Nonequilibrium Green’s Function Theory of Resonant Steady State Photoconduction in a Double Quantum Well FET subject to THz Radiation at Plasmon Frequency

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Abstract. Recent experimental observations by X.G. Peralta and S.J. Allen, et al. \[1\] of dc photoconductivity resonances in steady source-drain current subject to terahertz radiation in a grid-gated double-quantum well FET suggested an association with plasmon resonances. This association was definitively confirmed for some parameter ranges in our detailed electrodynamic absorbance calculations. \[2\],\[3\]. In this paper we propose that the reason that the dc photoconductance resonances match the plasmon resonances in semiconductors is based on a nonlinear dynamic screening mechanism. In this, we employ a shielded potential approximation that is nonlinear in the terahertz field to determine the nonequilibrium Green’s function and associated density perturbation that govern the nonequilibrium dielectric polarization of the medium. This “conditioning” of the system by the incident THz radiation results in resonant polarization response at the plasmon frequencies which, in turn, causes a sharp drop of the resistive shielded impurity scattering potentials and attendant increase of the dc source-drain current. This amounts to disabling the impurity scattering mechanism by plasmon resonant behavior in nonlinear screening.

1. Introduction

In the quest for a tunable terahertz (THz) detector, X.G. Peralta and S.J. Allen, et al. \[1\] performed experiments on a grid-gated double quantum well field effect transistor (FET) subjected to THz radiation. Figure 1 provides a schematic illustration of the experimental configuration. They measured linear source-to-drain dc photoconductance in the double quantum well channel, finding resonant behavior as a function of the gate voltage (which determines the equilibrium density distribution), suggesting that it might be associated with plasmon resonances. Some of their resonance measurements are shown in Figure 2. This resonant behavior was also found to depend on the strength of a magnetic field applied parallel to the quantum well layers. The resonances also exhibited unusual temperature dependence, shown in Figure 3, in that the peaks become narrower and higher as temperature increases up to 40°K, thereafter diminishing with further temperature increase. The double quantum well is nanometer scale, the fingers of the grating gate (and their separation) are micron scale. The gate voltage is sufficiently high that it perforates the mobile carrier charge distribution of the upper quantum well under the fingers of the gate, while the charge distribution of the lower
quantum well is mildly modulated. The incident THz radiation is in the range of 600 GHz with wavelength much greater than nanometer scale.

This paper presents a brief summary of our electrodynamic analysis of the grid-gated double quantum well FET involved, and a nonequilibrium Green’s function study of nonlinear screening that may account for the resonant photoconductance phenomenology by partial disabling of impurity scattering mechanisms.

2. Electrodynamic Analysis
To confirm the correspondence of the photoconductivity resonances with plasmon resonances we carried out an electrodynamic analysis of THz absorption. In a first stage [2], noting that the electromagnetic wavelength was much greater than the nanodimensions of the system, we assumed the behavior to be that of a single quantum well having a periodic equilibrium density distribution given by the sum of the densities of the lower and perforated upper quantum wells (Figure 4). In this, the "non-perforated" sum was taken as \( N_A = 4.27 \times 10^{11} \text{ cm}^{-2} \) and the "perforated" sum, \( N_B(V_g) \), varies with the gate voltage \( V_g \), decreasing as \(|V_g|\) increases. Further, we assumed a simple Drude surface conductivity. Needless to say, this treatment involving a single density modulated quantum well neglects the role of acoustic plasmons. However, the calculated absorption spectrum for this composite structure exhibits resonances (as a function of gate voltage) that are in good agreement with the Peralta-Allen observed resonances at temperature 25°K. The plasmon resonances determined in this electromagnetic absorption analysis [2] are shown side by side with the measured photoconductance resonances in Figure 5. Furthermore, we calculated the amplitudes of induced charge and current density oscillations in this single density modulated quantum well model assuming that "perforation" occurs over half the period: Examining the \( n = 2 \) plasmon at frequency 570 GHz, we found a clearly dipolar perturbed charge distribution in the "non-perforated" half-period, as shown in Figure 6.

