# Plasmon induced terahertz absorption and photoconductivity in a grid-gated double-quantum-well structure

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The terahertz (THz) absorption spectra of plasmon modes in a grid-gated double-quantum-well field-effect transistor structure is analyzed theoretically and numerically using the scattering matrix approach and is shown to faithfully reproduce strong resonant features of recent experimental observations of THz photoconductivity in such a structure.

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Recently, the terahertz (THz) photoconductivity of a double-quantum-well (DQW) field-effect transistor (FET) with a periodic metal grid gate was observed [1]. It appears that presence of a DQW is important to produce a strong photoresponse. The position and the strength of the peaks in the photoresponse, photoconductance (the change of channel conductance), are controlled by both the voltage $V_g$ applied to the gate and the period of the grating gate, which strongly suggests that the sharp resonant features are related to plasma oscillations in this composite device.The strongest THz photoresponse occurs when the upper QW is fully depleted under metallic portions of the grating gate, while the lower one remains connected but with laterally modulated electron density. The response is fast (less than $1\,\mu$s) and depends on the presence of the DQW channel. Remarkably, the sharpest and strongest resonance response does not occur at the lowest temperatures but rather between 25 and 40 K. The sign of the photoconductance is not unique and depends on the grating period. The physical mechanism responsible for the THz photoconductivity is not evident at this time. Particularly, the temperature dependence, as well as the reversal of sign of the photoconductance, are not understood.

A simple „transmission“ model discussed in [1] describing plasma oscillations in the grid-gated DQW FET provided a rough qualitative guide suggesting that the series of frequency dependent resonances observed in the THz photoconductance corresponds to the excitation of plasma waves the in structure. However, this simple model fails to explain the measured separation in $V_g$ between resonances at different frequencies and the variation of the lineshape with frequency.

Far infrared absorption in density modulated two-dimensional electron systems has been studied theoretically in a series of papers [2–9], including studies of density modulated bilayers [2,4,6] and structures with metal gratings [3,5].

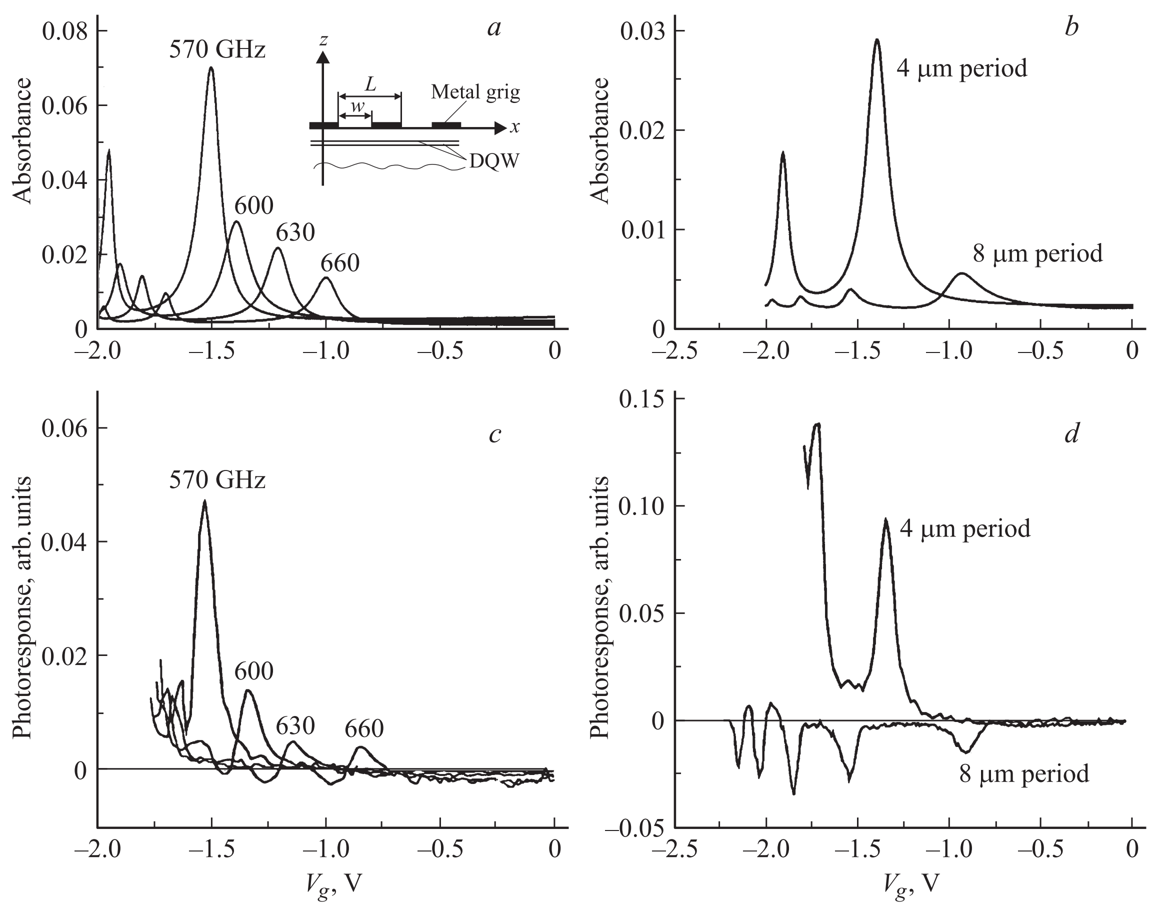
In this paper, we develop a detailed, realistic model of plasmon electrodynamics in the grid-gated DQW FET structure studied experimentally in [1]. We present the results of our electrodynamic simulation of the THz absorption spectrum taking acoount of all essential features of the actual composite structure studied in [1], which faithfully reproduce strong resonant features of experimental observations of THz photoconductivity in such a structure.

