The "metallic" regime of holes in GaAs/AlGaAs heterostructures corresponds to densities where two splitted heavy hole bands exist at a zero magnetic field. Using Landau fan diagrams and weak field magnetoresistance curves we extract the carrier density in each band and the interband scattering rates. The measured inelastic rates depend Arrheniusly on temperature with an activation energy similar to that characterizing the longitudinal resistance. The "metallic" characteristics, namely, the resistance increase with temperature, is hence traced to the activation of inelastic interband scattering. The data are used to extract the bands dispersion relations as well as the two particle-hole excitation continua. It is then argued that acoustic plasmon mediated Coulomb scattering might be responsible for the Arrhenius dependence on temperature. The absence of standard Coulomb scattering characterized by a power law dependence upon temperature is pointed out.

Non-interacting two dimensional electron gas (2DEG) is believed to be insulating in the sense that its resistance always diverges as the temperature, T, approaches zero. This trend is opposite to that characteristic of most 3D metals where for T → 0 the resistance becomes smaller and eventually saturates to a finite value. The insulating nature of 2D systems has been observed in many experiments [1,2] and was a generally accepted dogma until a few years ago, when Kravchenko et al. [3] discovered that in silicon based high mobility 2DEG at high enough carrier densities, the resistance decreases as T → 0 ("metallic" behavior). The same samples, at lower densities, displayed insulating characteristics and the crossover from positive to negative variation of the resistance with temperature was soon identified as a novel metal-insulator transition. Since the discovery of Kravchenko et al. [3], qualitatively similar dependencies of the resistance upon temperature were observed in two dimensional hole gas (2DHG) in GaAs/AlGaAs heterostructures [4,5], SiGe quantum wells [6], InAs quantum wells [8] as well as other silicon samples [9].

While the nature of the insulating phase may be roughly understood, the physics of the metallic phase remains obscure. A mere reduction of the resistance with temperature does not however necessarily imply delocalization. Within a Drude-like framework it may result from a temperature dependent carrier density, or as suggested by Altshuler and Maslov [11], by a temperature dependent scattering time.

Here we provide experimental evidence that indeed suppression of scattering with decreasing T is responsible for the metallic characteristics of 2DHG in GaAs/AlGaAs heterostructures. The scattering mechanism, interband Coulomb scattering between the two splitted heavy hole bands, is very different from the one proposed in ref. [10]. The correlation between the existence of two conducting bands and "metallic" behavior was put forward by Pudalov [11]. It has recently been convincingly substantiated by Papadakis et al. [12].

We find the same correlation but go beyond that and prove that the characteristic dependence upon temperature, \( \rho_x(T) = \rho(0) + \alpha \exp(-T_0/T) \) \((T_0\) and \( \alpha \) are some constants) follows from a similar dependence of the interband inelastic scattering rates upon temperature. We extract the bands' structures and their particle-hole excitation continua from the measurements. We then use the latter to show that the Arrhenius temperature dependence might result from activation of plasmon mediated interaction, similarly to plasmon enhanced Coulomb drag between coupled layers [13].

We believe that the same considerations may apply to the other "metallic" 2D systems since band degeneracy is generally lifted there due to spin-orbit coupling and the lack of inversion symmetry at the interface where the 2D layer resides.

The splitting between the two heavy hole bands in GaAs/AlGaAs heterostructures due to spin orbit coupling and lack of inversion symmetry have been extensively studied both theoretically [14,15] and experimentally [16,17]. The two bands are approximately degenerate up to a certain energy where they split and acquire very different effective masses (see inset to fig. 4). Thus, above a threshold density, current is carried by two bands as reflected in slope variation in the corresponding Landau fan diagram [16] or in the appearance of a second frequency in Shubnikov de Haas measurements [17]. Using either method, the hole densities in the lighter band, \( n_l \), and heavier band, \( n_h \), can be extracted [17] and fig. 1 depicts them as a function of the total density, \( n_{total} \).

For a total density below \( 1.7 \times 10^{11} cm^{-2} \), the bands are degenerate. For higher densities they split and for a total density above \( 2.8 \times 10^{11} cm^{-2} \), practically all additional carriers go to the heavier and less mobile band. The inset to fig. 1 depicts a characteristic Shubnikov de Haas curve and demonstrates the existence of two sets of oscillations corresponding to the two bands. The data were taken using a high mobility (\( \simeq 500,000 cm^2/V \cdot s \)
at 100 mK) 2DHG confined to a GaAs/Al_{0.8}Ga_{0.2}As interface in the 100 plane. The sample had a 2DEG front gate and silicon doped backgate, 40mm and 300nm from the 2DHG, respectively.