To proceed with a more accurate representation of the experimental configuration, we also carried out an electrodynamic analysis of absorption by the grid gated double quantum well subject to THz radiation, shown in Figure 7. The equilibrium densities involved were determined by examining the 2D density expressions as functions of the local chemical potential,

\[
n_{1(2)} = D_2 \left[ E_F - E_{1(2)} + e\phi_{QW1(2)} \right]
\]

where \( D_2 = m^*/\pi \hbar^2 = 2D \) state density and \( E_{1(2)} \) is the energy of the bottom of the energy subband in the quantum well \( 1(2) \), \( E_F \) is the Fermi energy. These relations were solved jointly with the Poisson relations (Figure 8):

\[
\phi_{QW1} = V_G - \chi h (n_1 + n_2)
\]

\[
\phi_{QW2} = V_G - \chi h (n_1 + n_2) - \chi d n_2
\]

where \( V_G \) is the gate voltage, \( h \) is the separation of the upper quantum well from the gate, \( d \) is the separation of the two quantum wells and \( \chi = 4\pi e/\varepsilon_h \). (\( \varepsilon_h \) is the background dielectric constant.) The resulting "non-perforated" electrostatic equilibrium density solutions are given for\( V_G > V_{QW1}^{(th)} \) by

\[
n_1 = N_1 + \frac{D_2 e V_G (1 + D_2 \chi e d)}{1 + 2D_2 \chi e h + D_2 \chi e d (1 + D_2 \chi e h)}
\]

\[
n_2 = N_2 + \frac{D_2 e V_G}{1 + 2D_2 \chi e h + D_2 \chi e d (1 + D_2 \chi e h)}
\]

and for \( V_{QW2}^{(th)} < V_G < V_{QW1}^{(th)} \) by
\begin{align}
\n1 &= 0, \quad (4a) \\
2 &= N_2 + \frac{D_2 e (\chi h N_1 + V_g)}{1 + D_2 \chi e (d + h)}, \quad (4b)
\end{align}

where

\begin{equation}
V_{QW1}^{(th)} = -N_1 \frac{1 + 2D_2 \chi e + D_2 \chi e d (1 + D_2 \chi e)}{D_2 e}.
\end{equation}

and

\begin{equation}
V_{QW2}^{(th)} = -N_2 \frac{1 + 2D_2 \chi e (d + h)}{D_2 e} - \chi h N_1.
\end{equation}

\(N_1\) and \(N_2\) are defined in the absence of a gate voltage \((V_G = 0)\) as

\begin{align}
\n1 &= N_1 = 1.7 \times 10^{11} \text{ cm}^{-2}, \quad (6a) \\
2 &= N_2 = 2.57 \times 10^{11} \text{ cm}^{-2}. \quad (6b)
\end{align}

The electromagnetic field equations for the two quantum well system may be written in the form (the incident THz wave is \(E^{(0)} \exp(-ik_0 y - i\omega t)\), polarized along the \(x\)-axis, where \(k_0 = \omega \sqrt{\varepsilon_h / c}\) : 

\begin{equation}
E_x(x, y) = -\frac{2 \pi i}{\varepsilon_h \omega} \sum_{m = -\infty}^{\infty} \alpha_m \left[ J_m^{(L)} \exp(\alpha_m |y|) + J_m^{(U)} \exp(\alpha_m |y - d|) \right] \\
\times \exp(i \beta_m x) + E^{(0)} \exp(-ik_0 y),
\end{equation}

\begin{equation}
E_y(x, y) = \frac{2 \pi}{\varepsilon_h \omega} \sum_{m = -\infty}^{\infty} \beta_m \left[ J_m^{(L)} \text{sign}(y) \exp(\alpha_m |y|) + J_m^{(U)} \text{sign}(y - d) \exp(\alpha_m |y - d|) \right] \\
\times \exp(i \beta_m x),
\end{equation}

and

\begin{equation}
H_z(x, y) = -\frac{2 \pi}{c} \sum_{m = -\infty}^{\infty} \left[ J_m^{(L)} \text{sign}(y) \exp(\alpha_m |y|) + J_m^{(U)} \text{sign}(y - d) \exp(\alpha_m |y - d|) \right] \\
\times \exp(i \beta_m x) + E^{(0)} \sqrt{\varepsilon_h} \exp(-ik_0 y),
\end{equation}

