation and put $\alpha = 1$, i.e., $W_0^{(D)} = f^{-1}(\alpha = 1)$. By means of this procedure the following values are found $W_0^{(1)} = -18.0$, $W_0^{(2)} = -27\pi/2$, and $W_0^{(3)} = -54\pi$, respectively. With these values one gets further $S^{(1)} = 2.0$, $\epsilon_0^{xx,(1)} = 5.0$; $S^{(2)} = 2\sqrt{2\pi}$, $\epsilon_0^{xx,(2)} = 2.0$; $S^{(3)} = 8\sqrt{\pi}$, $\epsilon_0^{xx,(3)} = 1.0$. In comparison with experiments we use the unperturbed xx binding energy $\epsilon_0^{xx,(D)}$ as input parameter. While this procedure treats the x-x interaction in a simplified way, it treats the polariton effects exactly.

The fraction of the phase space in which the polariton effects dominate over the Coulombic x-x interaction is roughly of the order of κ^D, where $\kappa = k_{\mathrm{opt}} a_{xx} \ll 1$. Therefore, the influence of the polariton effects on the xx internal state increases drastically if one reduces the dimensionality from $D = 3$ to $D = 1$. We start with 2D, because this case contains all the features in the most explicit form.

(2D) Quasi-2D GaAs QWL.—The quasi-2D xx is supposed to originate from the two Y-mode surface polaritons which are transverse with $\mathbf{p}$ in-plane and $\Omega_c^2 = R_c\sqrt{p^2 - \omega^2\varepsilon_0/c^2}$ [12]. In Fig. 1 the calculated xx energy $\Delta\tilde{\Omega}_K^{xx} = \tilde{\Omega}_K^{xx} - 2\omega_t = -\epsilon_0^{xx} + \Delta^{xx}(K) + K^2/2M_{xx}$ (solid line), the unperturbed energy $\Delta\Omega_K^{xx} = -\epsilon_0^{xx} + K^2/2M_{xx}$ (dashed line), and the radiative half-width $\Gamma^{xx}(K)/2$ (dash-dotted line) are plotted versus K. The following parameters have been used: $\omega_t = 1.592$ eV, $R_c = 30.7 \times 10^{-2}$ eV^2 Å, $M_x = 0.44m_0$, $\varepsilon_0 = 12.9$, $\epsilon_0^{xx} = 2$ meV, and $a_{xx} = 131.6$ Å. The most interesting feature of the renormalized xx dispersion $\tilde{\Omega}_K^{xx}$ is the "Mexican hat" structure at $K = 0$. The strong spatial dispersion of $\Delta^{xx}(K)$ is shown in the inset of Fig. 1. The radiative renormalization results in a negative effective xx mass M_{xx}^{eff} at $K = 0$ ($M_{xx}^{\mathrm{eff}} \simeq -0.007m_0$ for Fig. 1). In order to explain this result, one has to examine the joint density of polariton states (JDPS) given by the function $\delta\tilde{\Omega}(\mathbf{p},\mathbf{K})$, i.e., by the denominators of the integrands of Eq. (7). For $\Gamma^{xx} \to 0$ the point $\mathbf{K}_0$ (see Fig. 1) is a critical saddle point of the JDPS for which $\partial\delta\tilde{\Omega}(\mathbf{p},\mathbf{K})/\partial p_i = 0$ with $i = x, y$ (see Fig. 2). This critical point corresponds to the degenerate two-polariton (with $\mathbf{p} = \mathbf{K}_0/2$) xx absorption. In 2D it gives rise to a logarithmic *van Hove* singularity [13] in the JDPS. The critical points lie on a circle with the radius $K_0 \simeq 6.03 \times 10^5$ cm^{-1}. For $K < K_0$ the solutions of $\delta\tilde{\Omega}(\mathbf{p},\mathbf{K}) = 0$ lie on a topologically connected contour; for $K > K_0$ they are located on topologically disconnected contours. The critical saddle point $\mathbf{K}_0$ is a touching marginal point (see Fig. 2). In the vicinity of $\mathbf{K}_0$ the xx couples effectively to surface electromagnetic field resulting in a strong decrease of the xx energy together with a sharp increase of $\Gamma^{xx}(K)$ around K_0. The minimum of the xx dispersion at $K_{\min} \simeq 4.9 \times 10^5$ cm^{-1} is shifted from K_0 to a smaller value partly because of the competition between the positive kinetic energy $K^2/2M_{xx}$ and the negative Lamb shift $\Delta^{xx}(K)$.

