PHYSICAL REVIEW B **66**, 081105(R) (2002)

Frequency-dependent electrical transport under intense terahertz radiation

C. Zhang
Department of Physics and Institute for Superconducting and Electronic Materials, University of Wollongong, New South Wales 2522, Australia

(Received 20 May 2002; published 29 August 2002)

By using a formalism based on quantum transport equation, we study the current response of an electron gas strongly coupled to an intense terahertz laser. The effect of laser is treated exactly, which gives rise to the electron-photon sidebands. Our formalism includes up to the second order in electron-impurity scattering and to any order in optical transitions between the sidebands. Resonant scattering can occur when the frequency of the probing field equals a multiple of the terahertz laser frequency. The usual linear response theory is the zero-photon term in our formalism. This result can be applied to study the electrical transport in strongly coupled electron-photon systems. One such system is a two-dimensional semiconductor structure under an intense terahertz radiation.

DOI: 10.1103/PhysRevB.66.081105 PACS number(s): 71.10.Ca, 71.45.Gm, 78.30.Fs

Significant progress has been made in recent years in the terahertz phenomena in electronic materials. This is mainly due to the rapid development of high-power, long-wavelength, and tunable laser sources such as free-electron lasers. Recently, a new mechanism for population inversion in semiconductors under an applied electric field was predicted[1] and a widely tunable continuous-wave terahertz generation was achieved experimentally in strained Ge.[2] These radiation sources can provide linearly polarized laser radiation in the terahertz regime.[3–10] Terahertz lasers have been applied to the experimental investigation of nonlinear transport and optical properties in electron gases such as low-dimensional semiconductor systems. Many interesting terahertz phenomena have been investigated, including resonant absorption,[3] photon enhanced hot-electron effect,[6] terahertz photon-induced impact ionization,[7] LO-phonon bottleneck effect,[8] terahertz photon assisted tunneling,[9] terahertz cyclotron resonance,[10] terahertz switching effects in tunneling diodes.[11–13] The conductivity changes due to an ac field radiation were calculated using a perturbative method.[14]

Despite the rapid development of terahertz phenomena, a theoretical formalism describing the quantum transport in strongly coupled electron-photon systems is lacking. In this paper, we present a theoretical investigation of the electrical current response of such an electron-photon system to a probing field. We calculate the density matrix of the electron-photon system and its contribution to the electrical current. The total current consists of contributions of successively electron-photon sidebands. The transport scattering time exhibits the resonant scattering of the probing electromagnetic field due the transitions between electron-photon sidebands.

Let us consider a two-dimensional electron gas under an intense laser radiation. We choose the laser field to be along the x direction, $\mathbf{E}_\gamma(t)=E_\gamma\cos(\omega t)\mathbf{e}_x$, where E_γ and ω are the amplitude and frequency of the laser field. For notational convenience, both $\hbar$ and the speed of light c have been set to unity. We shall also neglect the effect of the magnetic-field component of the laser field, since it is smaller than the electrical-field component by a factor of c. Let us choose the vector potential for the laser field to be in the form

$$\mathbf{A}_\gamma=(E_\gamma/\omega)\sin(\omega t)\mathbf{e}_x. \tag{1}$$

The time-dependent Schrödinger equation for a single electron is given as

$$i\frac{\partial}{\partial t}\psi(\mathbf{r},t)=H\psi(\mathbf{r},t)=\frac{(\mathbf{p}-e\mathbf{A}_\gamma)^2}{2m^*}\psi(\mathbf{r},t). \tag{2}$$