## 1. Theoretical model

We analyze the THz absorption spectrum in the DQW multilayered laterally periodic structure of [1] (see the inset in Figure, $a$) using a scattring matrix approach [10–12]. The scattering matrix approach employed involves the following steps. Within each layer of the structure, the Maxwell equations are rewritten in an in-plane momentum matrix representation starting with decomposition of the magnetic field into a sum of plane waves

$$\mathbf{H}(x,z,t)=\sum_m \mathbf{H}_m \exp(-i\beta_m x)\exp(ik_z z)\exp(-i\omega t), \tag{1}$$

where $m=0,\pm1,\pm2,\ldots,\pm M$; $\beta_m=k_x+2\pi m/L$; $L$ is the period of the structure; $k_x$ is the reduced in-plain wavevector. In the case under consideration, when THz




Calculated absorption spectra: (*a*) at several different THz frequencies for $L = 4\,\mu$m, $w = L/2$ and (*b*) at frequency 600 GHz for two different periods of the structure with $w = L/2$, $V_{\rm th} = -2.6$ V is a fitted value. Other parameters of the structure are given in the text. Terahertz photoresponse measured as a function of gate voltage: (*c*) for four different frequencies for $L = 4\,\mu$m, $w = L/2$ and (*d*) at frequency 600 GHz for the different periods of the metal gate with $w = L/2$ at $T = 25$ K, as reported in [1].

radiation is incident normally onto the multilayered structure, $k_x = 0$ and there are only $E_{x0}$ and $H_{y0}$ non-zero components of the electric and magnetic field in the incident wave. Assuming the square profile of the lateral periodicity of the structure, the dielectric function $\varepsilon(x)$ in each layer transforms into matrix $\hat{\varepsilon}$ with elements

$$\varepsilon_{mm'} = \int_{-L/2}^{L/2} \varepsilon(x) \exp\left[-\frac{2\pi(m-m')}{L}\right] dx. \qquad (2)$$

Obviously, for homogeneous layers we have $\hat{\varepsilon} = \varepsilon\hat{I}$, where $\varepsilon$ is constant and $\hat{I}$ is the unit matrix.

As a result, in each layer, we arrive at the matrix eigenproblem

$$\left[\hat{\varepsilon}\left(\frac{\omega^2}{c^2} - \hat{\beta}\hat{\varepsilon}^{-1}\hat{\beta}\right)\right]\mathbf{A} = k_z^2\mathbf{A}, \qquad (3)$$

where $\hat{\beta}$ is the diagonal matrix with elements $\beta_{mm'} = \delta_{mm'}\beta_m$. The $(2M+1)$-dimensional eigen-vectors $\mathbf{A}^{(n)}$, where $n = 1, 2, \ldots, 2M+1$, have normalized amplitudes of the magnetic field $H_{ym}$ as their components. Eigen-vectors $\mathbf{A}^{(n)}$ and $\mathbf{A}^{(n')}$ corresponding to eigen-values $\left(k_z^{(n)}\right)^2$ and $\left(k_z^{(n')}\right)^2$ satisfy the orthogonality relationship $\mathbf{A}^{(n)}\mathbf{A}^{(n')} = \delta_{nn'}$. In particular, for homogeneous layers, the eigen-value problem (3) yields the eigen-vectors $\mathbf{A}^{(n)}$ with components $\delta_{nm}$ corresponding to eigen-values

$$\left(k_z^{(n)}\right)^2 = \varepsilon\,\frac{\omega^2}{c^2} - \delta_{nm}\beta_m.$$

