










Beatings of ratchet current magneto-oscillations in GaN-based grating gate structures: Manifestation of spin-orbit band splitting

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We report on the study of the magnetic ratchet effect in AlGa_N/Ga_N heterostructures superimposed with a lateral superlattice formed by a dual-grating gate structure. We demonstrate that irradiation of the superlattice with a terahertz beam results in the *direct* ratchet current, which shows giant magneto-oscillations in the regime of Shubnikov–de Haas oscillations. The oscillations have the same period and are in phase with the resistivity oscillations. Remarkably, their amplitude is greatly enhanced as compared with the ratchet current at zero magnetic field, and the envelope of these oscillations exhibits large beatings as a function of the magnetic field. We demonstrate that the beatings are caused by the spin-orbit (SO) splitting of the conduction band. We develop a theory which gives a good qualitative explanation of all experimental observations and allows us to extract the SO splitting constant $\alpha_{\text{SO}} = 7.5 \pm 1.5 \text{ meV\AA}$. We also discuss how our results are modified by plasmonic effects and show that these effects become more pronounced with decreasing the period of the grating gate structures down to submicrons.

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I. INTRODUCTION

One of the most important tasks of modern optoelectronics is to provide efficient conversion of high-frequency terahertz signals into a direct current (DC) electrical response, for reviews see, e.g., Refs. [1–9]. In the last decades, the focus of research in this direction was on the periodic structures like field effect transistor (FET) arrays, grating gate, and multigate structures. Such structures also attract growing interest as simple examples of tunable plasmonic crystals [10–14]. Plasmonic crystals already demonstrated excellent performance as terahertz detectors [15–20], in close agreement with the numerical simulations [21–25]. They are also actively studied as possible emitters or amplifiers of terahertz radiation [26–28].

Importantly, a DC photoresponse requires some asymmetry of the system, which would determine the direction of the produced DC current. Generation of electric DC current in response to an alternating (AC) electric field in systems with broken inversion symmetry is usually called the *ratchet effect*, which was studied both theoretically and experimentally in a great number of systems, for reviews see, e.g., Refs. [29–34]. For effective radiation conversion to DC signal in periodic structures, there should be strong built-in asymmetry inside

the unit cell of the plasmonic crystals. The ratchet DC current can be induced by the electromagnetic wave incident on the spatially modulated system provided that the wave amplitude is also modulated but is phase-shifted in space, for review, see Ref. [30]. On the theoretical level, the ratchet current arises already in noninteracting approximation (so-called *electronic ratchet*). Electronic ratchets were discussed in two-dimensional (2D) systems with lateral gratings [30,35–41] or arrays of asymmetric dots/antidots [42–45]. Electron-electron (ee) interaction can dramatically increase ratchet current due to plasmonic effects [24,46–53].

Although the ratchet effect was treated theoretically and observed experimentally in diverse low-dimensional spatially modulated structures, some basic issues of this effect still remain puzzling. One of the interesting questions that has not yet been discussed in the literature is the manifestation of the effects of spin-orbit (SO) interaction in the ratchet effect. In this paper, we address this question. We study the ratchet effect in a magnetic field (in what follows, we call it *magnetic ratchet effect*) in the regime of Shubnikov–de Haas (SdH) oscillations and demonstrate that it is dramatically modified by SO interaction. Specifically, we report on the observation of the magnetic ratchet effect in the lateral Ga_N-based superlattice formed by a dual-grating gate (DGG) structure. The specific property of the Ga_N systems as compared with other 2D structures including graphene is a very high value of

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Rashba SO coupling (at least 10 times larger than in GaAs-based structures) [54–62], which is caused by a high built-in electric field existing in such polar materials. This is, therefore, very favorable for observations of SO-induced effects. We demonstrate, both experimentally and theoretically, that in quantized magnetic fields, terahertz excitation results in the giant magneto-oscillations of the ratchet current coming from Landau quantization, which due to large SO band splitting, are strongly modulated as a function of the magnetic field. There are two possible approximations which allow one to describe the effect: the hydrodynamic approximation, in which ee collisions dominate, and the drift-diffusion approximation, in which ee interaction is fully neglected (ratchet effect within these two approximations was analyzed in Refs. [24,25,53] and Refs. [30,36,38,39,63], respectively). On a qualitative level, both approximations reproduce the main results, namely, the gigantic enhancement of the ratchet effect in the SdH regime, as well as the SO-induced beating of the SdH oscillations. Here, we use the hydrodynamic approach because it can describe the plasmonic effects in a sufficiently simple way.

Our results provide a method to study the band spin splitting. Currently, the most widely used techniques are direct measurements of magnetoresistance in the SdH oscillation regime [64], the weak antilocalization experiments [65], optical methods [66], and photogalvanic studies [67]. Since SdH oscillations in the magnetoresistance regime and ratchet current correlate, these two measuring methods are complementary, which gives us the opportunity to double check the results.

The possibility to increase the ratchet effect in the magnetic field deserves special attention. Therefore, we start the paper with a discussion of the key points of the magnetic ratchet effect, see Sec. II. The rest of the paper is organized as follows. In Sec. III, we present the experimental results on magnetic ratchet effects in GaN-based structures. In the following Sec. IV, we present the theory and compare its results with the experimental data. Section IV C is devoted to a discussion of the plasmonic effects. Finally, in Sec. V, we summarize the results.