It can be shown that Eq. (2) can be solved exactly. The time-dependent wave function can be written as

$$\psi_{\mathbf{k}}(\mathbf{r},t)=\exp(-i2\gamma_1\omega t)\exp\{i\gamma_0 k_x[1-\cos(\omega t)]\}\times\exp[i\gamma_1\sin(2\omega t)]\exp(-i\epsilon_k t)\exp(i\mathbf{k}\cdot\mathbf{r}), \tag{3}$$

where $\gamma_0=(eE_\gamma)/m^*\omega^2$ and $\gamma_1=(eE_\gamma)^2/(8m^*\omega^3)$. The effect of electrical field is included in this wave function exactly. We notice that these wave functions form an orthonormal set,

$$\langle\psi_{\mathbf{k}}(r,t)|\psi_{\mathbf{k}}(r,t)\rangle=\delta_{\mathbf{k},\mathbf{k}'}.$$

We now construct the quantum field operators

$$\Psi(\mathbf{r},t)=\sum_{\mathbf{k}} a_{\mathbf{k}}\psi_{\mathbf{k}}(\mathbf{r},t), \tag{4}$$

$$\Psi^\dagger(\mathbf{r},t)=\sum_{\mathbf{k}} a_{\mathbf{k}}^\dagger\psi_{\mathbf{k}}^*(\mathbf{r},t), \tag{5}$$

where $a_{\mathbf{k}}^\dagger$ ($a_{\mathbf{k}}$) is the creation (annihilation) operator for the electronic state with momentum $\mathbf{k}$. These field operators satisfy the equal-time commutation relation $\{\Psi^\dagger(\mathbf{r},t),\Psi(\mathbf{r}',t)\}=\delta(\mathbf{r}-\mathbf{r}')$. Alternatively, we can write these field operators in terms of eigenfunctions of a free electron,

$$\Psi(\mathbf{r},t)=\sum_{\mathbf{k}} b_{\mathbf{k}}(t)e^{i\mathbf{k}\cdot\mathbf{r}}, \tag{6}$$

$$\Psi^\dagger(\mathbf{r},t)=\sum_{\mathbf{k}} b_{\mathbf{k}}^\dagger(t)e^{-i\mathbf{k}\cdot\mathbf{r}}, \tag{7}$$

where

$$b_{\mathbf{k}}(t)=a_{\mathbf{k}}\exp(-i2\gamma_1\omega t)\exp(i\gamma_0 k_x[1-\cos(\omega t)])\times\exp[i\gamma_1\sin(2\omega t)].$$

We now calculate the current response of this electron-photon system to a weak electric field $\mathbf{E}(t)=\mathbf{E}e^{-i\Omega t}$, where Ω is the frequency of the probing field. The Hamiltonian of the system, in the second-quantized notation, can be written as

$$H=H_0+H_{ee}+H_{eI}, \tag{8}$$

where H_0 is the Hamiltonian of a noninteracting many-electron system,

$$H_0=\frac{1}{2m^*}\sum_{\mathbf{p}}[\mathbf{p}+e\mathbf{A}_\gamma+e\mathbf{A}]^2 b_{\mathbf{p}}^\dagger(t)b_{\mathbf{p}}(t), \tag{9}$$

where $\mathbf{A}=\mathbf{E}e^{-i\Omega t}/(i\Omega)$ is the vector potential of the probing field. H_{ee} is the electron-electron interaction

$$H_{ee}=\frac{1}{2}\sum_{\mathbf{p},\mathbf{p}',\mathbf{q}} V_q b_{\mathbf{p}+\mathbf{q}}^\dagger(t)b_{\mathbf{p}'-\mathbf{q}}^\dagger(t)b_{\mathbf{p}'}(t)b_{\mathbf{p}}(t)$$
$$=\frac{1}{2}\sum_{\mathbf{p},\mathbf{p}',\mathbf{q}} V_q a_{\mathbf{p}+\mathbf{q}}^\dagger(t)a_{\mathbf{p}'-\mathbf{q}}^\dagger(t)a_{\mathbf{p}'}(t)a_{\mathbf{p}}(t), \tag{10}$$

where $V_q=2\pi e^2/q$ is the Fourier transform of electron-electron interaction in two dimensions. H_{eI} is the interaction between the electrons and random impurities,

$$H_{eI}=-\sum_{\mathbf{p},\mathbf{q}} ZV_q b_{\mathbf{p}+\mathbf{q}}^\dagger(t)b_{\mathbf{p}}(t)\sum_i e^{i\mathbf{q}\cdot\mathbf{R}_i}, \tag{11}$$

where Z is the charge number of the impurity and $\mathbf{R}_i$ is the position of the ith impurity.