Eigen-problem (3) is degenerate in respect to the reversal of sign of the transverse wavevector $k_z$. Thus, the total field in each layer is a superposition of eigen-modes of the form (1) with $k_z = \pm\left|k_z^{(n)}\right|$, where $n = 1, 2, \ldots, 2M+1$, propagating along and counter to $z$-axis. Applying boundary conditions of continuity of $E_x$ and $H_y$ components of the total field at interfaces of the layers, we can relate the amplitudes of waves outgoing from the whole structure to those of incoming waves in the matrix form

$$\begin{pmatrix} \mathbf{A}^+_{\rm out} \\ \mathbf{A}^-_{\rm out} \end{pmatrix} = \mathbb{S}\begin{pmatrix} \mathbf{A}^+_{\rm in} \\ \mathbf{A}^-_{\rm in} \end{pmatrix},$$

where superscripts $+$ and $-$ refer to the waves propagating along and counter to $z$-axis, respectively, and $\mathbb{S}$ is the scattering matrix of the whole structure, which can be written in the $2 \times 2$ block matrix form

$$\mathbb{S} = \begin{pmatrix} \hat{S}_{11} & \hat{S}_{12} \\ \hat{S}_{21} & \hat{S}_{22} \end{pmatrix}.$$

In our particular problem under consideration we have only one incoming plane wave incident onto the metal-grid gate (i.e. $\mathbf{A}^{+}_{\text{in}} = 0$) and there is only the $m = 0$ non-zero component in the vector $\mathbf{A}^{-}_{\text{in}}$. Thus, we can calculate the amplitudes of reflected and transmitted waves as

$$\mathbf{A}^{+}_{\text{out}} = \hat{S}_{12}\mathbf{A}^{-}_{\text{in}}, \qquad \mathbf{A}^{-}_{\text{out}} = \hat{S}_{22}\mathbf{A}^{-}_{\text{in}}.$$

Since the period of the structure under consideration ($L \simeq 4\,\mu$m) is roughly two orders of magnitude shorter than the THz radiation wavelength, only $m = 0$ plane-wave components in the reflected and transmitted waves survive at distances much longer than $L$ away from the structure, while all $m \neq 0$ plane-wave components are evanescent. Ultimately, we can readily calculate reflectivity $R$, tramsmittivity $T$, and absorbance $A = 1 - R - T$. For more detailed description of the scattering matrix formalism for multilayered laterally periodic structures, see, for example, [12].

In carrying out our calculations, we assume that the effective thickness $d$ of the upper and lower QWs in the DQW structure is 10 nm each, the separation between the center planes of the QWs $\Delta = 27$ nm, and the distance between the upper (near to the metal-grid gate) QW and grating gate $h = 385$ nm. We assume that the upper QW is fully depleted under the metallic portions of the grating gate of period $L$, forming an array of disconnected electron strips of width $w$ with areal electron density $N_U = 1.7 \cdot 10^{11}\,\text{cm}^{-2}$. The lower QW is density modulated with $N_L = 2.57 \cdot 10^{11}\,\text{cm}^{-2}$, essentially unchanged under the electron strips of the upper QW (i.e. under the open parts of the grating gate). The response of electron gas in QWs to the terahertz electric field $E\exp(-i\omega t)$ is described by a local bulk conductivity in the Drude form

$$\sigma_{U(L)}(x) = \frac{e^2 N_{U(L)}(x)}{m^* d} \frac{\tau}{1 - i\omega\tau},$$

where $e$ and $m^*$ are the electron change and effective mass, $\tau$ is the phenomenological electron relaxation time.

The electron density in the lower QW under metallic portions of the grating gate as a function of the gate voltage $N_L(V_g)$ was obtained from the parallel plate capacitor model with the gate voltage depletion threshold value $V_{\text{th}}$ and zero-gate-voltage electron density in both QWs combined chosen as fitting parameters to the experimental data of [1]. The parallel plate capacitor model assumed in this paper is fairly well justified for characteristic parameters of the structure studied in [1]: $L \gg h > d, \Delta$. This electrostatics problem is solved separately and the obtained value $N_L(V_g)$ is used as an input parameter in electrodynamic modeling. The metal grating strips are composed of a 20 nm Ti-layer covered by a 50 nm Au-layer. We took the bulk conductivities as $\sigma_0 = 2.3 \cdot 10^4$ S/cm for Ti and $\sigma_0 = 4.85 \cdot 10^5$ S/cm for Au.

