Effect of nonuniform continuum density of states on a Fano resonance in semiconductor quantum wells

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We experimentally study the excitonic Fano resonance in $GaAs/Al_{0.3}Ga_{0.7}As$ multiple quantum wells under an external electric field. The resonance results from exciton transitions via localized states and continuum minibands of Wannier-Stark states. In photoreflectance spectra the line shape displays anomalous behavior as a function of increasing field. Our theoretical analysis reveals that this effect, which is absent for the conventional Fano resonance, arises from the field dependence of the nonuniform density of continuum states.

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I. INTRODUCTION

The Fano resonance (FR) is due to quantum-mechanical coupling between a discrete energy state and a degenerate energy continuum. Its line shape, known as Fano line shape, has the asymmetric form

$$I(\epsilon) = I_o \frac{(q+\epsilon)^2}{1+\epsilon^2}, \tag{1}$$

where I_o is a fictitious "background" spectrum resulting from the continuum states, q is the Fano shape parameter determining the asymmetry of the line shape, and $\epsilon=2(\omega-E_r)/\Gamma$ measures the energy difference from the resonance center E_r normalized by the half width of the level broadening Γ of the discrete state due to coupling to the continuum states.[1] The shape parameter q is understood[2] as a ratio between the rates of two transition pathways: one for the direct transition to the discrete state and the other for the two-step transition to the discrete state via the continuum [see, e.g., Eq. (6)]. The FR appears ubiquitously in a variety of systems including atoms, molecules, semiconductors, and even optical resonators.[3–12] Recently, direct control of the FR by tuning the resonance center E_r or the Fano coupling between the discrete state and the continuum has been reported.[9,10] In this paper, we report another aspect of controllable FR tuning whereby a bias field modifies the continuum states.

Importantly, coherent control of quantum states is a key step for quantum information processing.[13] For example, in optical experiments on atomic systems, the interference between resonant and nonresonant transitions of energy levels can be controlled using spectrally encoded broadband laser pulses.[14–16] Analogously, in semiconductors a similar form of coherent control of optical transitions toward high-performance quantum devices is enabled if the densities of the continuum energy states can be tuned. In conjunction with the interference in FR, this offers an alternative approach to use quantum-confined Stark effect for quantum information devices.[17,18]

In semiconductor quantum wells, resonant exciton transitions generate a FR via the interference between the two channels, a discrete state of one subband of an electron-hole pair and energetically degenerate exciton continua pertaining to other subbands.[19,20] Among the quantum wells, $GaAs/Al_xGa_{1-x}As$ superlattices are interesting, as the continuum states contributing to the FR are supported by the exciton minibands of the low-lying Wannier-Stark ladder (WSL) of the superlattice.[21–23] The continuum states are activated and controlled by an external Stark electric field. Therefore the Fano coupling parameter Γ and the Fano shape parameter q may strongly depend on the field. Indeed, in a recent experiment,[9] a monotonic decrease in Γ was observed as the Stark field increases. This behavior is attributed to the following fact:[9] as the field becomes stronger, the WSL states become more localized, and the momentum mismatch between the discrete exciton and the WSL continuum increases at higher momentum, resulting in the reduction in the coupling strength (the coupling matrix element) between them.

The broadening of the exciton transition described by Γ depends not only on the coupling matrix element but also on the density of states (DOS) of the continuum miniband or the spectral shape and width of the continuum states. Therefore, the Fano profile is expected to depend on the DOS of the continuum state, as depicted in Fig. 1, especially when the DOS is narrow or nonuniform near the resonance. The control of the FR by directly and continuously changing the DOS in quantum wells, by means of field biasing or carrier injection, may open a new way to control matter-light interaction. Alternatively, the tunable FR may be useful to read out the shape of the DOS. In previous studies, however, the effect of DOS has been ignored,[1] where either a flat and wide DOS is assumed, or studied in an indirect way using, for example, several bulk-semiconductor samples with different DOS (Ref. 5) (i.e., the DOS is fixed by the dopant concentration).