Now the total average two-dimensional (2D) current density of the system is defined as

$$\mathbf{j}=\left\langle\frac{\delta H}{\delta\mathbf{A}}\right\rangle=\frac{e}{m^*}\sum_{\mathbf{p}}\langle[\mathbf{p}+e\mathbf{A}_\gamma+e\mathbf{A}]b_{\mathbf{p}}^\dagger(t)b_{\mathbf{p}}(t)\rangle$$
$$=\mathbf{j}_1+\mathbf{j}_\gamma+\mathbf{j}_0. \tag{12}$$

Since $\mathbf{A}_\gamma=(E/\omega)\sin(\omega t)\mathbf{e}_x$,

$$\mathbf{j}_\gamma=\frac{e^2}{m^*}e\mathbf{A}_\gamma\sum_{\mathbf{p}}\langle b_{\mathbf{p}}^\dagger(t)b_{\mathbf{p}}(t)\rangle=\frac{e^2 n}{m\omega}\sin(\omega t)E_\gamma\mathbf{e}_x, \tag{13}$$

where n is the 2D electron concentration. The terahertz laser field is oscillating as $\cos(\omega t)$. Therefore $\mathbf{j}_\gamma$ is in the direction of the laser field but its phase is behind the phase of the laser field by π. Thus $\mathbf{j}_\gamma$ is purely reactive,

$$\mathbf{j}_0=\frac{e^2}{m^*}e\mathbf{A}\sum_{\mathbf{p}}\langle b_{\mathbf{p}}^\dagger(t)b_{\mathbf{p}}(t)\rangle=\frac{ine^2}{m\Omega}\mathbf{E}(t)=\sigma_0\mathbf{E}(t). \tag{14}$$

This is the well-known Drude conductivity in the absence of electron-impurity scattering. All effects on the electrical current due to the electron-electron interaction, electron-impurity interaction, and sideband transitions are contained in the current density $\mathbf{j}_1$,

$$\mathbf{j}_1=\frac{e}{m^*}\sum_{\mathbf{p}}\langle\mathbf{p}b_{\mathbf{p}}^\dagger(t)b_{\mathbf{p}}(t)\rangle. \tag{15}$$

The purpose of the current work is to develop a method to evaluate $\mathbf{j}_1$. Let us construct the single-electron density matrix,

$$F(\mathbf{p},\mathbf{p}+\mathbf{k})=\langle b_{\mathbf{p}}^\dagger(t)b_{\mathbf{p}+\mathbf{k}}(t)\rangle$$
$$=F_0(\mathbf{p},\mathbf{p}+\mathbf{k})\exp\{i\gamma_0 k_x[1-\cos(\omega t)]\}. \tag{16}$$

where $F_0(\mathbf{p},\mathbf{p}+\mathbf{k})=\langle a_{\mathbf{p}}^\dagger(t)a_{\mathbf{p}+\mathbf{k}}(t)\rangle$. The equation of motion for the single-electron density matrix is given as

$$i\frac{\partial}{\partial t}F(\mathbf{p},\mathbf{p}+\mathbf{k})=[\epsilon_{\mathbf{p}+\mathbf{k}}-\epsilon_{\mathbf{p}}+k_x\gamma_0\omega\sin(\omega t)]F(\mathbf{p},\mathbf{p}+\mathbf{k})$$
$$+\sum_{\mathbf{q}}V_q\left[n(\mathbf{q},t)-Z\sum_i e^{i\mathbf{q}\cdot\mathbf{R}_i}\right]$$
$$\times[F(\mathbf{p},\mathbf{p}+\mathbf{k}-\mathbf{q})-F(\mathbf{p}+\mathbf{q},\mathbf{p}+\mathbf{k})]. \tag{17}$$