Dielectric function $\varepsilon(x)$ entering expression (2) in QW layers is taken as

$$\varepsilon_{U(L)}(x) = \varepsilon_b + i\,\frac{4\pi\sigma_{U(L)}(x)}{\omega},$$

where $\varepsilon_b = 12.24$ in the background dielectric constant. Inside the metal-grid layer, we assume $\varepsilon(x) = i4\pi\sigma_0/\omega$ at the metal strips and $\varepsilon(x) \equiv 0$ at the grid openings.

## 2. Results and discussion

Figure, *a* shows the resonant absorption of incident THz electromagnetic radiation at several frequencies vs. gate voltage. It is seen in Figure, *a, c* that the variation of lineshape of plasma resonances with frequency is very close to that observed in photoconductivity resonances. In our calculations we used a fitting parameter $\tau = 10^{-11}$ s. This is roughly an order of magnitude less than the value extracted from the DC mobility [1]. Although the THz measurement temperatures are substantially above that used to measure the DC mobility, it is not likely that the DV mobility has dropped by an order of magnitude. On the other hand, one may consider an inhomogeneous broadening of the plasma resonance: the mean free time for a carrier to traverse a period of the modulation is $\sim L/v_F$, where $v_F$ is the Fermi velocity of electrons in the QWs, yielding $\sim 10^{-11}$ s for characteristic parameters of the structure under consideration, which is the same order of magnitude as the scattering used to fit the data.

The calculated THz absorption spectra and measured THz photoresponce for two different periods of the DQW FET structure are shown in Figure, *b, d*. It is evident, for both periods, that the positions of the resonances in photoresponce relate to those of the corresponding plasma resonances, although the signs of the photoresponse measurements are opposite for the different periods involved.

It can be seen in the Figure that the separation between the resonances of THz photoresponce exceeds the calculated sepatation by about 15%. This may be due to the fact that the parallel plate capacitor model we use to calculate the equilibrium electron density in the lower QW suffers from the neglect of fringing field effects.

While the resonant positions and strengths are well reproduced by our electrodynamic model, it has not dealt with the physical mechanism whereby excitation of the plasmon resonances induces changes in the DC transport. The sign change shown in Figure, *d*, underscores the fact that this mechanism is not at all understood.

More sophisticated modeling, which takes into account the interwell electron transfer in grid-gated DQW FET structures, is now in progress. We believe that this will provide an even better correspondence between theoretical and experimental results, as well as advance uor understanding of the physical mechanism that gives rise to the change in conductance at resonance.

## References


[1] X.G. Peralta, S.J. Allen, M.C. Wanke, N.E. Harff, J.A. Simmons, M.P. Lilly, J.L. Reno, P.J. Burke, J.P. Eisenstein. Appl. Phys. Lett. **81**, 1627 (2002).
[2] S. Katayama. J. Phys. Soc. Jpn. **60**, 1123 (1991).
[3] R.J. Wilkinson, C.D. Ager, T. Duffield, H.P. Hughes, D.G. Hasko, H. Ahmed, J.E.F. Frost, D.C. Peacock, D.A. Ritchie, G.A.G. Jones, C.R. Whitehouse, N. Absley. J. Appl. Phys. **71**, 6049 (1992).
[4] S. Katayama. Surf. Sci. **263**, 359 (1992).
[5] L. Schaich, P.W. Park, A.H. MacDonald. Phys. Rev. B **46**, 12 643 (192).
[6] C. Steinebach, D. Heitmann, V. Gudmundsson. Phys. Rev. B **56**, 6742 (1997).
[7] B.P. Van Zyl, E. Zaremba. Phys. Rev. B **59**, 2079 (1999).
[8] O.R. Matov, O.F. Meshkov, V.V. Popov. JETP **86**, 538 (1998).
[9] O.R. Matov, O.V. Polischuk, V.V. Popov. JETP **95**, 505 (2002).
[10] C.D. Ager, H.P. Hughes. Phys. Rev. B **44**, 13 452 (1991).
[11] D.M. Whittaker, I.S. Culshaw. Phys. Rev. B **60**, 2610 (1999).
[12] S.G. Tikhodeev, A.L. Yablonskii, E.A. Muljarov, N.A. Gippius, T. Ishihara. Phys. Rev. B **66**, 045 102 (2002).