Transverse Magnetosonic Waves and Viscoelastic Resonance in a Two-Dimensional Highly Viscous Electron Fluid

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In high-mobility materials, conduction electrons can form a viscous fluid at low temperatures. We demonstrate that in a high-frequency flow of a two-dimensional electron fluid in a magnetic field the two types of excitations can coexist: those of the shear stress (previously unknown transverse magnetosound) and those associated with the charge density (conventional magnetoplasmons). The dispersion law and the damping coefficient of transverse magnetosound originate from the time dispersion of the viscosity of the fluid. Both the viscoelastic and the plasmonic components of the flow exhibit the recently proposed *viscoelastic resonance* that is related to the own dynamics of shear stress of charged fluids in a magnetic field. We argue that the generation of transverse magnetosound, manifesting itself by the viscoelastic resonance, is apparently responsible for the peak in photoresistance and peculiarities in photovoltage observed in ultrahigh-mobility GaAs quantum wells.

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Introduction.—In materials with sufficiently weak disorder, quasiparticles can form viscous fluids provided that the interparticle momentum-conserving collisions are much more intensive than any other collisions which do not conserve momentum [1,2]. The hydrodynamic regime of phonon transport in liquid helium and dielectrics was studied in detail several decades ago [3]. However, only recently the hydrodynamic regime of charge transport was discovered in novel materials: graphene, Weyl semimetals, and the best-quality GaAs quantum wells [4–16]. These discoveries were accompanied by an extensive development of theory [17–35].

High-frequency dynamics of a viscous electron fluid in a zero magnetic field was theoretically considered in Ref. [36] within the Landau Fermi-liquid model. It was shown that, in the region of frequencies ω and wave vectors q satisfying the inequality $\omega \gg v_F q$ (v_F is the Fermi velocity), the dynamics of the fluid is described by the classical viscoelastic motion equations. A high-frequency flow of a viscous 2D electron fluid in a long sample in zero magnetic field was theoretically studied in Refs. [30,31]. The flow was shown to contain the plasmonic part formed by plasmons and the viscoelastic part formed by transverse zero sound [30,31]. The latter can be excited in a Fermi liquids with a sufficiently strong interaction between quasiparticles [37]. Transverse zero sound and plasmons in 2D strongly nonideal Fermi liquids in the absence of a magnetic field were considered in Refs. [32,33].

Hydrodynamic equations of high-frequency dynamics of a 2D electron fluid in a magnetic field were formulated in Refs. [23,34,35]. It was shown that the ac viscosity coefficients at the frequency ω, equal to the doubled cyclotron frequency $2\omega_c$, exhibit the *viscoelastic resonance*, originating from the own rotation of the shear stress tensor in a magnetic field [35]. The effects of viscosity on 2D magnetoplasmons were studied in Ref. [34,35].

In this Letter, we predict transverse zero magnetosound in a highly viscous (strongly nonideal) 2D electron fluid. Such waves are related to perturbations of shear stress of a charged Fermi liquid in a magnetic field. We use the hydrodynamic approach, which is valid at a strong interparticle interaction [38], in order to calculate the dispersion law of transverse magnetosound and the linear response of the fluid to a radio-frequency electric field. In sufficiently narrow samples, the response is formed predominantly by the standing waves of transverse magnetosound at frequencies $\omega > 2\omega_c$ and exhibits the viscoelastic resonance at $\omega = 2\omega_c$. Depending on the sample width, the resonance manifests itself in the ac conductance corresponding to the linear response or, vice versa, in the ac impedance.

We discuss the giant peak in photoresistance and the peculiarities in photovoltage that were recently observed in ultrahigh-mobility GaAs quantum wells at the frequencies near $\omega = 2\omega_c$ [39–41]. The explanation [16] of the giant negative magnetoresistance, discovered in similar GaAs quantum wells [4–8], demonstrates that 2D electrons in such structures form a viscous fluid. Here we provide the evidences that the high-frequency magnetotransport phenomena reported in Refs. [39–41] are explained by the excitation of transverse magnetosound in the 2D electron fluid. We also show that independence of the flow in narrow samples on the sign of the circular polarization of radiation, obtained in our theory, correlates with measurements of photoresistance of similar structures [42–45].