Here $\epsilon_{\mathbf{p}}=p^2/2m$ is the kinetic energy of an electron having momentum $\mathbf{p}$, and

$$n(\mathbf{q},t)=\sum_{\mathbf{p}}F(\mathbf{p},\mathbf{p}+\mathbf{k}) \tag{18}$$

is the induced charge density. Making use of Eqs. (15) and (17), we obtain the following time-dependent current density $\mathbf{j}_1(t)$:

$$\frac{d\mathbf{j}_1(t)}{dt}=\frac{-Ze}{m^*}\sum_{\mathbf{q}}V_q\mathbf{q}n(-\mathbf{q},t)\sum_i e^{i\mathbf{q}\cdot\mathbf{R}_i}. \tag{19}$$

In deriving the above equation, we have used the identity $n^*(\mathbf{q},t)=n(-\mathbf{q},t)$ and the fact that

$$\sum_q\frac{\mathbf{q}}{q}n(\mathbf{q},t)n(-\mathbf{q},t)=0. \tag{20}$$

The result for the current density given in Eq. (19) is exact. $n(\mathbf{q},t)$ contains the information of electron-impurity interaction and sideband transition. This quantity can be written as a product of a field-dependent term and a field-independent term,

$$n(-\mathbf{q},t)=n_0(-\mathbf{q},t)\exp(i\gamma_0 q_x[1-\cos(\omega t)]), \tag{21}$$

where

$$n_0(-\mathbf{q},t)=\sum_{\mathbf{p}}F_0(\mathbf{p},\mathbf{p}+\mathbf{q}) \tag{22}$$

is the density fluctuation in the presence of impurities and the probing field E, but in the absence of laser field E_γ. Making use of the generating function for the Bessel function, $\exp[i\alpha\cos(x)]=\Sigma_m i^m J_m(\alpha)\exp(imx)$ (where J_m is the Bessel function of the first kind), we can decompose the density matrix, the induced density, and the total current density into a sum of successive harmonics,

$$F(\mathbf{p},\mathbf{p}+\mathbf{k})=\sum_m F^{(m)}(\mathbf{p},\mathbf{p}+\mathbf{k}), \tag{23}$$

$$F^{(m)}(\mathbf{p},\mathbf{p}+\mathbf{k})=(-i)^m J_m(k_x\gamma_1)e^{ik_x\gamma_0}F_0(\mathbf{p},\mathbf{p}+\mathbf{k})e^{-im\omega t}, \tag{24}$$

$$n^{(m)}(\mathbf{q},t)=\sum_{\mathbf{p}}F^{(m)}(\mathbf{p},\mathbf{p}+\mathbf{k}), \tag{25}$$

and

$$\mathbf{j}_1(t)=\sum_m\mathbf{j}_1^{(m)}(t) \tag{26}$$

and for the mth-order term, we have

$$\frac{d\mathbf{j}_1^{(m)}(t)}{dt}=\frac{-Ze}{m^*}\sum_{\mathbf{q}}V_q\mathbf{q}n^{(m)}(-\mathbf{q},t)\sum_i e^{i\mathbf{q}\cdot\mathbf{R}_i}. \tag{27}$$

Since $n_0(-\mathbf{q},t)$ only responds to E, we assume

$$n_0(-\mathbf{q},t)=n_0(-\mathbf{q},\Omega)e^{-i\Omega t},$$

$\mathbf{j}_1^{(m)}(t)$ must take the form $\mathbf{j}_1^{(m)}(t)=\mathbf{j}_1^{(m)}(\Omega,\omega)\exp[-i(\Omega+m\omega)t]$ $(\Omega=\Omega+i0)$. We now obtain

$$\mathbf{j}_1^{(m)}(\Omega,\omega)=\frac{-Ze}{m^*(\Omega+m\omega)}\sum_{\mathbf{q}}V_q\mathbf{q}n^{(m)}(-\mathbf{q},\Omega,m\omega)$$
$$\times\sum_i e^{i\mathbf{q}\cdot\mathbf{R}_i}. \tag{28}$$