The Mexican hat structure around $K = 0$ results in a considerable enhancement of the effective xx binding energy for K values of the optical range. Starting with $\epsilon_0^{xx} = 2$ meV (the upper limit of the unperturbed xx binding energy in GaAs QWL according to the variational calculations [6]) one gets $\tilde{\Omega}_K^{xx} - 2\omega_t = -3.27$ meV for

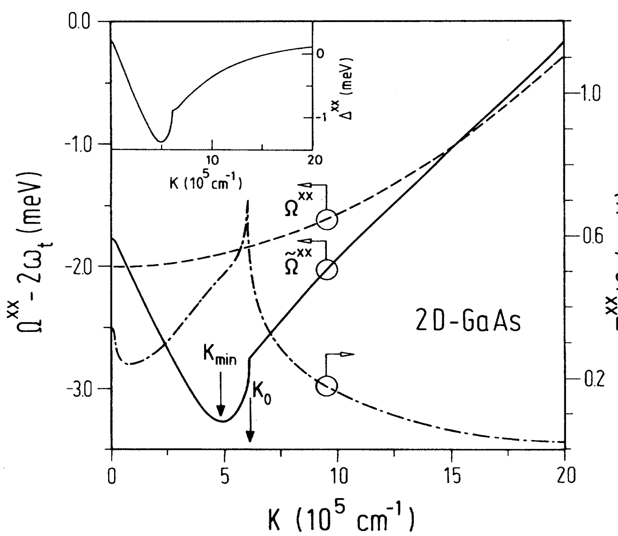

FIG. 1. Quasi-2D GaAs QWL: Renormalized (solid line) and unperturbed (dashed line) xx dispersions together with xx radiative half-width $\Gamma^{xx}(K)/2$ (dash-dotted line). Inset: xx Lamb shift Δ^{xx} versus momentum K.

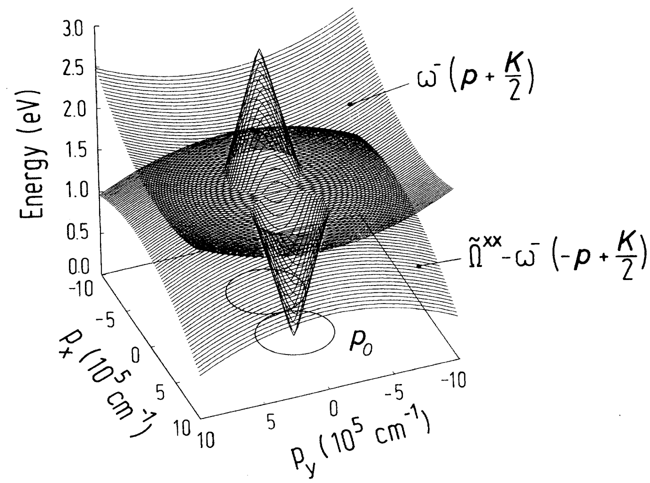

FIG. 2. Graphic solution of the energy-momentum conservation $\Delta\tilde{\Omega}(\mathbf{p},\mathbf{K}) = 0$ of the xx optical decay (resonant creation) into two polaritons for $K = K_0$.

$K = K_{\min}$. The latest optical experiments [4] indicate $\epsilon^{xx} \simeq 2.7$ meV in contradiction to the previous observations [3] ($\epsilon^{xx} \simeq 1$ meV) and cannot be explained without xx radiative renormalizations.