High-frequency hydrodynamics of 2D electron strongly nonideal Fermi liquid.—A high-frequency flow of a 2D electron fluid is described by the particle density $n(\mathbf{r},t) = n_0 + \delta n(\mathbf{r},t)$ and the hydrodynamic velocity $\mathbf{V}(\mathbf{r},t)$ [here $\mathbf{r} = (x,y)$ is the coordinate in the 2D layer and n_0 is the equilibrium density]. We decompose $\delta n(\mathbf{r},t)$ and $\mathbf{V}(\mathbf{r},t)$ by time harmonics proportional to $e^{-i\omega t}$ with the complex amplitudes $\delta n(\mathbf{r})$ and $\mathbf{V}(\mathbf{r})$. In the regime linear by the perturbation of the density and by the velocity, the continuity equation and the Navier-Stokes equation in a magnetic field $\mathbf{B} = B\mathbf{e}_z$ have the form [35]:

$$-i\omega\delta n + n_0 \mathrm{div}\mathbf{V} = 0, \tag{1}$$

$$-i\omega\mathbf{V} = e\mathbf{E}/m + \omega_c \mathbf{V}\times\mathbf{e}_z + \eta_{xx}\Delta\mathbf{V} + \eta_{xy}\Delta\mathbf{V}\times\mathbf{e}_z, \tag{2}$$

where the hydrodynamic pressure term, $-\nabla P/m$, and the bulk momentum relaxation term, $-\gamma\mathbf{V}$ are omitted [46]; $\mathbf{E} = \mathbf{E}(\mathbf{r})$ is the complex amplitude of the time harmonic of an ac electric field $\mathbf{E}(\mathbf{r},t)$, e and m are the electron charge and mass, and the ac viscosity coefficients $\eta_{xx} = \eta_{xx}(\omega)$ and $\eta_{xy} = \eta_{xy}(\omega)$ are [35]

$$\left.\begin{matrix}\eta_{xx}\\ \eta_{xy}\end{matrix}\right\} = \frac{\eta_0}{1 + (-\omega^2 + 4\omega_c^2)\tau_{ee}^2 - 2i\omega\tau_{ee}} \left\{\begin{matrix}1 - i\omega\tau_{ee}\\ 2\omega_c\tau_{ee}\end{matrix}\right. . \tag{3}$$

Here η_0 is the viscosity the absence of magnetic field and τ_{ee} is the electron-electron scattering time. If $\omega, \omega_c \gg 1/\tau_{ee}$, then the viscosity coefficients (3) exhibit the viscoelastic resonance at $\omega = 2\omega_c$ [35].

The electric field $\mathbf{E}(\mathbf{r},t)$ consists of the field of incident radiation $\mathbf{E}_0(t)$ and the internal field $\mathbf{E}_{\mathrm{int}}(\mathbf{r},t)$ is induced by a perturbation of the electron density $\delta n(\mathbf{r},t)$. We do not consider the retardation effects that are important in the region of small wave vectors in ungated structures (see Refs. [57,58]). Thus for $\mathbf{E}_{\mathrm{int}}$ we use the electrostatic equation: $\mathbf{E}_{\mathrm{int}} = -\nabla\delta\varphi$. For the structures with a metallic gate located at the distance d from the 2D layer the perturbation of the electrostatic potential is

$$\delta\varphi = (4\pi e d/\kappa)\delta n, \tag{4}$$

where κ is the background dielectric constant.