The method to calculate the quantity $n^{(m)}(\mathbf{q},t)$ was developed by the author and his co-worker.[15] A detailed derivation can be found in the Ref. 15. Here we only briefly summarize

the steps involved in this method. We first assume $\mathbf{E}=0$ and calculate the change of the electron density matrix up to the second order in electron-impurity interaction, F_1. The zeroth-order density matrix is assumed to be of the form of the Fermi-Dirac distribution $f_{\mathbf{p}}$,

$$F_0(\mathbf{p},\mathbf{p}+\mathbf{k})=f_{\mathbf{p}}\delta_{\mathbf{k},0}. \tag{29}$$

The equation of motion for $F_1^{(m)}$ is given as

$$-[\epsilon_{\mathbf{p}+\mathbf{k}}-\epsilon_{\mathbf{p}}-m\omega]F_1^{(m)}(\mathbf{p},\mathbf{p}+\mathbf{k}) = V_k\left[n_1^{(m)}(\mathbf{k})-Z\sum_i e^{i\mathbf{k}\cdot\mathbf{R}_i}\right][f_{\mathbf{p}+\mathbf{k}}-f_{\mathbf{p}}]. \tag{30}$$

In the above equation, each term has a common factor $J_m(k_x\gamma_0)\exp(ik_x\gamma_0)$, which is not explicitly included. The solution of $F_1^{(m)}$ is given as

$$F_1^{(m)}=-\frac{ZV_k}{2\pi D^{(m)}(k,\omega)}\frac{f_{\mathbf{p}+\mathbf{k}}-f_{\mathbf{p}}}{\epsilon_{\mathbf{p}+\mathbf{k}}-\epsilon_{\mathbf{p}}-m\omega}\sum_i e^{i\mathbf{k}\cdot\mathbf{R}_i}, \tag{31}$$

where

$$D^{(m)}(q,\omega)=1-V_qQ(q,m\omega) \tag{32}$$

is the dielectric function of two-dimensional electrons in the random-phase approximation and

$$Q(q,\omega)=\int\frac{d\mathbf{p}}{(2\pi)^2}\frac{f_{\mathbf{p}+\mathbf{q}}-f_{\mathbf{p}}}{\epsilon_{\mathbf{p}+\mathbf{q}}-\epsilon_{\mathbf{p}}-\omega} \tag{33}$$

is the polarizability of free electrons.

Next we calculate the time-dependent part of the density matrix under the influence of the probing field and the impurities $F_2^{(m)}$. Only terms linear in the probing field were retained. The equation of motion for $F_2^{(m)}$ is given as

$$-[\epsilon_{\mathbf{p}+\mathbf{q}}-\epsilon_{\mathbf{p}}-\Omega-m\omega]F_2^{(m)}(\mathbf{p},\mathbf{p}+\mathbf{q}) = V_kn_2^{(m)}(\mathbf{q})[f_{\mathbf{p}+\mathbf{q}}-f_{\mathbf{p}}]+\frac{ie\mathbf{q}\cdot\mathbf{E}}{m^*\Omega}F_1^{(m)}(\mathbf{p},\mathbf{p}+\mathbf{q}). \tag{34}$$

Solving Eq. (34) we obtain the following result for $n^{(m)}(\mathbf{q},\Omega)$:

$$n^{(m)}(\mathbf{q},\Omega)=-\frac{iZe\mathbf{q}\cdot\mathbf{E}V_qJ_m(q_x\gamma_0)}{m^*\Omega(\Omega+m\omega)}\times\left[\frac{1}{D(q,m\omega)}-\frac{1}{D(q,\Omega+m\omega)}\right]\times\exp(i\gamma_0q_x)\sum_i e^{i\mathbf{q}\cdot\mathbf{R}_i}. \tag{35}$$