(3D) Bulk CuCl.—In Fig. 3 the corresponding 3D results are shown for the CuCl parameters: $\omega_t = 3.203$ eV, $\omega_{\mathrm{lt}} = 5.7$ meV, $M_x = 2.6m_0$, $\varepsilon_0 = 5.6$, $\epsilon_0^{xx} = 34$ meV. These values give $\Omega_c = 191.1$ meV and $a_{xx} = 9.3$ Å. Two important features can be distinguished.

(i) Strong renormalization of the xx dispersion around $K = 0$ ($K \leq K_c \simeq 1.5 \times 10^5$ cm^{-1}). This range is shown in detail as the inset of Fig. 3. There is a very shallow minimum in the xx dispersion at $K_{\min} \simeq 0.8 \times 10^5$ cm^{-1} with the depth of about 0.01 μeV (in 3D K space the minimum lies on a sphere with the radius $K_{\min}$). This value is considerably smaller than the corresponding $\Gamma^{xx}(K_{\min})$. Thus, one obtains a nearly "horizontal" xx dispersion with $M_{xx}^{\mathrm{eff}} \rightarrow \infty$ for $K \leq K_c$. This xx dispersion will strongly oppose a possible xx Bose-Einstein condensation in $K = 0$, which has been searched for in the last two decades [14].

(ii) For $K > K_c$ the xx Lamb shift results in an effective renormalization of the "bare" xx mass M_{xx}. For this range we estimate $M_{xx}^{\mathrm{eff}} \simeq 2.15M_x$. Such a modification of M_{xx} has indeed been observed experimentally [15] ($M_{xx}^{\mathrm{eff}} \simeq 2.3M_x$) and has not been understood up to now. The K dependence of $\Gamma^{xx}(K)/2$ is in good agreement with the recent measurements [2]. For bulk CuCl, the van Hove feature at $K_0 \simeq 8.85 \times 10^5$ cm^{-1} is classified as the S_2 critical point (saddle point of second type in notations of [13]), its influence on $\Delta^{xx}(K)$ is considerably weaker than in 2D.

(1D) Quasi-1D GaAs QWR.—For 1D the influence of the polariton effects on a xx state is most pronounced.

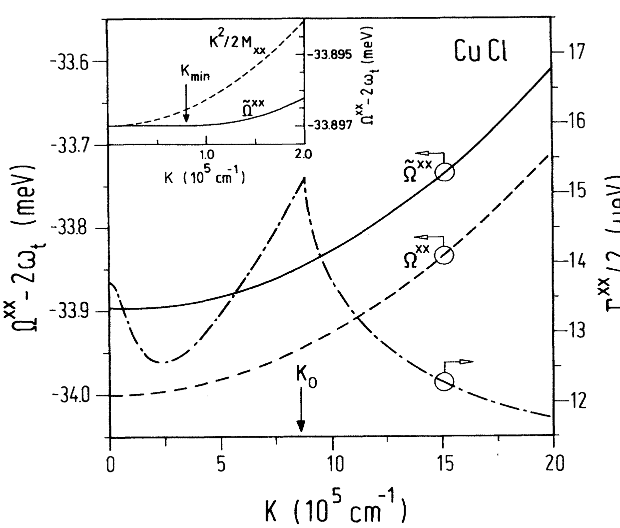

FIG. 3. Bulk CuCl: Renormalized (solid line) and unperturbed (dashed line) xx dispersions, and radiative half-width $\Gamma^{xx}(K)/2$ (dash-dotted line). Inset: magnified region around $K = 0$. The unperturbed dispersion $K^2/2M_{xx}$ is shifted for direct comparison with xx renormalized energy $\Delta\tilde{\Omega}_K$.