A "true hydrodynamic" flow is characterized by the almost equilibrium distributions of electrons by their velocities in the coordinate systems moving with the velocities $\mathbf{V}(\mathbf{r},t)$. Such a flow is described by the Navier-Stokes equation (2) and is formed under the condition that the interparticle scattering length $l_{ee} = v_F\tau_{ee}$ is the shortest spacescale. However, the hydrodynamiclike description within Eq. (2) can be also applicable when a fluid flow is driven by a high-frequency electric field with the frequencies $\omega \gg 1/\tau_{ee}$ and, therefore, the quasiequilibrium distribution of electrons in the moving frame does not have enough time to be formed. More precisely, if the characteristic spacescale Δx of flow inhomogeneities is much greater than the cyclotron radius, $R_c = v_F/\omega_c$, or the path that a free electron passes during the period of $\mathbf{E}_0(t)$, $l_\omega = v_F/\omega$, the first and second angular harmonics dominate in the distribution of electrons by velocities and Eq. (2) turns out to be valid [2,16,36].

If the interaction between electrons is weak and electrons can be regarded as an almost ideal Fermi gas, transverse sonic waves, in which $\mathbf{V}\perp\mathbf{q}$, cannot propagate in the system [37,59]. In this case the dc zero-field viscosity coefficient η_0 in Eq. (3) has the form $\eta_0 = v_F^2\tau_{ee}/4$ [60].

Excitation of transverse zero sound becomes possible if the interparticle interaction is sufficiently strong and 2D electrons form a strongly nonideal highly viscous Fermi liquid [37,59]. Dynamics of an electron Fermi liquid in the absence of magnetic field was theoretically studied in Ref. [36]. It was shown that the eigenmodes of the Fermi liquid allow the hydrodynamiclike description provided by $l_\omega \ll \Delta x$. This inequality is surely realized for transverse zero sound in a strongly nonideal electron liquid, in which the parameter r_s, characterizing the strength of the Coulomb interaction, is much greater than unity [36].

The analysis [38], following Refs. [36,61,62], demonstrates that, in a magnetic field, a strong interaction between quasiparticles also justifies the applicability of hydrodynamics for the description of the transverse eigenmode. In Ref. [38], Eq. (2) was consistently derived from the kinetic equation for strongly interacting Fermi-liquid quasiparticles (the Landau interaction parameters F_0 and F_1 are much greater than unity). Herewith the value of the viscosity coefficient η_0 becomes proportional to $(1+F_0)(1+F_1)$ [38]; thus the parameter $v_F^\eta = 2\sqrt{\eta_0/\tau_{ee}}$ becomes much larger than the Fermi velocity v_F.

For the structures investigated in Refs. [39–41], the interparticle interaction parameter r_s is about unity. The recent experiments [63–66] evidence that the effective mass of 2D electrons in similar high-mobility GaAs quantum wells is substantially renormalized due to the interparticle interaction. These two facts indicate that the conditions $F_{0,1} \gtrsim 1$ or, even possibly, $F_{0,1} \gg 1$ seem to be fulfilled for 2D electrons in the structures examined in Refs. [39–41] and that the hydrodynamic model (1)–(4) is applicable for a qualitative description of ac magnetotransport in them [67].

Formation of the linear response.—First, we study the wave solutions of the hydrodynamic equations in the absence of an external field $\mathbf{E}_0(t)$.

Substitution of $\mathbf{V}(\mathbf{r}) = \mathbf{V}_0 e^{i\mathbf{q}\cdot\mathbf{r}}$ and $\delta n(\mathbf{r}) = \delta n_0 e^{i\mathbf{q}\cdot\mathbf{r}}$ into Eqs. (1), (2), and (4) leads to the algebraic equations for the amplitudes $\mathbf{V}_0$ and δn_0 and for the eigenvalues $\omega_{1,2}(q)$ (see their exact form in the Supplemental Material [46]). We suppose that the parameters of the system satisfy the inequalities presented in Table I. Apparently, such conditions can be fulfilled, in some degree, for 2D electrons in the best-quality samples of high-mobility GaAs quantum wells. In this case, the largest (at $q \gg \omega_c/s$) of the

TABLE I. Values of the parameters allowing the formulated hydrodynamic description.