Substituting Eq. (35) into Eq. (28), making use of the normalization condition, $\Sigma_{i,j}e^{i\mathbf{q}\cdot(\mathbf{R}_i-\mathbf{R}_j)}=n/Z$, and then summing over m, we obtain

$$\mathbf{j}_1(\Omega,\omega)=\frac{-iZne^2}{m^*\Omega}\sum_{\mathbf{q}}\sum_m\frac{(-i)^mV_qJ_m(q_x\gamma_0)}{(\Omega+m\omega)^2}\times\mathbf{q}\mathbf{q}\cdot\mathbf{E}\left[\frac{1}{D(q,m\omega)}-\frac{1}{D(q,\Omega+m\omega)}\right]\exp(i\gamma_0q_x). \tag{36}$$

If the terahertz laser field is zero, $\gamma_0=0$, only the $m=0$ term contributes; in this case we recover the result of current response of an electronic system.[15,16]

The total current density due to the probing field can now be written as

$$\mathbf{j}_T=\mathbf{j}_0+\mathbf{j}_1=\sigma\mathbf{E}. \tag{37}$$

By comparing σ with the Drude conductivity $ine^2/(m^*\Omega)(I+i\tau^{-1}/\Omega)$, we can identify the transport scattering time,

$$\left[\frac{1}{\Omega\tau}\right]_{ij}=\frac{Z}{m^*}\sum_{\mathbf{q}}\sum_m\frac{V_qJ_m(q_x\gamma_0)}{(\Omega+m\omega)^2}q_iq_j\times\mathrm{Im}\left\{(-i)^m\left[\frac{1}{D(q,m\omega)}-\frac{1}{D(q,\Omega+m\omega)}\right]\times\exp(i\gamma_0q_x)\right\}. \tag{38}$$

For isotropic systems, the off-diagonal part of the inverse scattering time vanishes. Equation (38) can be simplified by making use of the following properties.

(a) The dielectric function is only dependent on the magnitude of q.

(b) The Bessel functions are symmetric for even m and antisymmetric for odd m with respect to the argument q_x. Therefore the integration over the direction of q will be nonzero for the products $\cos(q_x\gamma_0)J_{2m}(q_x\gamma_0)$ and $\sin(q_x\gamma_0)J_{2m+1}(q_x\gamma_0)$.

(d) ω now represents gap of electron-photon sidebands which must be real and thus $D(q,m\omega)$ is real. Only $D(q,\Omega+m\omega)$ has an imaginary part. We obtain

$$\left[\frac{1}{\Omega\tau}\right]_{ij}=\frac{Z}{m^*}\sum_{\mathbf{q}}V_qq_iq_j\left\{\mathrm{Im}\left[-\frac{1}{D(q,\Omega)}\right]\frac{J_0(q_x\gamma_0)}{\Omega^2}\times\cos(\gamma_0q_x)+\mathrm{Im}\left[-\frac{1}{D(q,\Omega+\omega)}\frac{1}{(\Omega+\omega)^2}-\frac{1}{D(q,\Omega-\omega)}\frac{1}{(\Omega-\omega)^2}\right]J_1(q_x\gamma_0)\sin(\gamma_0q_x)\cdots\right\}. \tag{39}$$

We have numerically calculated scattering time. The parameters used in our calculation are those of GaAs, $m^*=0.067m_0$, $r_s=m^*e^2/(\hbar^2k_F)=1.0$. We assume that impurities are singly ionized, $Z=1$. All energies and frequencies are in units of E_F and all wave numbers are in units of k_F. $\alpha=k_FeE/(m^*\omega^2)$ is the dimensionless electron-photon coupling parameter.