The simultaneous transport of two types of carriers with different mobilities and densities is also manifested in our system by a weak field classical positive magnetoresistance \( \varepsilon \). A set of resistance curves for \( \rho_{\text{total}} = 4.25 \times 10^{-11} \text{cm}^2 \) for different temperatures is depicted in the inset to fig. 2. Measurements as a function of density show that the positive magnetoresistance at weak fields appears when the total density is beyond the split off density, namely, when the two bands have different masses and mobilities. It then grows larger with density, as the bands deviate further. The Lorentzian shaped magnetoresistance expected from two band transport is obtained by subtracting the weak parabolic negative magnetoresistance (attributed to Coulomb interactions \( B \)) from the full magnetoresistance curve. It is depicted in fig. 2. Note the excellent agreement with the predicted shape, eqs. 2 below. Below \( \approx 0.6K \), the resistance is practically independent of \( T \) while for temperatures above \( \approx 2K \), the Lorentzian is hardly visible. The suppression of the classical two band magnetoresistance results from interband scattering. At low temperatures this scattering is mainly elastic. As the temperature is increased, inelastic scattering commences, the drift velocities of carriers in the two bands gradually approach each other, and the magnetoresistance is consequently diminished.

The data presented in figs. 1 and 2 can be used to extract the interband scattering rates. To that end, the standard two band transport formulae \( B \) should be generalized to include interband scattering. The starting point is coupled Drude equations for transport in the two bands which straightforwardly give the resistance tensor

\[
\begin{bmatrix}
\varepsilon_x \\
\varepsilon_y \\
\varepsilon_x \\
\varepsilon_y
\end{bmatrix}
= 
\begin{bmatrix}
S_i & -H R_i & -Q & 0 \\
H R_i & S_i & 0 & -Q \\
-Q & 0 & S_h & -H R_h \\
0 & -Q & H R_h & S_h
\end{bmatrix}
\begin{bmatrix}
J_{ix} \\
J_{iy} \\
J_{hx} \\
J_{hy}
\end{bmatrix}
\]

(1)

Here, \( \varepsilon \) and \( H \) are the electric and magnetic fields, respectively, \( R_i = 1/\varepsilon n_i \) is the Hall coefficient of the \( i \)-th band, \( n_i \) is the hole density in that band and \( J_i \) is the current density there. The elements \( S_i \), \( S_h \), and \( Q \) can be expressed in terms of conductances, \( S_i = \sigma_{ii}^{-1} + \sigma_{ih}^{-1} ; S_h = \sigma_{ih}^{-1} + \sigma_{hh}^{-1} ; Q = \lambda \sigma_{ih}^{-1} = \lambda^{-1} \sigma_{hh}^{-1} \) where \( \lambda \) is a function of the velocities and densities. The diagonal conductances, \( \sigma_{ii} \), pertain to scattering (elastic and inelastic) within each band while \( \sigma_{ij} \) accounts for interband scattering. The latter processes may include carrier transfer from one band to the other as well as drag-like processes where one particle from one band scatters off a particle in the other band and both carriers maintain their bands. The function \( \lambda \) depends on the nature of the dominant interband scattering mechanism.

Setting \( J_{iy} + J_{hx} = 0 \) one obtains the Lorentzian shape longitudinal magnetoresistance depicted in fig. 2

\[
\rho_{xx}(H) = \frac{i x}{J_{ix} + J_{hx}} = \rho_{xx}(H \to \infty) + \frac{L}{1 + (H/W)^2}
\]

(2)

The hole densities in the two bands, and hence \( R_i, R_h \), are known from fig. 1. Fitting the data depicted in fig. 2 to eqs. 2, one obtains \( \rho_{xx}(H \to \infty), L \), and \( W \) which are in turn used to extract \( S_i, S_h \), and \( Q \) as a function of \( \lambda \). The Landau fan diagram combined with the low field magnetoresistance, hence provide a unique opportunity to directly measure inter and intraband scattering. The results of such analysis are shown in fig. 3 where \( S_i, S_h, Q \), and \( \rho_{xx}(H = 0) \) are depicted vs. \( T \) for the same total density as in fig. 2. At low temperatures, all these quantities saturate to some residual values which we attribute to inter and intraband elastic scattering. As the temperature is increased, inelastic scattering commences and these quantities grow.

The remarkable and central result of our work is the observation that the inelastic scattering rates follow the same temperature dependence as \( \rho_{xx}(H = 0) \), namely, \( S_i(T) = S_i(0) + \alpha S_i \exp(-T_0/T) \), \( Q(T) = Q(0) + \alpha Q \exp(-T_0/T) \), where \( \alpha S_i, \alpha Q \) are some constants. Remarkably, the characteristic temperature, \( T_0 \), is similar to all these quantities, including \( \rho_{xx} \). For \( S_i \) and \( S_h \), we find \( T_0 = 5.0K \), for \( Q \), \( T_0 = 4.8K \), and for \( \rho_{xx} \), \( T_0 = 4.3K \). Similar correlations are found for other total densities in the "metallic" regime. Our results hence strongly suggest that the universal Arrhenius temperature dependence of the resistance in the "metallic" regime merely reflects the increase of inelastic interband scattering with temperature and probably have nothing to do with phase transition into a delocalized state.