Strong interparticle interaction:	$v_F^\eta \gg v_F$
High-frequency regime:	$\omega_c \sim \omega \gg 1/\tau_{ee}$
Rigid spectrum of plasmons:	$s/v_F^\eta \gg \omega\tau_{ee}$
Microscopically wide samples:	$W \gg \omega/v_F$
Weak bulk disorder:	$\gamma/\omega \ll (v_F^\eta/s)^2/(\omega\tau_{ee})$

two eigenvalues $\omega_{1,2}(q)$ describes the magnetoplasmon waves. Neglecting all the relaxation processes ($1/\tau_{ee} \to 0$ and $\gamma \to 0$), we obtain the well-known dispersion law of magnetoplasmons: $\omega_p(q) = \sqrt{\omega_c^2 + s^2q^2}$, where $s = 2\sqrt{\pi e^2 n_0 d/m\kappa}$.

Damping of magnetoplasmons arises due to viscosity as well as due to a finite bulk momentum relaxation with the rate γ. At $\gamma = 0$ the damping coefficient takes the following form [35]:

$$\Upsilon_p(q) = \frac{\omega_c^2 + \omega_p^2}{2\omega_p^2} q^2 \mathrm{Re}\eta_{xx} + \frac{\omega_c}{\omega_p} q^2 \mathrm{Im}\eta_{xy}. \qquad (5)$$

Here the viscosity coefficients $\eta_{xx}(\omega)$ and $\eta_{xy}(\omega)$ are taken at $\omega = \omega_p(q)$. The damping of magnetoplasmons due to scattering on disorder in high-quality samples with a very small rate γ is important only in the very vicinity of the cyclotron resonance, $|\omega - \omega_c| \lesssim \gamma$ [46].

The other of the eigenvalues $\omega_{1,2}(q)$ corresponds to the magnetosonic waves, whose amplitude $\mathbf{V}_0$ is perpendicular to the wave vector $\mathbf{q}$. Such transverse waves are due to perturbations of the shear stress tensor and are analogous to transverse sound in amorphous solids. The dispersion law of magnetosound $\omega_s(q)$ and its damping coefficient $\Upsilon_s(q)$ originate from the time dispersion of the diagonal viscosity $\eta_{xx}(\omega)$. At high frequencies, ω_c, $\omega \gg 1/\tau_{ee}$, one obtains the following [46] [see also Fig. 1(a)]:

$$\omega_s(q) = \sqrt{4\omega_c^2 + \frac{(v_F^\eta)^2 q^2}{4}}, \qquad \Upsilon_s(q) = \frac{4\omega_c^2 + \omega_s^2}{2\omega_s^2 \tau_{ee}}. \qquad (6)$$

As it was discussed above, for applicability of the current hydrodynamic approach one needs for $v_F^\eta \gg v_F$. Therefore the magnetosound wavelength, having the order of magnitude $l_s = v_F^\eta/\omega$ at $\omega_c \sim \omega$, is much larger than the path that a quasiparticle passes during one period of an ac field, $l_\omega = v_F/\omega$. So we can consider the flows with the minimal spacescale Δx as small as $v_F/\omega \ll \Delta x \ll v_F^\eta/\omega$.

In order to study how magnetoplasmons and transverse magnetosound can be excited, one needs to consider a spatially nonuniform flow that is driven by an uniform ac electric field $\mathbf{E}_0(t)$ in a finite-size sample. So, second, we calculate the linear response of a fluid in a long sample on an circularly polarized ac field $\mathbf{E}_0(t) = \mathbf{E}_0 e^{-i\omega t} + \text{c.c.}$