In Fig. 1 we plot the scattering time as a function of Ω at fixed field strength. Several interesting features can be observed. In the absence of the intense terahertz laser,[15] the scattering time initially decreases as Ω increases. It reaches a shallow minimum at $\Omega\approx E_F$. At large Ω, it increases with Ω. Therefore the high frequency approximation $\Omega\tau\gg1$ is usually valid for $\Omega>E_F$ is an electron gas. In the present case, this behavior is qualitatively altered due the photon sidebands. At low Ω, τ decreases initially with Ω. It has a series of minima at $\Omega=m\omega$. This is due to the simultaneous absorption of one photon with energy $\hbar\Omega$ and making a

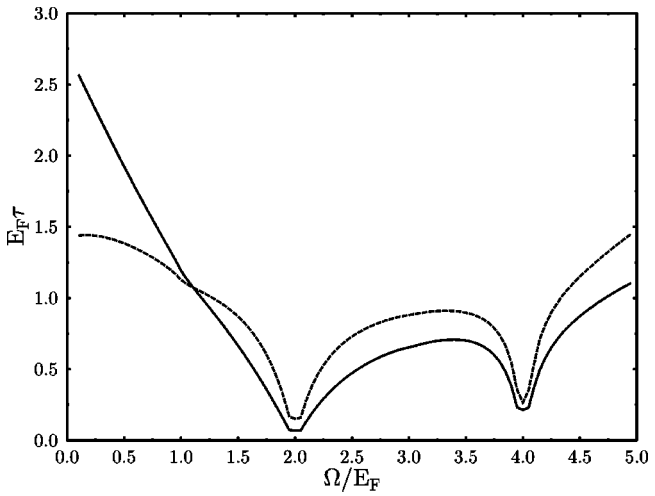

FIG. 1. Transport scattering time versus the frequency of the probing field. The electron-photon coupling $\alpha=1.5$, the terahertz laser frequency is $2E_F$. Parameters of GaAs semiconductor are used as specified in the text. The solid line is for τ_{xx} and the broken line is for τ_{yy}.

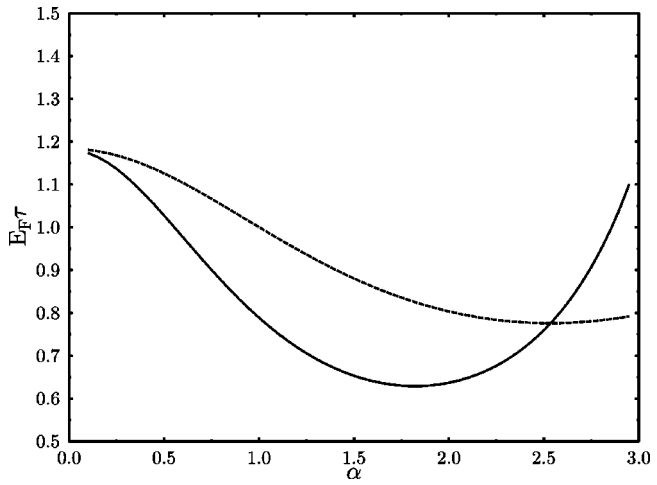

FIG. 2. Transport scattering time versus the electron-photon coupling α. The frequency of the probing field is $3E_F$, the terahertz laser frequency is $2E_F$. The solid line is for τ_{xx} and the broken line is for τ_{yy}.

transition between two electron-photon sidebands separated by $m\hbar\omega$. Such processes can be called resonant scattering. The scattering time for resonant scattering is extremely short as shown in Fig. 1. Between resonant frequencies, the scattering time first increases and then decreases. Since there are infinite number of electron-photon sidebands, the scattering time cannot reach a sufficiently large value. The small difference between τ_{xx} and τ_{yy} is due to our choice that the laser field is polarized along the x direction. Under this geometry, the scattering in the x direction is more frequent except at very low frequencies.

