Surface acoustic wave induced magnetoresistance oscillations in a 2D electron gas

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We study the geometrical commensurability oscillations imposed onto the resistivity of 2D electrons in a perpendicular magnetic field by a propagating surface acoustic wave (SAW). We show that, for \( \omega < \omega_c \), this effect is composed of an anisotropic dynamical classical contribution increasing the resistivity and the non-equilibrium quantum contribution isotropically decreasing resistivity, and we predict the appearance of zero-resistance states associated with geometrical commensurability at large SAW amplitude. We also describe how the commensurability oscillations modulate the resonances in the SAW-induced resistivity at multiples of the cyclotron frequency.

High-mobility two-dimensional electron gases (2DEGs) display interesting effects under intense microwave irradiation [1] or the influence of surface acoustic waves in the microwave frequency range [6, 7]. One class of effects is induced magneto-oscillations that originate from the geometric commensurability between the cyclotron radius \( R_c \) and the period \( 2\pi/q \) of spatially periodic perturbations [10], which has been observed in statically modulated 2D systems [11, 12, 15]. This effect was recently studied in the attenuation and renormalization of surface acoustic wave (SAW) velocity due to interactions with electrons [11, 14] and in the drag effect [1]. For a spatially periodic propagating field of the SAW, the commensurability effect can also be viewed as the resonant SAW interaction with collective excitations of 2D electrons at finite wavenumbers [3, 11], enabling one to excite modes otherwise forbidden by Kohn’s theorem [10].

In this Letter, we study the non-linear dynamical effect in which SAWs induce changes in the magneto-resistivity of a high quality electron gas in the regime of classically strong magnetic fields, \( \omega_c \tau \gg 1 \) and high temperatures \( k_BT \gg \hbar \omega_c \). We show that the resistivity changes reflect both frequency and geometrical resonances in the surface acoustic wave attenuation and are formed from two competing contributions.

The first contribution originates from the SAW-induced guiding center drift of the cyclotron orbits – a purely classical effect. For a SAW with frequency \( \omega \), and wavenumber \( q \), propagating in the \( x \) direction with speed \( s = \omega/q \), there is an anisotropic increase in the resistivity \( \rho_{xx} \) (at high fields \( \omega_c \tau \gg 1 \), this is equivalent to an increase of conductivity in the transverse direction \( \sigma_{yy} \)), which oscillates as a function of inverse magnetic field when the Fermi velocity \( v_F \gg s \). We show that at \( \omega \gg \omega_c \) the resistivity change displays resonances at multiples of the cyclotron frequency \( \omega \approx N\omega_c \).

The second contribution arises from the modulation of the electron density of states (DOS), \( \tilde{g}(\epsilon) = [1 - \Gamma \cos(2\pi\epsilon/\hbar\omega_c)] \gamma \) (where \( \gamma = m/\pi \hbar^2 \)), and consequently, from the energy dependence of the non-equilibrium population of excited electron states caused by Landau level quantization. We follow the idea proposed in Ref. [6] to explain the formation of zero-resistance states [1, 2, 4] under microwave irradiation with \( \omega \gg \omega_c \). We show that in the frequency range \( \tau^{-1} \ll \omega \ll \omega_c \), the quantum contribution suppresses resistivity both in \( \rho_{xx} \) and \( \rho_{yy} \) and persists up to temperatures \( k_BT \gg \hbar \omega_c \) and filling factors \( \nu \gg 1 \) where no Shubnikov-de Haas oscillations would be seen in the linear-response conductivity.

When \( \tau^{-1} \ll \omega \ll \omega_c \), the commensurability oscillations take the form

\[
\frac{\delta \rho_{xx}}{\rho_{xx}} \approx 2J_0^2(qR_c) \mathcal{E}^2 \left[ \frac{\nu_s}{\nu_c} - \frac{\tau_{in}}{\tau} \omega \right] (2\pi \nu \gamma)^2, \]

\[
\frac{\delta \rho_{yy}}{\rho_{xx}} \approx -2J_0^2(qR_c) \mathcal{E}^2 \times \frac{\tau_{in}}{\tau} (2\pi \nu \gamma)^2, \tag{1}
\]