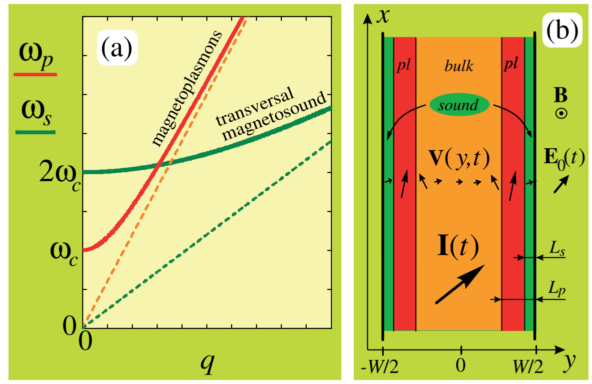

FIG. 1. (a) Dispersion laws of magnetoplasmons $\omega_p(q)$ and transverse magnetosound $\omega_p(q)$. Dashed lines show the dispersion laws in zero magnetic field: $\omega_p^0(q) = sq$ and $\omega_s^0(q) = v_F^\eta q/2$ (here $s \gg v_F^\eta$; see Table I). At $q \to 0$ the frequency $\omega_s(0) = 2\omega_c$ corresponds to the own rotation of the shear stress tensor of the charged fluid in magnetic field. (b) Long sample in the external fields $\mathbf{E}_0(t)$ and $\mathbf{B}$ in the regime $\omega > 2\omega_c \gg 1/\tau_{ee}$. In wide samples, $W \gg L_p, L_s$, the three regions are formed: the central region (orange) where the flow is controlled by the bulk momentum relaxation, the near-edge layers (dark-green) with the widths L_s in which the flow is governed by viscosity, and the regions $L_s < |y - W/2| < L_p$ (red) where the magnetoplasmon contribution dominates.

[Fig. 1(b)]. Here $E_{0,x} = E_0/2$, $E_{0,y} = \mp iE_0/2$, and the signs "−" and "+" correspond to the right and the left circular polarizations. We suppose the longitudinal sample edges to be rough that leads to the no-slip boundary conditions: $\mathbf{V}(y = \pm W/2) = 0$.

From Eqs. (1)–(4) we obtain a system of linear differential equations for the complex amplitudes of the velocity $\mathbf{V}(y) = [V_x(y), V_y(y)]$ [46]. In the main order by the parameter $v_F^\eta/s \ll 1$, the solution of the equations is

$$\mathbf{V}(y) = \frac{eE_0}{2m}[\mathbf{A}_0^\pm + h(\lambda_p, y)\mathbf{A}_p^\pm + h(\lambda_s, y)\mathbf{A}_s], \qquad (7)$$

where the first term $\mathbf{A}_0^\pm = [i, \pm 1]/(\tilde{\omega} \pm \omega_c)$ describes the bulk Drude contribution, the second and the third terms with the amplitudes $\mathbf{A}_p^\pm = \pm[i\omega_c/\omega, -1]/(\tilde{\omega} \pm \omega_c)$ and $\mathbf{A}_s = [-i, 0]/\omega$ are the plasmonic and the viscoelastic contributions, the value $\tilde{\omega} = \omega + i\gamma$ takes into account a weak scattering on disorder, the function $h(\lambda, y) = \cosh(\lambda y)/\cosh(\lambda W/2)$ describes the profile of the flow, the eigenvalue $\lambda_p = iq_p - \omega\Upsilon_p(q = q_p)/(s^2 q_p)$ corresponds to magnetoplasmons ($q_p = \sqrt{\omega^2 - \omega_c^2}/s$ is the magnetoplasmon wave vector), and the eigenvalue $\lambda_s = \sqrt{-i\omega/\eta_{xx}(\omega)}$ corresponds to magnetosound.

If the parameters of the system satisfy the conditions of Table I, the layers formed by magnetoplasmons and magnetosound are separated in space and have the widths $L_p = 1/|\mathrm{Re}\lambda_p|$ and $L_s = 1/|\mathrm{Re}\lambda_s|$ of the different orders of magnitudes [see Fig. 1(b) and the Supplemental Material [46]].

It is noteworthy that the viscoelastic part of Eq. (7) is independent on the sign of the circular polarization.