where \(\alpha_m = \pm \sqrt{\beta_m^2 - k_0^2}\), \(\beta_m = 2 \pi m / L\) with \(m\) an integer, and \(J_m^{(L)}\) and \(J_m^{(U)}\) are the amplitudes of the Fourier components of the induced sheet current densities in the lower and upper 2D layer planes \(y = 0\) and \(y = d\), respectively, which are given by

\begin{equation}
J_m^{(L,U)} = \frac{1}{L} \int_0^L J_m^{(L,U)}(x) \exp(-i \beta_m x) \, dx.
\end{equation}

On the other hand, the induced currents are given by
\[ J^{(L)}(x) = \sigma L E_x(x,0) \text{ and } J^{(U)}(x) = \sigma U E_x(x,d), \]  
with the Drude conductivity as
\[ \sigma_{L,U} = \frac{e^2 N_{L,U}}{m^*} \frac{\tau}{1 - i\omega \tau}. \]

Substitution of \( E_x(x,y) \) given above yields integral equations for the currents as
\[ J^{(L)}(x) = \sigma L \int_0^L \left[ G_L(x,x') J^{(L)}(x') + G_U(x,x') J^{(U)}(x') \right] dx' + \sigma_L E^{(0)}, \]
and
\[ J^{(U)}(x) = \sigma L \int_0^L \left[ G_U(x,x') J^{(L)}(x') + G_L(x,x') J^{(U)}(x') \right] dx' + \sigma_U E^{(0)} \exp(-ik_0d), \]
with kernels given by
\[ G_L(x,x') = -\frac{2\pi i}{\varepsilon_k \omega L} \sum_{m=-\infty}^{\infty} \alpha_m \exp \left[ i \beta_m (x-x') \right], \]
\[ G_U(x,x') = -\frac{2\pi i}{\varepsilon_k \omega L} \sum_{m=-\infty}^{\infty} \alpha_m \exp \left[ \alpha_m d + i \beta_m (x-x') \right]. \]

These equations are solved numerically by the Galerkin procedure.

The calculated results for THz absorption by the double quantum well are in close agreement with those of the single quantum well shown in Figure 5. The correspondence between their resonances with the observed photoconductance resonances are exhibited in Figure 9, confirming the identity of the photoconductivity resonances as plasmon resonances. We also examined the in-plane current density oscillations in the upper and lower quantum wells for the second \( n = 2 \) plasmon mode. The results are shown in Figure 10. Due to perforation by the gate voltage, there is no current in the upper quantum well under the gate fingers. While the absorbance is relatively insensitive to the fact that the wells are separated, the finite separation is significant in giving rise to the acoustic and optical plasmons and is important in the photoconductivity. Because of the asymmetry of the density distributions in the two wells, the DQW admixes the optical and acoustic plasmons, so that acoustic oscillations appear on the principally optical-like plasmon background under the openings of the metal grid-gate. The optical plasmon component plays the major role, controlling the level of excitation of the standing plasma mode, but the acoustic plasmon admixture adds important structure in the spatial distribution of the current densities in the QWs.

The optical-plasmon-like resonances are characterized by powerful in-phase charge oscillations in the two different layers with a fairly weak admixture of acoustic (out-of-phase) oscillations. The optical-plasmon-like resonances are quite pronounced in the spectra because the optical-plasmon component has a strong net dipole moment ensuring effective excitation of the optical-plasmon-like mode by incoming THz radiation. The much weaker acoustic-plasmon-like resonances are characterized predominantly by out-of-phase charge oscillations in different layers with a small admixture of optical (in-phase) oscillations. Note that in a symmetric bilayer \( (N_L = N_U) \) optical and acoustic plasmons are uncoupled and the acoustic plasmons are optically inactive because of their vanishingly small net dipole moment in the plane of the bilayer, which is typical for acoustic plasmons in structures with \( d \ll w \).
Acoustic-plasmon-like resonances move toward lower frequencies when the separation between the 2D electron layers decreases. On the contrary, optical-plasmon-like resonances move toward higher frequencies when \( d \) decreases, with frequency approaching the frequency of the corresponding optical-plasmon resonance in a single 2D electron layer with a combined electron density \( N = N_L + N_U \). In the process of such opposing movements of the resonances the acoustic-plasmon-like resonances pass through the optical plasmon-like resonances, and in this case a clear anticrossing between these two types of plasma resonances is exhibited. Figure 11 shows the THz transmission spectra in the frequency range of the anticrossing between the fundamental optical-plasmon-like resonance and the 4th acoustic-plasmon-like resonance. In the anticrossing regime, these two resonances effectively exchange their oscillator strengths. As a result, the more powerful optical-plasmon-like resonance supplies the weaker acoustic-plasmon-like resonance with additional oscillator strength. Exactly at the anticrossing point (\( d = 19.8 \) nm) both resonances have equal intensity.