As expected, the light band is more susceptible to scattering and hence \( \alpha S_i > \alpha S_h \). Note the inelastic contributions to \( S_i \) and \( Q \) at \( T = 2K \) are larger than their elastic counterparts. Moreover, they are all larger than \( \rho_{xx} \), indicating the two bands are strongly coupled by the scattering. Eventually, for \( T > 2K \), the coupling equilibrates the two drift velocities and the classical magnetoresistance is fully suppressed. However, the resistance, \( \rho_{xx} \), continues to grow. The latter growth results from inelastic interband scattering that affects the resistance even when the bands are already fairly strongly coupled. In fact, in some of the data published in the literature, the Arrhenius dependence of the resistance is not accompanied by the Lorentzian magnetoresistance, probably...
indicating substantial interband scattering.

We find experimentally that $S(T) \simeq S_0(0) + 2.1Q; S_h(T) \simeq S_h(0) + 0.48Q$, thus yielding for the density of figs 2, 3, $\lambda = 0.48$. Since the prefactors multiplying $Q$ are reciprocal, the resistances, $S_0(0); S_h(0)$, are identified as $\sigma^{-1}_l$ and $\sigma^{-1}_h$, respectively. The diagonal resistances, pertaining to intraband scattering, are hence found to be practically temperature independent. The Arrhenius $T$ dependence originates from interband scattering alone.

We turn now to discuss possible reasons for the Arrhenius temperature dependence of $S_l, S_h,$ and $Q$ which in turn lead to a similar temperature dependence of $\rho_{xx}$. These scattering rates (expressed as resistances) crucially depend on the bands’ dispersions relations, $E_i(k)$, and their resulting excitation spectra. To extract the bands dispersions depicted in the inset to fig. 4, we approximate the light band by a parabolic relation with a mass $m_l = 0.38m_0$ ($m_0$ is the bare electron mass). The variation of $p_l, p_h$ with $p_{total}$ depicted in fig. 1, is then used to calculate the ratio between the two bands compressibilities. Neglecting band warping as well as differences between density of states and compressibility, we use the ratio of the two compressibilities to extract the dispersion of the heavy band. This dispersion then allows, within the random phase approximation, the calculation of the excitation spectrum of the system. The spectrum is composed of two particle-hole continua, one for each band, and two plasmon branches. The particle-hole continua correspond to regions in the $q, \omega$ plane where the imaginary part of the polarization operator, $\Pi(q, \omega)$, is non-zero. The plasmon branches are the poles of the dielectric constant, $\epsilon(q, \omega)$. Both are shown in fig. 4 for zero temperature. Due to the very different masses, the optical plasmon branch corresponds mostly to motion of the light holes while the acoustic branch originates from the heavy holes screened by the lighter ones (analogous to acoustic phonons in metals). At small wave-vectors, the acoustic plasmon velocity is $v_{ap} = \sqrt{\frac{m_h}{2m_l}} v_{Fh}$, where $v_{Fh}$ is the heavier hole Fermi velocity.

Some interband scattering may be attributed to electron-phonon interaction but the calculated magnitude of this effect is more than an order of magnitude too small to account for the measured rates. A more plausible candidate is interband Coulomb scattering that leads to resistance through either Coulomb drag or particle transfer.

The rate of interband scattering is proportional to the screened potential squared, $|\frac{2\pi e^2}{q\epsilon(q, \omega)}|^2$. In the absence of plasmons, the dielectric function is regular and at low temperatures may be approximated by its static value. The resulting rate is usually proportional to $T^2$. The absence of such contribution in our data is puzzling and calls for a detailed calculation of the scattering rates with the specific particle-hole continua and band structure depicted in fig. 4. It may happen that the very different masses, as well as the concave shape of the heavy band, limit that contribution to small values.

The Coulomb scattering rates may be substantially enhanced in the presence of plasmons due to the divergent screened interaction in their vicinity, provided the plasmon branch overlaps the particle-hole continua. This effect is very pronounced in the calculation of the Coulomb drag between coupled quantum wells by Flensberg and Hu [13]. As indicated by fig. 4, at $T = 0$ the acoustic plasmon does not intersect the heavy holes continuum. As the temperature is raised, $\text{Im}[\Pi_h(q, \omega)]$ is thermally activated beyond its zero $T$ boundaries, leading to a finite overlap with the acoustic plasmon branch. The value of $\text{Im}[\Pi_h(q, \omega)]$ is Arrheniusally small there but the diverging screened interaction compensates for that. The corresponding scattering rates depend Arrheniusly on temperature.