with \( \mathcal{E} = e\chi \omega_{SAW}^2 / \nu_F \), \( \omega_{SAW}^2 \) the SAW longitudinal electric field, \( \omega_{SAW} = \chi / [2\pi\nu^2 \gamma] \) the 2D screening radius, \( \nu_F \) the Fermi energy, \( \tau \) and \( \tau_{in} \) the momentum and inelastic relaxation times [17], and \( \nu = 2\nu_F / \hbar \omega_c \), the filling factor. The geometrical oscillations in Eq. (1) are described by the Bessel function \( J_0 \). For different sample parameters and measurement conditions, the observed oscillatory change in the resistivity can be dominant in the SAW propagation direction [\( \eta \equiv (2\pi \nu_s / \gamma \nu_F) \sqrt{\tau_{in}/\tau} < 1 \), isotropic \( \eta \gg 1 \), or only in the component perpendicular to the SAW wavevector [\( \eta \approx 1 \)].

We now present our detailed analysis, starting with the dynamical classical contribution. At high magnetic fields \( \omega_c \tau \gg 1 \), the resistivity change \( \delta \rho_{xx} \) can be tracked back to the SAW induced drift \( Y(t) \) (along the \( y \)-axis) of the guiding centre of an electron cyclotron orbit [12] and the resulting enhancement of the transverse \( (y) \) component of the electron diffusion coefficient,

\[
\frac{\delta \rho_{xx}}{\rho_{xx}} = \frac{\delta \mathcal{D}_{yy}}{\mathcal{D}_{yy}} \sim \left( \left( \frac{Y}{R_c} \right)^2 \right).
\]

The drift is caused by an electric field \( E_{q,x}\cos(qx - \omega t + \phi) \mathbf{x} \), with \( x(t) = R_c \sin (\omega t - \psi) \). Between two impurity
scattering events the guiding centre is displaced by
\[ Y(t) \sim \frac{eE_{\text{SAW}}}{m_F} \int_0^t dt' \cos \left[ qR_c \sin(\omega_c t' - \psi) - (\omega t' - \phi) \right]. \]
A particular guiding centre displacement \( Y(t) \) depends on the initial phase \( \psi \) of the electron revolution along the cyclotron orbit and the phase \( \phi \) of the SAW field which change randomly each time electron scatters form impurities, if \( R_c > 2q/q_t \), thus leading to a random change in the value and direction of \( Y \). The mean square value of such a displacement (averaged over \( \psi \), \( \phi \) and the drift time between two scattering events),
\[ \langle Y^2 \rangle = \int_0^\infty \frac{dt}{\tau} e^{-t/\tau} \int_0^{2\pi} \int_0^{2\pi} \frac{d\psi}{2\pi} \frac{d\phi}{2\pi} Y^2(t), \]
can be used to obtain the frequency and wave number dependence of the effect,
\[ \frac{\delta^c \rho_{xx}}{\rho_{xx}} = 2(qlE^2) \sum_{N=-\infty}^{\infty} \frac{J_N^2(qR_c)}{1 + (\omega - N\omega_c)^2 \tau^2}, \]
where \( E = e\alpha \frac{E_{\text{SAW}}}{\sqrt{\epsilon_F}} \) and \( l = v_F\tau \). Equation \( 3 \) includes the Thomas-Fermi screening of the SAW field by 2D electrons, \( E_{\text{SAW}} = q\alpha \frac{e^2 \epsilon_{\text{SAW}}}{\sqrt{\epsilon_F}} \). A typical change in the magnetoresistance \( \delta^c \rho_{xx} \) for the regime \( qR_c \gg \omega/\omega_c \gg 1 \) (possible at normal electron densities for which \( \tau = v_F \gg s \)) is illustrated in Fig. 5 and compared with the Weiss oscillations in a static potential.

There are a sequence of resonances at integer multiples of the cyclotron frequency, \( \omega \approx N\omega_c \). The oscillations for even harmonics are in phase with the Weiss oscillations of the static potential, whilst those of odd harmonics are \( \pi \) out of phase. This is because the main contributions to the drift occur when the electrons are moving parallel (or anti-parallel) to equipotential lines. For odd harmonics the phase of the potential at the half-orbit point is opposite to that for a static potential, and hence cancellation and reinforcement effects that lead to minima and maxima in the resistance are interchanged. In the regime \( \omega_c \tau \gg 1 \), the resonances are very narrow and appear to display a random sequence of heights, reflecting the dependence on the geometric resonance conditions.