In this paper, we report an experimental observation of the FR, modified by nonuniform DOS of the WSL continuum states in a $GaAs/Al_{0.3}Ga_{0.7}As$ superlattice biased by a Stark field. We observe an anomalous dependence of the Fano coupling parameter Γ on the field strength. In particular, we observe an increase in Γ which contrasts with the expected monotonic decrease in Γ with increasing field strength.[9] The discrepancy is attributed to the effect of the nonuniform DOS

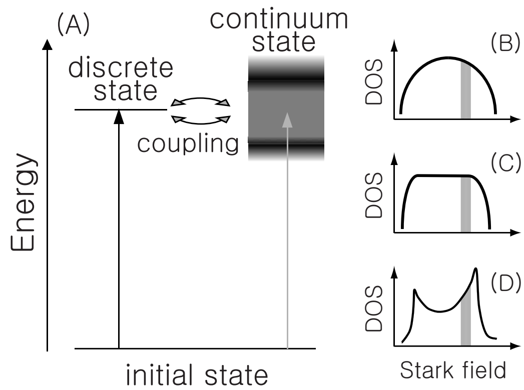

FIG. 1. (a) Schematic exciton energy diagram of a localized state and a Wannier-Stark miniband in a quantum-well superlattice under an Stark field. The localized state couples to the miniband states. As the field strength varies over a wide range, the miniband increases in energy and splits [see Eq. (2)]. Therefore, the density of the miniband states can be significantly modified near the energy of the localized state as, e.g., in (b)–(d). The change in the DOS affects the Fano resonance.

of the WSL continuum states, which can be tuned by the Stark field. We theoretically support this interpretation, by extending the original Fano formula (1) to include the dependence of the DOS on the energy of the continuum and the Stark field.

II. EXPERIMENT

We investigate a molecular-beam-epitaxy-grown GaAs/$Al_{0.3}Ga_{0.7}As$ superlattice with 35 GaAs quantum wells of 97 Å thickness (34 monolayers) separated by 17-Å-thick (six monolayers) $Al_{0.3}Ga_{0.7}As$ barriers, clad on both side by 2500 Å $Al_{0.3}Ga_{0.7}As$ buffer layers, *p*- and *n*-doped GaAs layers (for Schottky contact), and metallic contacts, respectively. The variation in the quantum-well thickness is less than one monolayer. This superlattice is similar to the structure previously studied for Bloch oscillations in Ref. 24. It has a width of 19 meV for the lowest electron miniband and 2 meV for the heavy-hole miniband. Photoreflectance measurements are carried out at 4 K using a Fourier transform infrared spectrometer.[25]

In the absence of the Coulomb interaction, we know how the WSL states evolve from a continuum into resonant states as the bias Stark field increases. Assuming that the coupling does not occur among the WSL states with different miniband indices, the field dependence of the single-particle WSL energy is given by Ref. 26 as

$$E_n = E_0 + neFd \quad (n = 0, \pm 1, \pm 2, \cdots), \tag{2}$$

where n is the index of the quantum wells of the superlattice and d (=114 Å) is the superlattice period. This formula is applicable to our case as the first miniband gap energy (≈60 meV) is larger than the maximum energy shift due to the bias field (≈28.9 meV). Every parameter discussed hereafter implies the first Wannier-Stark ladder index.

In the presence of the Coulomb interaction, a bound exciton state is formed below the lower miniband edge at zero bias field. The continuum miniband has the width of 19 meV, estimated from the Kronig-Penny model,[27] and is located 4.9 meV above the bound exciton state. As the Stark field increases, the energy states in WSL split into continuum subbands, following Eq. (2). As a result, the DOS of the WSL states has a structure of resonance peaks[28] (see also Fig. 4). At the same time, the bound exciton state sequentially couples with the lower-lying WSL states in the neighboring quantum wells. This coupling mechanism leads to asymmetric FR.[20]

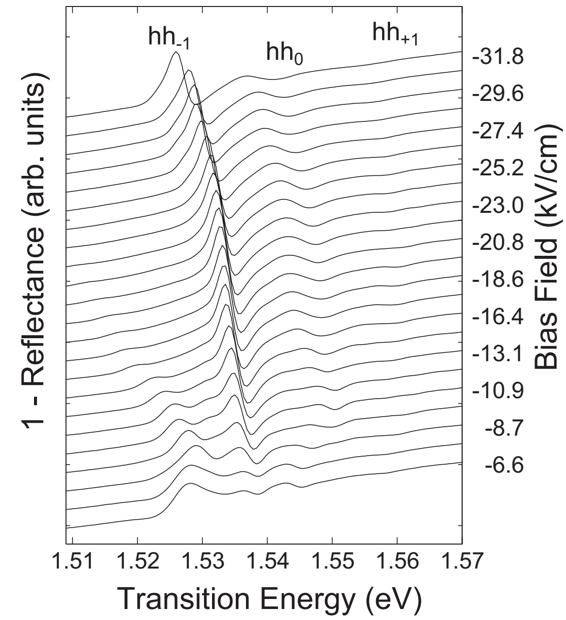

FIG. 2. Excitonic absorption spectra of a 97/17 Å superlattice for different bias fields, measured in reflection.