For $\epsilon_0^{xx} < 30$ meV the numerical estimates give positive $\Delta^{xx}(K) > \epsilon_0^{xx}$, so that no xx bound state exists because $\Delta\tilde{\Omega}_K^{xx} = \tilde{\Omega}_K^{xx} - 2\omega_t > 0$. For unperturbed $\epsilon_0^{xx} \geq$ 30 meV, we receive a xx bound state with $\Delta^{xx}(K) < 0$ and a camel-back structure at $K = 0$. However, the very large $\Gamma^{xx}(K) \sim \Delta\tilde{\Omega}_K^{xx}$ shows that the xx bound state exists only as a broad resonance. In 1D the x inside a xx cannot make a complete oscillation because it cannot traverse the bottleneck region of the lower polariton branch. The x motion can also not involve the upper polariton branch, because $\Omega_c \gg \epsilon_0^{xx}$. As a result, the xx decays optically very quickly, before a complete oscillation period $1/\epsilon_0^{xx}$. For 1D we have to conclude the absence of a well-defined xx bound state. The influence of the polariton effects is so strong that $\Omega_{c\kappa} \sim \epsilon_0^{xx}$. This result may explain why an xx in quasi-1D GaAs QWR has not been observed, although the variational calculations for Eq. (1) give a very large value $\epsilon_0^{xx} \leq 40$ meV [7].

We appreciate valuable discussions with L. V. Keldysh and Y. Toyozawa. This work has been supported by the Volkswagen Stiftung.

[1] W.E. Lamb and R.C. Retherford, Phys. Rev. **72**, 241 (1947).
[2] H. Akiyama, M. Kuwata, T. Kuga, and M. Matsuoka, Phys. Rev. B **39**, 12 973 (1989); M. Hasuo, N. Nagasawa, and T. Itoh, Opt. Commun. **85**, 219 (1991).
[3] R.C. Miller, D.A. Kleinmann, A.C. Gossard, and O. Munteanu, Phys. Rev. B **25**, 6545 (1982); D.C. Reynolds, K.K. Bajaj, C.E. Stutz, R.L. Jones, W.M. Theis, P.W. Yu, and K.R. Evans, Phys. Rev. B **40**, 3340 (1989).
[4] S. Bar-Ad and I. Bar-Joseph, Phys. Rev. Lett. **68**, 349 (1992).
[5] O. Akimoto and E. Hanamura, J. Phys. Soc. Jpn. **33**, 1537 (1972); W.F. Brinkman, T.M. Rice, and B. Bell, Phys. Rev. B **8**, 1570 (1973).
[6] D.A. Kleinman, Phys. Rev. B **28**, 871 (1983).
[7] L. Banyai, I. Galbraith, C. Ell, and H. Haug, Phys. Rev. B **36**, 6099 (1987).
[8] E. Hanamura, Solid State Commun. **12**, 951 (1973).
[9] M.I. Sheboul and W. Ekardt, Phys. Status Solidi B **73**, 165 (1976); H. Haug and S. Schmitt-Rink, Progr. Quantum Electr. **9**, 3 (1984).
[10] A.L. Ivanov and H. Haug, Phys. Rev. B **48**, 1490 (1993).
[11] L.D. Landau and E.M. Lifshitz, *Course of Theoretical Physics* (Pergamon, New York, 1965), Vol. 3 , Sec. 132.
[12] M. Nakayama, Solid State Commun. **55**, 1053 (1985).
[13] L. van Hove, Phys. Rev. **89**, 1189 (1953).
[14] L.L. Chase, N. Peyghambarian, G. Gryndberg, and A. Mysyrowicz, Phys. Rev. Lett. **42**, 1231 (1979); M. Hasuo, N. Nagasawa, T. Itoh, and A. Mysyrowicz, Phys. Rev. Lett. **70**, 1303 (1993).
[15] M. Ueta, H. Kanzaki, K. Kobayashi, Y. Toyozawa, and E. Hanamura, *Excitonic Processes in Solids,* Springer Series in Solid-State Sciences Vol. 60 (Springer, Berlin, 1986), Chap. 3.