In the intermediate frequency domain, \( \tau^{-1} \ll \omega \ll \omega_c \), a natural regime for GaAs structures with densities \( n_c \gtrsim 10^{10} \text{cm}^{-2} \) at sufficiently high magnetic fields, the classical oscillations take the form
\[ \frac{\delta^c \rho_{xx}}{\rho_{xx}} \approx \frac{2v_F^2 \epsilon^2 \omega^2 J_0^2(qR_c)}{s^2}. \]
Due to the competition between electron screening effects (\( \propto q^2 \)), the dynamical suppression of commensurability by the SAW motion (\( \propto \omega^{-2} \)), and the relation \( \omega/q = s \), the form of these oscillations is independent of the absolute value of the SAW frequency, provided that conditions \( \tau^{-1} \ll \omega \) and \( v_F \gg s \) are satisfied.

The dynamical mechanism just described dominates in a classical electron gas. No redistribution of electron kinetic energy (due to SAW absorption) will additionally change the magnetoresistance until the 2DEG is heated to a temperature \( k_B T_e \gtrsim h\omega_c k_F/q \) where geometrical oscillations become smeared. The essential assumption leading to this statement is that electron single-particle parameters (velocity and \( \tau^{-1} \)) vary slowly with energy at the scales comparable to the Fermi energy and can thus be approximated by constants.

To demonstrate this point, we analyze the classical kinetic equation for a 2DEG at temperature \( k_B T \lesssim h\omega_c k_F/q \) irradiated by SAWs. We solve the kinetic equation for the electron distribution function,
\[ f(t, x, \phi, \epsilon) = f_T + \int_{\omega_{\text{SAW}}} e^{-i\omega_{\text{SAW}} t + i\epsilon x} \int_{m} f_{m}^{f}(\epsilon) e^{imf}, \]
using the method of successive approximations. Here, \( f_T(\epsilon) \) is the homogeneous equilibrium Fermi function, and the angle \( \phi \) and kinetic energy \( \epsilon \) parametrize the electron state in momentum space. Each component \( f_{m}^{f} \) describes the \( m \)-th angular harmonic of the time- and space-dependent non-equilibrium distribution. To describe local values of the electron current and the accumulated charge density, we use the energy-integrated functions, \( g_{\omega_{\text{SAW}}}^{m} = \int_0^\infty d\epsilon f_{m}^{f} \). The relaxation of the local non-equilibrium distribution towards a Fermi function characterized by the value of local Fermi energy, \( \epsilon_F(t, x) \) (determined by the local electron density \( n(t, x) \propto g_0^0(t, x) = \int_0^\infty d\epsilon f_{0}^{0}(t, x) \), where \( f_{0}^{0}(t, x) = \int_0^\infty d\epsilon_{F}^{0}(t, x) \)) is described in the relaxation-time approximation by
\[ \dot{\tilde{L}} f = -\frac{f - f^{0}}{\tau} - \frac{f^{0} - f_T(\epsilon - \epsilon_F(t, x))}{\tau_{in}}, \]
\( \hat{\mathcal{L}} = \partial_t + v \cos \varphi \partial_x + \left[ \omega_c - \frac{eE}{p} \sin \varphi \right] \partial_x + evE \cos \varphi \partial_x, \)

where we distinguish between the elastic scattering rate \( \tau^{-1} \) and energy relaxation rate \( \tau^{-1}_{\text{in}} \). The electric field \( E \) in Eq. (6) is the combination of a homogeneous DC field \( E_{00} \) and the screened electric field of the SAW, \( E(t, x) = \sum_\omega E_{\omega q} e^{iqx - i\omega t} \hat{x} \), found from the unscreened SAW field via \( E_{\omega q} = E_{\text{SAW}} / \kappa(q, \omega) \), where \( \kappa(q, \omega) \) is the dielectric function of the whole 2D structure.