From Eq. (7) we get the result for the complex amplitude $\mathbf{I} = (I_x, I_y) = en_0 \int_{-W/2}^{W/2} \mathbf{V}(y)dy$ of the ac current:

$$\mathbf{I} = \frac{e^2 n_0 E_0 W}{2m}[\mathbf{A}_0^{\pm} + f_p \mathbf{A}_p^{\pm} + f_s \mathbf{A}_s], \tag{8}$$

where $f_{p,s} = \tanh(\lambda_{p,s}W/2)/(\lambda_{p,s}W/2)$.

Properties of the linear response.—Equations (7) and (8) indicate that the Drude part of the flow dominates in very wide samples, $W \gg L_p$, while magnetoplasmons give the main contribution to $\mathbf{I}$ in moderately wide samples, $l_p \ll W \ll L_p$ ($l_p = s/\omega$ is the characteristic wavelength of magnetoplasmons). In the last case, the standing waves of magnetoplasmons are formed in the y direction and the viscoelastic resonance manifests itself via the width of the magnetoplasmon resonances [see Eq. (5) and details in the Supplemental Material [46]]. In the current work, we are mainly interested in the case of medium and narrow samples, $W \ll l_p$, where the viscoelastic contribution provides a substantial or even the main contribution to the linear response.

In medium samples, $L_s \ll W \ll l_p$, at the frequencies above the viscoelastic resonance, $\omega > 2\omega_c$, the eigenvalue λ_s is mainly imaginary. Correspondingly, the flow profile $\mathbf{V}(y)$ in the near-edge regions, $W/2 - |y| \lesssim L_s$, is the standing waves of transverse magnetosound. Below the resonance, $\omega < 2\omega_c$, the eigenvalue λ_s becomes mainly real. Thus the viscoelastic part of the flow acquires an exponential profile and is located in the narrower near-edge layers $W/2 - |y| \lesssim l_s$, (here $l_s \sim v_F^\eta/\omega \ll L_s$ is the characteristic wavelength of magnetosound). Because of such a change of the flow, the dependence $\mathbf{I}(\omega_c)$ drastically modifies at $\omega_c = \omega/2$ and the viscoelastic resonance arises (Fig. 2). The absorbtion power $\mathcal{W} = 2\mathrm{Re}(\mathbf{E}_0^* \mathbf{I})$ exhibits a peak, asymmetric relative to the point $\omega_c = \omega/2$ [46]:

$$\mathcal{W}(\omega_c) \propto \mathrm{Re}\left(\frac{1}{\sqrt{i - (2\omega_c - \omega)\tau_{ee}}}\right). \tag{9}$$

In narrow samples, $l_s \ll W \ll L_s$, the standing waves of magnetosound are formed in the whole sample at $\omega > 2\omega_c$. In this case the plasmonic contribution is relatively small. When the sample width is equal to a half-integer number of magnetosound wavelengths, $W = (2m+1)\pi/q_s(\omega_c)$, resonances at $\omega_c = \omega_c^{s,m}$, related to amplification of the standing magnetosound waves, appear in the dependence $\mathcal{W}(\omega_c)$ [see Eq. (8) and Fig. 2; $q_s = 2\sqrt{\omega^2 - 4\omega_c^2}/v_F^\eta$ is the magnetosound wave vector].

In narrowest samples, $W \ll l_s$, the velocity amplitude $\mathbf{V}(y)$ has a parabolic profile in a whole sample [see Eq. (7)].