Far away from the optical-plasmon-like resonance (e.g. at \( d = 27 \) nm), the acoustic-plasmon-like resonance is virtually unnoticeable and it appears in the spectra only when it becomes stronger due to interaction with a powerful optical-plasmon-like resonance. On the other hand, the acoustic-plasmon-like resonance brings a greater admixture of the acoustic component into the optical-plasmon-like resonance in the anticrossing regime. As a result, the optical-plasmon-like resonance, still retaining large optical strength (i.e., net dipole moment), acquires a strong superimposed acoustic oscillatory component leading to huge enhancement of the THz interlayer electric field between the QWs (which may assist tunnelling between them).

Figures 12 and 13 show the distributions of amplitudes of the in-plane charge density oscillations in the individual 2D electron layers (Fig. 12) and that of the interlayer electric field (Fig. 13) at the fundamental optical-plasmon-like resonance far from (\( d = 27 \) nm) and within (\( d = 19.8 \) nm) the anticrossing regime. In the anticrossing regime, both the acoustic component and amplitude of the interlayer electric field dramatically increase; the latter becomes two and a half orders of magnitude greater than the amplitude of electric field of the incident THz wave.

### 3. Nonequilibrium Green’s Functions and Screening

While it is clear that the photoconductance resonances are associated with plasmon resonances, the mechanism by which the plasmons influence the Peralta-Allen dc transport experiments is not yet understood. The likelihood is that their linear source-drain conductance measurements are nonlinearly affected by the incident THz field. With this in mind, we have explored the role of the THz field in screening, starting from the recognition that the screened effective potential \( V(1) \) at a space-time point \( 1 = r_1,t_1 \) generated by an impressed potential \( U(2) \) at \( 2 = r_2,t_2 \) is given by

\[
V(1) = U(1) + \int d(4)3 \nu(1 - 3)\rho(3),
\]

where \( \nu(1 - 3) \) is the instantaneous Coulomb potential among the plasma electrons and \( \rho(3) \) is the perturbed density. A dynamic, nonlocal, spatially inhomogeneous and nonlinear variational derivative screening function, \( K_{v-d}(1,2) \), may be defined as \[4\]

\[
K_{v-d}(1,2) = \frac{\delta V(1)}{\delta U(2)} = \delta^{(4)}(1 - 2) + \int d(4)3 \nu(1 - 3)\frac{\delta \rho(3)}{\delta U(2)},
\]

alternatively,

\[
K_{v-d}(1,2) = \delta^{(4)}(1 - 2) - i \int d(4)3 \nu(1 - 3)D(3,2),
\]
where \( D(3, 2) = \left( \langle [\hat{\rho}(3) - \rho(3)][\hat{\rho}(2) - \rho(2)] \rangle \right) \) is the density-density correlation function (time-ordered) with \( \hat{\rho} \) as the density operator and \(< \ldots > \) denotes the nonequilibrium average. Using the chain rule of variational differentiation, we have [4]

\[
K_{v-d}(1, 2) = \delta^{(4)}(1 - 2) + \int d^{(4)}3 \int d^{(4)}4 \nu(1 - 3) \frac{\delta \rho(3)}{\delta V(2)} K_{v-d}(4, 2).
\]