The dynamical perturbation of the distribution function can be found from time/space Fourier harmonics of Eq. (6) at the frequency/wave number of the SAW,

\[
\begin{align*}
|\partial_\varphi + \frac{1}{\omega_c} & - i \frac{\omega}{\omega_c} + i q R_c \cos \varphi | f_{\omega q} = \Psi(\varphi), \\
\Psi = - & \frac{\delta^{0}}{\omega_c} (\partial_t f_{\omega q}) + \frac{\tau^{-1} - \tau^{-1}_{\text{in}}}{\omega_c} f_{\omega q}^0 + \frac{eE_{\omega q}}{\omega_c} \left[ \cos \varphi \partial_x - \frac{\sin \varphi}{p} \partial_x \right] (f_{00} + f_{\omega}),
\end{align*}
\]

where we include the unknown perturbation of the time/space averaged function, \( f_{00} \), related to the DC current to lowest order in \( E_{\omega q} \). Equation (7) can be formally solved using the Green function \( G(\varphi, \tilde{\varphi}) \)

\[
f_{\omega q}(\varphi) = \int_0^\varphi G(\varphi, \tilde{\varphi}) \Psi(\tilde{\varphi}) d\tilde{\varphi},
\]

\[
G(\varphi, \tilde{\varphi}) = e^{\frac{\delta^{0}}{\omega_c} (\varphi - \tilde{\varphi}) + i q R_c [\sin \varphi - \sin \tilde{\varphi]},
\]

which allows for an infinite range of variation of \( \varphi \) whilst guaranteeing periodicity of the solution \( f_{\omega q}(\varphi) \).

Thomas-Fermi screening of the SAW field is produced by the density modulation \( n_{\omega q} = \gamma g_{00}^{\omega q} \) via the induced field \( E_{\omega q}^{\text{ind}} = - \frac{2e q \gamma}{\omega} \hat{E}_{\omega q}^{\text{SAW}} \), so that \( E_{\omega q} = E_{\omega q}^{\text{SAW}} + E_{\omega q}^{\text{ind}} \). In the analysis of screening, the DC part of the electric field can be ignored (\( f_{00} = 0 \)), and self-consistency yields

\[
g_{0\omega}^{\omega q} = \frac{eE_{\omega q} 1 - (1 - i \omega \tau) K}{1 - K},
\]

\[
K = \int_0^{2\pi} \frac{d\varphi}{2\pi} \int_0^{\varphi} \frac{d\tilde{\varphi}}{\omega_c} G(\varphi, \tilde{\varphi}) = \sum_{N=-\infty}^{\infty} \frac{J^2_N(q R_c)}{1 + i \tau (N \omega_c - \omega)},
\]

where \( K(\omega, q) \) obeys the relations \( K(\omega, q) = K^*(-\omega, q) = K(\omega, -q) \) and is obtained from Eqs. (8) using the identity \( e^{\pm i \sin \varphi} = \sum_{N=-\infty}^{\infty} J_N(z) e^{\pm i N \varphi} \), and

\[
\kappa = 1 + \frac{1}{\omega_c |q|} - \frac{1 - (1 - i \omega \tau) K}{1 - K}.
\]

To find the steady state current, we analyze the time/space average of the kinetic equation in Eq. (6) and take into account the dynamical perturbation \( f_{\omega q} \),

\[
\begin{align*}
|\partial_\varphi + \frac{1}{\omega_c} & - i \frac{\omega}{\omega_c} + i q R_c \cos \varphi | f_{\omega q} = \Psi(\varphi), \\
\Psi = - & \frac{\delta^{0}}{\omega_c} (\partial_t f_{\omega q}) + \frac{\tau^{-1} - \tau^{-1}_{\text{in}}}{\omega_c} f_{\omega q}^0 + \frac{eE_{\omega q}}{\omega_c} \left[ \cos \varphi \partial_x - \frac{\sin \varphi}{p} \partial_x \right] (f_{00} + f_{\omega}),
\end{align*}
\]