In Fig. 2, we plot the scattering time as a function of the dimensionless coupling parameter α. At fixed laser frequency ω, α can be viewed as the dimensionless electric field of the laser. In the region where the laser field is weak, electron energy increases with increasing laser field. This leads to an increased scattering or a short scattering time. At high laser field, the effect of strong coupling of electronic states with the photon field will be dominant. This strong coupling will make the electrons less mobile. As a result, the scattering time increases. This effect has a similar origin as that of the suppression of dc conductivity[6] and the reduction of collective excitation energies.[17]

In conclusion, we have studied theoretically and numerically the electron transport under an intense terahertz radiation. Resonant scattering due to transitions between various photon sidebands can occur in strongly coupled electron-photon systems.

This work was initiated when the author was a visiting professor at the National Center for Theoretical Sciences in Hsinchu, Taiwan. It is a pleasure to acknowledge financial support from the Center. Part of the work was carried out during a scientific exchange visit to Tokai University. The visit was financially supported by the Australian Academy of Sciences and the Japan Society for Promotion of Science. We thank T. Toyoda and W. Xu for interesting discussions.

[1] M.A. Odnoblyudov, I.N. Yassievich, M.S. Kogan, Yu.M. Gelperin, and K.A. Chao, Phys. Rev. Lett. **83**, 644 (1999).
[2] Yu.P. Gousev, I.V. Altukhov, K.A. Korolev, V.P. Sinis, M.S. Kogan, E.E. Haller, M.A. Odnoblyudov, I.N. Yassievich, and K.A. Chao, Appl. Phys. Lett. **75**, 757 (1999).
[3] N.G. Asmar, A.G. Markelz, E.G. Gwinn, J. Cerne, M.S. Sherwin, K.L. Campman, P.E. Hopkins, and A.C. Gossard, Phys. Rev. B **51**, 18 041 (1995).
[4] W. Xu and C. Zhang, Appl. Phys. Lett. **68**, 3305 (1996).
[5] W. Xu and C. Zhang, Phys. Rev. B **55**, 5259 (1997).
[6] N.G. Asmar, J. Cerne, A.G. Markelz, E.G. Gwinn, M.S. Sherwin, K.L. Campman, and A.C. Gossard, Appl. Phys. Lett. **68**, 829 (1996).
[7] A.G. Markelz, N.G. Asmar, B. Brar, and E.G. Gwin, Appl. Phys. Lett. **69**, 3975 (1996).
[8] B.N. Murdin, W. Heiss, C.J.G.M. Langerak, S.-C. Lee, I. Galbraith, G. Strasser, E. Gornik, M. Helm, and C.R. Pidgeon, Phys. Rev. B **55**, 5171 (1997).
[9] C.J.G.M. Langerak, B.N. Murdin, B.E. Cole, J.M. Chamberlain, M. Henini, M. Pate, and G. Hill, Appl. Phys. Lett. **67**, 3453 (1995).
[10] T.A. Vaughan, R.J. Nicholas, C.J.G.M. Langerak, B.N. Murdin, C.R. Pidgeon, N.J. Mason, and P.J. Walker, Phys. Rev. B **53**, 16 481 (1996).
[11] C. Zhang, Appl. Phys. Lett. **78**, 4187 (2001).
[12] P. Orellana, F. Claro, and E. Anda, Phys. Rev. B **62**, 9959 (2000).
[13] P. Orellana and F. Claro, Appl. Phys. Lett. **75**, 1643 (1999).
[14] T. Brandes, Europhys. Lett. **33**, 629 (1996).
[15] N. Tzoar and C. Zhang, Phys. Rev. B **32**, 1146 (1985).
[16] N. Tzoar, P.M. Platzman, and A. Simon, Phys. Rev. Lett. **36**, 1200 (1976).
[17] C. Zhang, Phys. Rev. B **65**, 153107 (2002).