We turn now to evaluate the corresponding activation temperature, $T_0$. Since the temperature, $T \leq 2K$ is small compared with the Fermi energy, we restrict ourselves to small $q$. A direct calculation for a concave dispersion relation then gives, $\text{Im}[\Pi_h(q, v_{ap}q)] = \frac{k_p}{2\sqrt{2\pi e^2} \sqrt{\hbar^2 E/\partial k^2}} \exp[(-E_h - E_p)/T]$. Here, $k_p < k_{Fh}$ is the momentum for which the heavy hole velocity matches the plasmon velocity, the curvature, $\partial^2 E/\partial k^2$, is evaluated at $k_p$ and $E_p \equiv E(k_p)$. Note the result is independent of $q$! The activation temperature, $T_0$, is simply $E_h - E_p$. At our density, $m_h = 5.15m_l$, leading to $v_{ap} \simeq 1.6v_{Fh}$. The corresponding $E_p$ is marked in the inset to fig. 4 together with the Fermi energy. The resulting activation energy $T_0 = E_h - E_p \simeq 5.5K$, in good agreement with the value characterizing the inelastic rates, $Q, S_l, S_h$, and the resistance, $\rho_{xx}$. We are presently calculating the resulting scattering rates.

We briefly note another possible source of Arrhenius temperature dependence of the resistance. As the temperature is increased, carriers are transferred from the light to the heavy band, due to the larger entropy of the latter (larger density of states). This coexistence of two bands is hence analogous to the standard liquid (light band) - gas (heavy band) coexistence. Under certain conditions, the density changes calculated from the corresponding Clausius-Clapeyron equation may follow an Arrhenius dependence which is analogous to the vapor pressure in the liquid-gas problem. Since the heavy hole mobility is lower than that of the light holes, such carrier transfer should result in resistance increase. We are presently studying the implications of this effect.

In summary. We have shown experimentally that the “metallic” behavior of the resistance of holes in a GaAs/AlGaAs heterostructure results from inelastic scattering between the two splitted heavy hole bands. The measured interband scattering rates depend Arrheniusly on temperature with almost the same activation energy as the longitudinal resistance. Using Shubnikov
de Haas data, we mapped the band structure and the corresponding excitations. We found that due to the dispersion relation of the heavy band, the system supports a weakly damped acoustic plasmon. The Arrhenius small overlap between the heavy hole excitation continuum and the plasmon branch leads to a similar dependence of the interband scattering rates, and hence of $\rho_{xx}$ upon $T$. Using the measured band structures we calculate $T_0$ for the plasmon mediated process and obtain good agreement with the values deduced experimentally from the inelastic scattering rates. The absence of a substantial power law contribution to the inelastic interband scattering remains to be explained.

Acknowledgment
This work was supported by the Binational Science Foundation (BSF), Israeli Academy of Sciences, German-Israeli DIP grant, Minerva foundation, Technion grant for promotion of research, and by the V. Ehrlich career development chair.

Figure Captions
Figure 1 - Hole densities of the two splitted heavy hole bands as a function of total density. Inset - one of the Shubnikov de Haas traces used to determine the hole densities.

Figure 2 - The two-band classical magnetoresistance contribution to the resistance for different temperatures. Note the perfect Lorentzian shape. Inset - Full magnetoresistance curves for the same temperatures.

Figure 3 - The various scattering rates expressed in terms of resistances (left axis) and the zero field longitudinal resistance (right axis) vs. $T$. Solid lines depict best fit to Arrhenius dependence. Inset - Full magnetoresistance curves for the same temperatures.

Figure 4 - Solid lines - heavy and light particle-hole excitation continua as a function of momentum scaled to the heavy hole Fermi wave vector. Shaded area corresponds to the range where drag-like interband scattering is possible at very low $T$. Dashed lines - optical (op) and acoustic (ap) plasmon dispersions. Inset - The measured bands dispersion relations. The energy $E_F$ corresponds to the hole Fermi energy and $E_p$ is the energy where the heavy hole velocity matches that of the acoustic plasmon. The difference, $E_F - E_p$ gives the activation temperature, $T_0$ (see text).

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Figure 1

$P_1, P_h \text{ [cm}^2\text{]}$

$P_{\text{total}} \text{ [cm}^2\text{]}$

$T = 100\text{mK}$

$P_{\text{sat}} = 3.65 \times 10^{-1}$

$P_{\text{total}}$ [cm$^2$] $F_{\text{total}}$ [cm$^2$]

Figure 1
Figure 2

- Experimental data
- Fit to eq. (2)

$\rho_{xx} - \rho_{\text{background}}$

$P_{\text{total}} = 4.25 \times 10^{11}$
Figure 4