We substitute the solution Eq. (6) into Eq. (13), keeping track of the effect of the perturbation of the time/space averaged function \( f_{00} \) on \( f_{\omega q} \). We thus include SAW-induced non-linear effects. We multiply Eq. (13) by \( (2\omega_c/v)e^{-i\varphi} \), integrate with respect to \( \epsilon \) and \( \varphi \), then use the relation between the \( x \) and \( y \) components of the DC current, \( j_x - ij_y = e \gamma \nu_\parallel g_{00}^{\omega q} \) and the harmonic \( g_{00}^{\omega q} \) (note that electrical neutrality requires \( g_{00}^{\omega q} = 0 \), to get

\[
\begin{align*}
2\omega_c & \int \frac{dk}{2\pi} e^{-i\varphi} \int d\epsilon \left\{ \frac{1}{\omega_c} |v f_{00} + \sum_\omega \frac{v E_{\omega q}}{\omega_c} \left[ \cos \varphi \partial_x - \frac{\sin \varphi}{p} \partial_x \right] f_{\omega q} \right\} \\
& = \frac{2\omega_c e}{v} \nu_\parallel \gamma \int \left( \frac{2 \omega_c \nu_\parallel \gamma^{\omega q}}{v} \right) \left[ j_x - ij_y \right] = E_{\omega q}^{\nu_\parallel} - i E_{\omega q}^{\nu_\parallel},
\end{align*}
\]

which can be used to determine the SAW induced change of the resistivity tensor, \( \delta \rho \). The relation between the electric field and current is \( \mathbf{E} = \rho \mathbf{j} + \delta \rho \mathbf{j} \), where \( \rho \) is the Drude resistivity tensor, thus we find the resistivity corrections

\[
\delta \rho_{xx} = \frac{1}{2} \sum_\omega \left| \frac{\epsilon l E_{\omega q}}{\epsilon F} \right|^2 \Re \left\{ \frac{K}{1 - K} \right\},
\]

\[
\delta \rho_{yy} = 0,
\]

\[
\delta \rho_{xy} = \frac{1}{2} \sum_\omega \left| \frac{\epsilon l E_{\omega q}}{\epsilon F} \right|^2 \Im \left\{ \sum_\omega \frac{K}{1 - K} \right\} = 0,
\]

which, with the use of \( E_{\omega q} = E_{\omega q}^{\text{SAW}} / \kappa(q, \omega) \) and Eq. (14), yields the result in Eq. (15).

The magnetic field dependence of the resistivity change reflects the form of the SAW attenuation by the 2D electrons determined by the real part of the longitudinal dynamical conductivity \( \sigma_{\omega q} \),

\[
\Re \sigma_{\omega q} = \gamma s^2 \tau e^2 \Re \left\{ \frac{K}{1 - K} \right\}.
\]

The finite wavenumber \( q > 8 \omega_{\text{H}} / R_c^2 \) of the SAW allows the system to bypass Kohn’s theorem and generates absorption resonances at \( \omega = \omega_{\text{H}} + \Delta N(q) \),

\[
\left\langle \mathbf{E} \cdot \mathbf{j} \right\rangle = \frac{\gamma |e q a_{\text{sc}} E_{\omega q}^{\text{SAW}}|^2 s^2 \tau}{1 + (2 (\omega - \omega_{\text{H}} + \Delta N))^2 J^2_N(q R_c)}. \]

Absorption of SAWs by the 2D electrons changes their steady state distribution over energy, though for energy-independent characteristics this does not lead to additional changes in magnetoresistance beyond those described in Eq. (15). However, Landau level quantization, which is unavoidable in a phase-coherent electron system.
gas when \( \omega_c \tau \gg 1 \), gives rise to the quantum contributions to the geometrical commensurability oscillations which persist up to high temperatures. The oscillatory energy-dependence of the electron DOS, \( \tilde{\gamma}(\epsilon) \), imposes oscillations on the electron elastic scattering rate, \( \tau^{-1}(\epsilon) = \tau^{-1} \tilde{\gamma}/\gamma \), and the contribution to the observable conductivity \[^3\] is

\[
S(\epsilon) \approx \sigma_{xx} \tilde{\gamma}^2/\gamma^2, \quad \sigma = \int \mathrm{d}\epsilon S(\epsilon) [-\partial_{\epsilon} f^0_{\uparrow \downarrow}].
\]

At low temperatures, \( k_B T \lesssim \hbar \omega_c \), the DOS oscillations lead to Shubnikov-de Haas oscillations in conductivity. At high temperatures \( k_B T \gg \hbar \omega_c \), thermal broadening smears out oscillations, but the quantum contribution can remain in non-linear effects after energy averaging.