The absorption spectra in Fig. 2 show the FR discussed above. Below the onset bias field (≈−10.9 kV/cm), where the WSL starts to develop (not as a meaning of onset bias field where exciton state starts to meet continuum band), three delocalized exciton lines in the superlattice miniband regime are identified. They are associated with the fundamental heavy-hole exciton (hh) and the light-hole exciton (lh) followed by the excited heavy-hole exciton, as previously observed.[24,29] Above the onset, the FR peak is the most prominent at the $n=-1,0$ transitions (hh_{-1} and hh_0). This is because the electronic portion of the $n=-1,0$ exciton states maximally overlap with the heavy-hole state localized at the $n=0$ quantum well, resulting in the strongest absorption strength.[21,23]

As in Fig. 2, our experiment does not show the zero-field limit, where, if the data at the bottom in Fig. 2 were measured (exactly at zero bias field), it should reveal a broad plateau due to a continuum miniband in a high-energy side. The first peak from the left in Fig. 2 originates from a sharp miniband exciton peak at zero field. It looks rather broadened because the spectral resolution is not enough to decompose the peak into subpeaks due to excitonic transition to different quantum wells in the low-field regime,[23] exactly the same superlattice structure with ours was used, and a similar spectrum pattern was observed in the low-field regime. The second and third peaks (not a plateau) in Fig. 2 are remnants

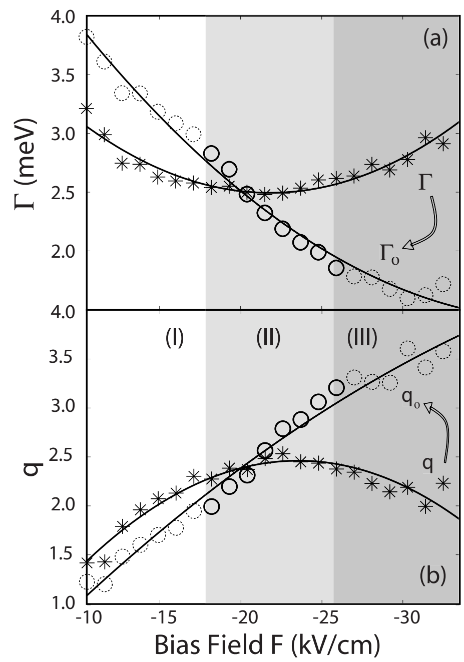

FIG. 3. The Fano coupling parameter Γ and the Fano shape parameter q (asterisks), extracted from the spectra of the hh_{-1} transition in Fig. 2 and by using Eq. (1), are drawn as a function of the bias field. For comparison, the bare parameters Γ_o and q_o are extracted (circles) by using Eq. (11), in which the effect of the dependence (assumed to be parabolic) of the density of the continuum states on the field is removed.

of so-called Franz-Keldysh (FK) oscillations in the low-field regime. The FK oscillations occur because the effective-mass motion at the miniband edges interfere with the WSL which starts at the miniband center.[28] Our experiment mainly focused on the regime above the onset field where the resonance shape is governed by Fano coupling mechanism which describes the interference between a discrete channel due to the bound exciton and a continuum channel due to the WSL continuum.

III. RESULTS AND ANALYSIS

From the spectra of the hh_{-1} transition and by using the line shape in Eq. (1), we extract the Fano parameters Γ and q (see Fig. 3).[30] At small bias fields, the extracted value of Γ (q) becomes smaller (larger), as the field increases, following the conventional behavior;[9] for this weak-field regime, the field dependence of the line shape of the Fano resonance was studied in Ref. 31. However, as the bias field is increased above approximately −21.6 kV/cm, the behavior of Γ and q deviates from the conventional behavior.