The corresponding direct dielectric function, \( \varepsilon_{v-d}(1, 2) \), is defined as the space-time inverse of \( K_{v-d}(1, 2) \),

\[
\int d^{(4)}3 \varepsilon_{v-d}(1, 3) K_{v-d}(3, 2) = \delta^{(4)}(1, 2),
\]

whence

\[
\varepsilon_{v-d}(1, 2) = \frac{\delta U(1)}{\delta V(2)} = \delta^{(4)}(1 - 2) - \int d^{(4)}4 \nu(1 - 3) \frac{\delta \rho(3)}{\delta V(2)}.
\]

These simple remarks identify the associated dynamic, nonlocal, inhomogeneous and nonlinear polarizability \( \alpha(1, 2) \) in position-time representation as

\[
\alpha(1, 2) = -\int d^{(4)}3 \nu(1 - 3) R(3, 2),
\]

with the density perturbation response function \( R(3, 2) \) given by

\[
R(3, 2) = \frac{\delta \rho(3)}{\delta V(2)}.
\]

The nonlinear determination of \( R \) proceeds from the relation of \( \rho \) to the nonequilibrium one-electron Green’s function, \( G_1 \),

\[
\rho(3) = -iG_1(3, 3^+),
\]

jointly with the use of a time-dependent Hartree-approximation for the Schrödinger/Dyson \( G_1 \)-equation (space-time matrix notation)

\[
G_1 = G_0 + G_0 V G_1,
\]

where \( G_0 \) is the equilibrium thermodynamic Green’s function and we take \( V \) as the screened effective potential (Eq. 13) generated by polarization of the medium induced by the external impressed THz potential \( U \). Iterating this equation twice to obtain \( G_1 \) to second order in \( V \), and then forming \( \delta G_1 / \delta V \), the leading nonlinear expression for \( R \) is obtained as (space-matrix notation; imaginary times in \([0, \tau]\)):

\[
iR(t_3, t_2) = iR_0(t_3, t_2) + G_0(t_3, t_2) \left( \int_0^\tau dt' G_0(t_2, t') V(t') G_0(t', t_3^+) \right)
+ \left( \int_0^\tau dt' G_0(t_3, t') V(t') G_0(t', t_2) \right) G_0(t_2, t_3^+).
\]

The first term on the right describes the equilibrium random phase approximation (RPA) density perturbation response function that pervades linear theory [4]: it is the ”ring ”
diagram $iR_0(t_3, t_2) = G_0(t_3, t_2) G_0(t_2, t_3^+)$, yielding the bulk Lindhard polarizability at zero temperature. The two remaining terms on the right provide the first nonlinear corrections. To first order in the THz potential, $U$, the effective screened potential $V$ may be written in terms of the full (nonlinear) screening function, $K$, as

$$V = K U$$

and approximated as $K \rightarrow K_0^\tau$, where $K_0^\tau$ is the retarded equilibrium screening function, such that

$$V \rightarrow K_0^\tau U \quad \text{and} \quad V(t) = \int_{-\infty}^{\infty} dt' K_0^\tau(t-t') U(t') = U_0 e^{-i\omega_E t} K_0^\tau(\omega_E).$$

Performing the analytic continuation from $[0, \tau]$ to the real time axis (but retaining matrix notation for the positional/momentum variables), we analyze the result for the retarded $\mathcal{R}^\tau$ diagram that $K$ and approximated as $K_0^\tau$. Furthermore, considering

$$U(t_3) = U_0 \exp(-i\omega_E t_3)$$

Employing the appropriate time/frequency transforms we obtain (superscripts $r$, $a$, $>$, $<$ indicate retarded, advanced, greater, lesser functions, respectively):

$$iR^\tau(t_3, t_2) =$$

$$iR_0^\tau(t_3 - t_2) + U_0 G_0^\tau(t_3 - t_2) \left[ \int_{-\infty}^{\infty} dt' G_0^a(t_2 - t') e^{-i\omega_E t'} K_0^\tau(\omega_E) G_0^a(t' - t_3^+) \right]$$

$$\quad + U_0 G_0^\tau(t_3 - t_2) \left[ \int_{-\infty}^{\infty} dt' G_0^\tau(t_2 - t') e^{-i\omega_E t'} K_0^\tau(\omega_E) G_0^\tau(t' - t_3^+) \right]$$