It was shown in Ref. \[^6\] that the electron energy distribution acquires an oscillatory part via the availability of final states for energy absorption processes. When the electron gas is excited by SAWs, carriers are redistributed between energy intervals with varying transport efficiencies, changing the overall resistivity of the 2DEG in proportion to the SAW attenuation. We extend the study of the balance equation performed in Refs. \[^5, 6\] for the non-equilibrium electron distribution (at \( \omega \gg \omega_c \)) to low-frequencies \( \omega \ll \omega_c \). To lowest order in \( |E_{\omega q}|^2 \), the oscillatory part is

\[
\delta f^0 = \tau_{1n} \sum_{\pm \omega q} |E_{\omega q}|^2 \frac{\sigma_{\omega q}}{\gamma^2} \partial_\epsilon \left[ \frac{\tilde{\gamma}^2}{\gamma^2} \partial_\epsilon (\xi_{00} + \xi_{0T}) \right],
\]

where \( \sigma_{\omega q} = \gamma \Gamma (2\pi/\hbar \omega_c) \sin(2\pi/\hbar \omega_c) \) appears under the square, so that its averaged broad thermal smearing gives \( \langle \partial_\epsilon \tilde{\gamma}^2/\gamma^2 \rangle = 1/2 \langle \partial_\epsilon \gamma^2/\hbar \omega_c \rangle^2 \). This results in a non-vanishing addition to both diagonal components of the conductivity (and, therefore, also of the resistivity) even when \( k_B T \gg \hbar \omega_c \). This generates isotropic magneto-oscillations

\[
\frac{\delta^q \rho_{\alpha\alpha}}{\rho_{xx}} = \frac{\delta^q \sigma_{\alpha\alpha}}{\sigma_{xx}} = -\frac{2\tau_{1n}}{\tau} \frac{4\pi \Gamma e F^2}{\hbar \omega_c} \xi^2 J^0_0(\xi R_c), \quad (16)
\]
in addition to the anisotropic classical commensurability effect and, together, they yield the result in Eq. \[^{lp}\].

In conclusion, we have demonstrated a new class of magnetoresistance oscillations caused in a 2DEG by SAW. We have shown that the effect consists of contributions with competing signs: (i) a classical geometric commensurability effect analogous to that found in static systems with positive sign, and (ii) a quantum correction, with negative sign. The latter result suggests that SAW propagation through a high mobility electron gas may generate a sequence of zero-resistance states (ZRS) linked to the maxima of \( J_0^2(q R_c) \) for strong enough SAW fields. Whilst this prediction concerns the low-frequency domain \( \omega \lesssim \omega_c \), such ZRS would be formed via the same mechanism \[^{lp}\] as the microwave-induced ZRS at \( \omega \gg \omega_c \). A large enough SAW-induced change \( |\delta \sigma_{xx}| > \sigma_{xx} \) resulting in negative local conductivity would require formation of electric field/Hall current domains. Since the anisotropy in Eq. \[^{lp}\] suggests that such conditions can be achieved the easiest in the conductivity component along the SAW wavevector, we expect that domains would form with current flowing perpendicular to the direction of SAW propagation, and their stability would depend on the sample geometry \[^{lp}\].

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\[^17\] We choose short-range disorder for computational convenience, and expect that long-range disorder will introduce quantitative changes to our results.
\[^18\] \( \Delta_N \approx N \omega_c J_{2n}^0(\xi R_c) / \xi R_c \sim J_{2n}^0(q R_c) - N A_N / \omega_c \), with \( A_N = \sum_{p=1}^N J_{2n+2p}^0(q R_c) - J_{2n}^0(q R_c) \).
\[^19\] We can only treat the quantum correction using the classical \( \sigma_{xx} \) when \( \omega \ll \omega_c \).
For a SAW with the wavevector directed across the axis of a Hall bar, current domains can be stabilized by ending in ohmic contacts. For a wave propagating (or standing) along the Hall bar, current domains would have to orient across the bar direction and terminate at the sample edges (destabilizing them), leading to a finite resistance. Finally, the anisotropy would not support a zero-conductance regime in a Corbino geometry.