The unexpected response of Γ and q can be attributed to the Stark-field dependence of the DOS of the WSL continuum states; the field dependence modifies the Fano parameters, which we will derive below. The DOS has a structure of resonance peaks because of the field-induced splitting of the continuum band into the WSL minibands [see Eq. (2)]; each peak is associated with the WSL states centered at a quantum well. The peak height is determined by the overlap

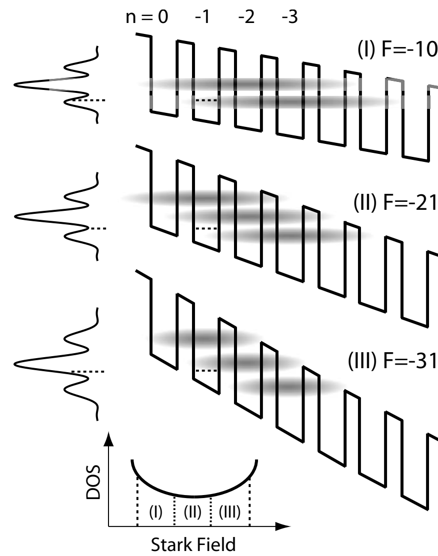

FIG. 4. Schematic view of the energy dependence of the density of the continuum WSL states (left panel) and spatial extension of the lower-lying three WSL states in the quantum-well superlattice (right) at different Stark-field strengths. The bound exciton state localized in the $n=-1$ well is depicted as dashed lines. This figure shows that near the energy of the bound exciton, the density of the continuum states varies with the field. Between the resonant couplings with the next nearest-neighboring WSL state in (I) and with the nearest-neighboring WSL state in (III), the exciton state resonantly sweeps through the energy interval, having the minimal coupling somewhere in between as in (II). Inset: The dependence of the density of the WSL states, coupled to the bound exciton, on the field strength.

between the states in the miniband and the heavy hole localized in the $n=0$ quantum well; the DOS is called as optical DOS in the literature,[28] as it reflects the optical transition from the localized heavy hole to the continuum states. As the energy of the WSL states (relative to the energy of the bound exciton) and the overlap between the WSL states and the heavy hole depend on the Stark field, the DOS is not constant but varies depending on the Stark field as depicted in the left panel of Fig. 4. In the operating range of the Stark field from −10.9 to −31.8 kV/cm, the bound exciton state passes through the valley between two peaks of the continuum DOS, associated with the WSL states with $n=-3$ and those with $n=-2$, respectively. For example, at the Stark field of −10.9 kV/cm, the hh_{-1} bound exciton couples with the WSL states centered at the $n=-3$ quantum well while it couples with those at the $n=-2$ quantum well at −31.8 kV/cm [see Figs. 4(I) and 4(III)]. As a result, the DOS has a parabolic dependence on the Stark field (see the inset of Fig. 4).

Next, we theoretically derive Γ and q, taking account of the dependence of the DOS on the bias field. We start with a model consisting of one discrete bound exciton state, $|\phi\rangle$, and a continuum of WSL states, $|\psi_\omega\rangle$, both of which are optically excited from an initial state $|i\rangle$ in a valence band. The DOS $\rho(\omega,F)$ of the continuum depends on the bias Stark field F and on the transition energy ω in general. The effective Hamiltonian describing electrons in the conduction band is written as

$$H = E_r|\phi\rangle\langle\phi| + \sum_\omega \omega|\psi_\omega\rangle\langle\psi_\omega| + \sum_\omega (V_F|\psi_\omega\rangle\langle\phi| + \text{H.c.}). \tag{3}$$

Here V_F is assumed to be independent of ω and is chosen to be real as the time-reversal symmetry is preserved. The absorption spectra is obtained[32] as

$$I(\omega) = -(1/\pi)\text{Im}\langle i|\hat{T}(\omega - H + i0^+)^{-1}\hat{T}|i\rangle, \tag{4}$$

where $\hat{T}$ is the transition operator from the valence band to the conduction band. After straightforward calculations, one arrives at the Fano formula (1) for the absorption spectra.

The Fano parameters Γ and q now depend on ω and F,

$$\Gamma(\omega,F) = 2\pi\rho(\omega,F)V_F^2, \tag{5}$$

$$q(\omega,F) = \frac{1}{\pi\rho(\omega,F)V_F}\frac{\langle\Phi|\hat{T}|i\rangle}{\langle\psi_\omega|\hat{T}|i\rangle}, \tag{6}$$

where

$$|\Phi\rangle = |\phi\rangle + \mathcal{P}\int d\omega'\rho(\omega',F)\frac{V_F}{\omega-\omega'}|\psi_{\omega'}\rangle \tag{7}$$

is the state modified from $|\phi\rangle$ by the coupling with the continuum $|\psi_\omega\rangle$, and $\mathcal{P}$ represents the principal value.