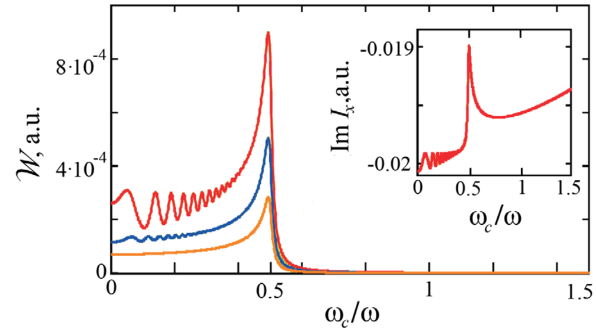

FIG. 2. Absorbtion power $\mathcal{W}$ as a function of ω_c at fixed ω for medium and narrow samples with the widths $W/l_{ee} = 5.7$, 3.2, 1.8 (orange, blue, and red curves) at $\omega\tau_{ee} = 50$ and $s/v_F^\eta = 300$. Inset presents the imaginary part of the ac current $I_x(\omega_c)$ for a medium sample ($W/l_{ee} = 1.8$). The viscoelastic resonance at $\omega_c = \omega/2$ for all curves and the "magnetosonic resonances" at $\omega_c = \omega_c^{s,m}$ for red curves are seen.

This regime can be regarded as a high-frequency Poiseuille flow in magnetic field. The viscoelastic resonance manifests itself not in the current, but in the ac impedance $Z = E_0/I_x$. Near the resonance Eq. (8) yields the Lorentzian profile of the real part of the impedance:

$$\mathrm{Re}Z(\omega_c) \propto \frac{1}{1 + (2\omega_c - \omega)^2 \tau_{ee}^2}. \tag{10}$$

At $\omega\tau_{ee} \gg 1$ the imaginary part of the current (8) does not depend on any relaxation parameter [46]. This corresponds to ac field–induced elastic oscillations of the highly viscous electron fluid ("electron honey"), glued to the longitudinal sample edges at $y = \pm W/2$.

Manifestation of the viscoelastic resonance in nonlinear effects.—In high-mobility GaAs quantum wells, bright surprising effects were observed at the frequencies near $\omega = 2\omega_c$ in nonlinear ac magnetotransport: a strong peak in photoresistance [39,40], and a peculiarity in the photovoltaic effect [41]. Below we present the arguments that the viscoelastic resonance is responsible for these effects.

It was stressed in Ref. [39] that the strong peak ("spike") in photoresistance and a very well pronounced giant negative magnetoresistance are observed in the *same best-quality* GaAs structures. In Ref. [16] the giant negative magnetoresistance was explained as a manifestation of forming a viscous flow of a 2D electron fluid. Thus peculiarities in any ac transport effect in those structures near the frequency $\omega = 2\omega_c$, actually observed in experiments [39–41], must inevitably have an explanation within the hydrodynamic model.

In Ref. [41] the photovoltage effect was measured on several GaAs quantum well structures of different geometries. The peculiarity at $\omega = 2\omega_c$ was much better observed in the sample having a meander-shaped gate, compared with the uniformly gated and the ungated samples. Apparently, the meander-shaped gate induces inhomogeneous perturbations of an electron flow. Thus, in the meander-gated sample the role of shear stress and viscosity must be greater than in the other samples and the viscoelastic resonance is expected to be more pronounced, as it was actually observed.

There can exist various versions of the hydrodynamic theories of the photoresistance and the photovoltaic effects [46], in which the nonlinear response originates from different sources of nonlinearity. In any case, a nonlinear response contains a linear response as a block and, thus, must exhibit the peak or another peculiarity at $\omega = 2\omega_c$. We have shown in this work that the peak of the viscoelastic resonance in the linear response can be asymmetric as well as symmetric, depending on the sample width. Both these two types of the peak, apparently, were observed in photoresistance in Refs. [39,40] on different samples.

In the experiments [42–45] it was demostrated that photoresistance of high-mobility GaAs quantum wells is usually almost independent of the sign of circular polarization of radiation at $\omega_c \lesssim \omega$. Although the giant peak in photoresistance was observed up to now only for the linear polarization of ac field, the independence of the viscoelastic part of the linear response (8) on the sign of the circular polarization and our analysis of the results of experiments [42–45] correlates with the statement that the peak has the hydrodynamic origin (see the Supplemental Material [46]).

Conclusion and acknowledgments.—We have pointed out that transverse magnetosonic waves, accompanied by the viscoelastic resonance, can be excited in a highly viscous electron fluid and, apparently, have been observed in GaAs quantum wells in experiments [39–41].

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