$$\quad + U_0 \left[ \int_{-\infty}^{\infty} dt' G_0^\tau(t_3 - t') e^{-i\omega_E t'} K_0^\tau(\omega_E) G_0^a(t' - t_2) \right] G_0^\tau(t_2 - t_3^+),$$

where $K_0^\tau(\omega_E)$ is in frequency representation. Furthermore, considering $\int_{-\infty}^{\infty} dt'...$ on the right of Eq. (24), we have $(\eta, \beta = r, a, >, <)$

$$= \int_{-\infty}^{\infty} dt' \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} e^{-i\omega(t_2 - t')} G_0^\tau(\omega) e^{-i\omega_E t'} K_0^\tau(\omega_E) \int_{-\infty}^{\infty} \frac{d\bar{\omega}}{2\pi} e^{-i\bar{\omega}(t' - t_3^+)} G_0^\beta(\bar{\omega})$$

$$= \int_{-\infty}^{\infty} d\omega \int_{-\infty}^{\infty} \frac{d\bar{\omega}}{2\pi} e^{-i\omega t_2} \delta(\omega - \omega_E - \bar{\omega}) e^{i\omega t_2} G_0^\tau(\omega) K_0^\tau(\omega_E) G_0^\beta(\bar{\omega})$$
The nonlinear density perturbation response function,

\[ \delta \rho(\omega) = e^{-i\omega t_2} \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} e^{-i\omega(t_2-t_3)} G^0_0(\omega_E + i\omega) K^r_0(\omega_E) G^\delta_0(\omega). \]  

(25)

In the long time limit, the principal contribution comes from low frequencies \( \omega_E > \omega \rightarrow 0 \), and we obtain a time-oscillatory result as

\[ \int_{-\infty}^{\infty} dt' \cdots = e^{-i\omega t_2} G^0_0(\omega_E) K^r_0(\omega_E) G^\delta_0(t_2-t_3). \]

(26)

Correspondingly, the long time steady state result for \( R^r(t_3,t_2) \) is [5]

\[ iR^r(t_3,t_2) = iR^r_0(t_3,t_2) + U_0 e^{-i\omega t_2} G^\delta_0 \left( t_3-t_2 \right) G^0_0(\omega_E) K^r_0(\omega_E) G^\delta_0 \left( t_2-t_3^+ \right) \]

\[ + U_0 e^{-i\omega t_2} G^\delta_0 \left( t_3-t_2 \right) \left\{ G^0_0(\omega_E) K^r_0(\omega_E) G^\delta_0 \left( t_2-t_3^+ \right) + G^\delta_0(\omega_E) K^r_0(\omega_E) G^\delta_0 \left( t_2-t_3^+ \right) \right\} \]

\[ + U_0 e^{-i\omega t_3^+} G^\delta_0(\omega_E) K^r_0(\omega_E) G^\delta_0 \left( t_3^+ - t_2 \right) G^0_0 \left( t_2 - t_3^+ \right) \]

\[ + U_0 e^{-i\omega t_3^+} G^\delta_0(\omega_E) K^r_0(\omega_E) G^\delta_0 \left( t_3^+ - t_2 \right) G^\delta_0 \left( t_2 - t_3^+ \right), \]  

(27)

where we have noted that \( G^\delta_0 \left( t_3^+ - t_2 \right) G^\delta_0 \left( t_2 - t_3^+ \right) \equiv 0 \). (\( \omega \)-arguments indicate frequency representation and \( t \)-arguments indicate direct time representation). Since \( K^r_0 \) is the equilibrium screening function it exhibits resonant response when \( \omega_E \) matches a plasma/normal-mode frequency, and this resonant increase will be transmitted to \( \delta p_2/\delta V \) as well as \( R \) and \( \varepsilon_{v-d} \), increasing the dielectric screening denominator and, correspondingly, reducing the shielded impurity scattering potentials, partially disabling them and giving rise to a resonant source-drain current increase.