The dependence of ρ on ω and F can be approximately obtained for regime II in Fig. 4. In the regime $j \in \{\text{I},\text{II},\text{III}\}$, the dependence on ω may be simplified around the energy E_r of the bound exciton as

$$\rho(\omega,F) \simeq \alpha_{F,j}(\omega - E_r) + \beta_{F,j} + O(\omega^2), \tag{8}$$

where $\alpha_{F,j}$ and $\beta_{F,j}$ depend on F. As shown in Fig. 4, in regimes I and III, the peak structure of the DOS causes $\alpha_{F,\text{I}} < 0$ and $\alpha_{F,\text{III}} > 0$ while in regime II, the valley of the DOS can be approximated as $\alpha_{F,\text{II}} \simeq 0$. On the other hand, the parabolic dependence of ρ on the field strength F results as

$$\beta_{F,\text{II}} \sim (F - F_0)^2 + \beta_o, \tag{9}$$

where F_0 is the field strength at which the DOS has its minimum value of $\beta_o > 0$. One can determine $\rho(\omega,F)$ for regimes I and III in a similar way.

One may notice that the dependence of ρ on ω modifies the Fano line shape in general. When $\rho(\omega,F)$ depends linearly on ω as in Eq. (8), the form (1) of the Fano line shape is still preserved but with redefined $\Gamma(F)$ and $q(F)$; the ω dependence of Γ and q is absorbed in the form (1). Therefore, the extraction of $\Gamma(F)$ and $q(F)$ in Fig. 2, from the data in Fig. 1 and by using Eq. (1), may be valid. By substituting Eqs. (5)–(8) into Eq. (1), we find $\Gamma(F)$ and $q(F)$ in regime II,

$$\Gamma(F) \simeq 2\pi\beta_{F,\text{II}}V_F^2, \quad q(F) \simeq \frac{c}{\pi\beta_{F,\text{II}}V_F}, \tag{10}$$

where the dependence of $c \equiv \langle\phi|\hat{T}|i\rangle/\langle\psi_\omega|\hat{T}|i\rangle$ on ω and F may be ignored. The dependence of Γ and q on the field strength F is governed not only by V_F but also by $\beta_{F,\text{II}}$ which comes from the parabolic dependence of the DOS on F. This behavior can explain the result in Fig. 3.

For further analysis, we introduce the "bare" Fano parameters, $\Gamma_o(F) \equiv 2\pi\beta_o V_F^2$ and $q_o(F) \equiv c(\pi\beta_o V_F)^{-1}$. In regime II, we can extract Γ_o and q_o from the experimental result of Γ and q in Fig. 3, by using the relations

$$\Gamma_o(F) = \frac{\beta_o}{\beta_F}\Gamma(F), \quad q_o(F) = \frac{\beta_F}{\beta_o}q(F). \tag{11}$$

Then, as shown in Fig. 3, the bare Fano parameters follow the conventional monotonous behavior,[9] i.e., the conventional behavior is retrieved. In the bare parameters, the DOS contribution is removed and its dependence on F is determined only by V_F, which decreases for stronger F because of the momentum mismatch and the localization of the WSL states.

In summary, the behavior of the Fano coupling and shape parameters in Fig. 3 is, therefore, related to the shape of DOS of WSL states. If the DOS effect is neglected, the Fano coupling decreases monotonically as the bias field increases, because of the two reasons: the axial localization of the wave function of the WSL state and the increasing momentum mismatch of energy transitions, which explains the (I) and (II) regimes of our experiments, as in Fig. 4 and Ref. 9. We notice that the 67/17 Å superlattice used in Ref. 9 has a larger bandwidth (43 meV) of the continuum miniband than in our experiment. In Ref. 9, the bound exciton level does not seem to bypass the minimum of the parabolic DOS (regime II), therefore, the effect of the nonuniform DOS on the FR may not be clearly identified. However, in the high-field regime (−21.9 kV/cm∼−31.8 kV/cm), the increase in the DOS of the WSL state can be more dominant than the decrease in the coupling matrix element so that the Fano coupling increases as in Fig. 3(III).

IV. CONCLUSION

In conclusion, we report the experimental observation of anomalous behavior of the FR parameters Γ and q in a biased semiconductor superlattice. We attribute this observation to the effect of the nonuniform shape of the density of the continuum WSL states, which is caused by resonant sweeping of an exciton bound state in between two neighboring extended WSL states.

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