4. Conclusions; nonlinear screening and THz photoconductivity resonances

The nonlinear density perturbation response function, \( R \), involves the equilibrium inverse dielectric function, \( K_0(\omega_E) \), in frequency \( (\omega) \) representation. Although the construction of the nonlinear inverse dielectric function (space-time matrix inverse) is nontrivial, one can expect that the long-time behavior associated with the shielding of impurity scattering potentials (involved in steady-state conduction linear in the source-drain electric field) will be strongly affected by resonant behavior of the equilibrium inverse dielectric screening function, \( K_0(\omega_E) \), in the nonlinear terms of the polarizability when \( \omega_E \) approaches a plasmon resonance frequency. Clearly, this results in a resonant increase in the nonlinear polarizability, \( \alpha(1,2) \), and, with that, a resonant increase of the nonlinear direct dynamic dielectric function, \( \varepsilon_{v-d}(t_1,t_2) \) and corresponding decrease of its space-time matrix inverse, \( K_{v-d}(1,2) \), the variational-differential screening function. In turn, this causes a sharp drop of the resistive shielded impurity scattering potentials at these frequencies with attendant resonant increase of dc source-drain current. This amounts to a partial disabling of the impurity scattering mechanism by plasmon resonant behavior in nonlinear screening. Moreover, one can expect important magnetic field effects in semiconductor transport since the cyclotron frequency, \( \omega_c \), can be competitive with the plasma frequency in determining such resonances.

Needless to say, this analysis is applicable to any dielectric medium having plasma (or other) resonances, and is not restricted to any particular geometry. Thus far, the attendant resonant dc transport phenomenology has been observed in double quantum wells and single quantum wells.
We note in closing that while this discussion of resonant effects can accommodate conduction increases of 10%-20% (as observed), it cannot be pressed much further since the iteration approximation would break down as \( V(K_0) \) gets large. Another limitation to be noted is that the full screening function, \( K \), differs from the variational differential screening function, \( K_{v-d} \), upon which the present analysis is based:

\[
V(1) = \int d^4 2 \ K(1,2) U(2).
\]  

Differentiating with respect to \( U(2) \), we have

\[
K(1,2) = K_{v-d}(1,2) - \int d^4 \frac{\delta K(1,3)}{\delta U(2)} U(3).
\]

Iteration of this

\[
K(1,2) \approx K_{v-d}(1,2) - \int d^4 \frac{\delta K_{v-d}(1,3)}{\delta U(2)} U(3) + ...\]

shows that further nonlinear effects, beyond those arising from the nonlinear structure of the density-density correlation function embodied in \( K_{v-d} \), involve even higher-order density correlation functions (starting with the density-density-density correlation function, a 3-particle Green’s function) which will not be addressed in this work.

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Figure 1
Schematic of the Experimental Configuration

Figure 2
Terahertz photoresponse as a function of gate voltage at four different frequencies. The temperature was $T = 25$ K and the source-drain current was $I_{SD} = 100 \, \mu A$. Grating period is 4 $\mu m$. 
Figure 3

Terahertz photoreponse at 570 GHz at four different temperatures (solid). Grating period is 4 μm. Schematic cross section of the devices (inset).

Figure 4
Figure 5
Calculated absorption spectra for the composite structure at various THz frequencies.

Figure 6
Amplitudes of charge (solid curves) and current (dash-dotted curves) density oscillations over a period of the density-modulated single 2DEG for plasmon mode n=2 at frequency 570 GHz.

Measured terahertz photoresponse as a function of gate voltage at four different frequencies (T = 25K).
Figure 7

Schematic of the density modulated DQW system with perforation of the density distribution in the upper QW under the gate fingers.

Figure 8
Figure 9
Full dots mark normalized photoconductance

Figure 10
Amplitudes of the in-plane current density oscillations in the upper QW (solid curve) and in the lower QW (dash-dotted) for the second ($n=2$) plasmon mode.
Amplitudes of the in-plane charge density oscillations in (a) upper and (b) lower 2D electron strips for $d=19.8$ nm (solid curves) and $d=27$ nm (dashed curves).
Amplitudes of the normal-to-plane interwell electric field oscillations at the center plane between the two layers of 2D electron strips for $d=19.8$ nm (solid curves) and $d=27$ nm (dashed curves). The amplitude of the interwell electric field is normalized to the amplitude of electric field in the incident terahertz wave.