II. RATCHET EFFECT IN MAGNETIC FIELDS: STATE OF THE ART

Physics of the ratchet effect becomes much richer if one applies the magnetic field. Magnetic ratchet effect, which is in some publications called magnetophotogalvanic effect, was widely studied in different semiconductor systems. The magnetic ratchet effect can be induced even in the case of homogeneous graphene with structure inversion asymmetry, see, e.g., Refs. [68–72]. The effect is sensitive to disorder and can be tuned by the gate voltage [72]. Furthermore, the theoretical consideration of the magnetic ratchet effect in graphene and bilayer graphene showed that it can be substantially enhanced under the cyclotron resonance (CR) condition [73]. Most recently, it has been shown theoretically and observed experimentally that the magnetic ratchet effect also can be drastically enhanced by deposition of asymmetric lateral potential introduced by an asymmetric periodic metallic structure on top of a structure with 2D electron gas (2DEG) [63,74–

76]. Remarkably, the magnetic ratchet strongly increases (by more than two orders of magnitude) in the SdH oscillation regime. Physically, this happens due to very fast oscillations of the resistivity with the Fermi energy and, consequently, with the electron concentration. As a result, inhomogeneous (dynamical and static) density modulations induced by the electromagnetic wave lead to a very strong response.

Importantly, the responsivity in the regime of SdH oscillations increases not only in grating gate structures but also in single FETs [77–80]. Although the general physics of enhancement in both cases is connected with fast oscillations of resistivity, there is an essential difference. In grating gate structures, the shape of typical DC photoresponse roughly reproduces resistance oscillations, while in single FETs, the typical response is $\pi/2$, shifted with respect to resistance oscillations. The latter shift was explained theoretically by the hydrodynamic model in Ref. [78] and demonstrated experimentally in Ref. [80]. The key idea is as follows. The transport scattering rate $\gamma(n)$ in the SdH regime sharply depends on the dimensionless electron concentration $n = (N - N_0)/N_0$ (here, N_0 is background concentration, and N is the concentration in the channel). Expanding $\gamma(n) \approx \gamma(0) + \gamma'(0)n$ with respect to small n , one finds that a nonlinear term $\gamma'(0)n\mathbf{v}$ appears in the Navier-Stokes equation, where \mathbf{v} is the drift velocity. This is sufficient to give a nonzero response, which in a single FET arises in the second order with respect to external terahertz field (both n and \mathbf{v} are linear with respect to this field). Therefore, in this case, the response is proportional to the first derivative of $\gamma'(0)$ with respect to concentration (i.e., with respect to the Fermi energy E_F); hence, it is $\pi/2$, shifted in respect to the conductivity oscillations. By contrast, in the grating gate structures, the DC response appears only in the third order with respect to perturbation [30]. As a consequence, the ratchet current is proportional to the second derivative $\gamma''(0)$ (see discussion in Sec. IV) and therefore roughly (up to a smooth envelope) reproduces resistance oscillations.

We will show that, like in other structures [63,74,75], the amplitude of the magneto-oscillations in studied devices is greatly enhanced as compared with the ratchet effect at zero magnetic field. We experimentally demonstrate that the photocurrent oscillates in phase with the longitudinal resistance and, therefore, almost follows the SdH oscillations multiplied by a smooth envelope. This envelope encodes information about cyclotron and plasmonic resonances. The most important experimental result is the demonstration of the beatings of the ratchet current oscillations. We interpret these beatings assuming that they come from SO splitting of the conduction band. The value of SO splitting extracted from the comparison of the experiment and theory is in good agreement with independent measurements of SO band splitting [54–62].

An important comment should be made about the role of the ee interaction. The effect of the interaction is twofold. First of all, sufficiently fast ee collisions drive the system into the hydrodynamic regime. We assume that this is the case for our system and use the hydrodynamic approach. Second, ee interaction leads to plasmonic oscillations, so that a new frequency scale, the plasma frequency $\omega_p(q)$ appears in the problem, where q is the inverse characteristic of the spatial scale in

the system. For a device with a short length, for example, for a single FET, q is proportional to the inverse length of the device. For periodic grating gate structures, $q = 2\pi/L$, where L is the period of the structure. At zero magnetic field, the DC response is dramatically enhanced in the vicinity of plasmonic resonance $\omega = \omega(q)$, both for a single FET with asymmetric boundary conditions [81] and for periodic asymmetric grating gate structures [53]. Also, the response essentially depends on the polarization of the radiation.

Here, we calculate analytically the DC response in the quantizing magnetic field within the hydrodynamic approximation for arbitrary polarization of the radiation and analyzed plasmonic effects. One of our main predictions is that, for linearly polarized radiation, the dependence of the ratchet current on the direction of the polarization appears only due to the plasmonic effects. We use the derived expression to prove that, for specific parameters of our structures, the plasmonic effects are negligible, and as a consequence, the DC response does not depend on the polarization direction. The latter issue is very important for us since the direction of linear polarization used in our experiment was not well controlled. We also argue how to modify the structures to observe plasmonic resonances.

III. EXPERIMENT

A. Experimental details

We chose the AlGaIn/GaN heterostructure system for the experimental study of the effect of SO splitting on the magnetic ratchet effect. Important unique properties of the GaN system are the ability to form high-density, high-mobility 2DEG on the AlGaIn/GaN interface and large Rashba spin splitting of the conduction band [54–62]. Density of 2DEG and the band spin splitting in this system are about an order of magnitude higher than that in the AlGaAs/GaAs system. High carrier density is an important factor because, as will be shown later, the amplitude of the photoresponse in the regime of the SdH oscillations is proportional to the square of electron density.

AlGaIn/GaN heterostructures were grown by the metalorganic vapor phase epitaxy (MOVPE) method in the closed coupled showerhead 3×2 inch Aixtron reactor (Aixtron, Herzogenrath, Germany). The epi-structure consisted of a 25 nm $\text{Al}_{0.25}\text{Ga}_{0.75}\text{N}$ barrier layer, a 1.5 nm $\text{Al}_{0.66}\text{Ga}_{0.37}\text{N}$ spacer, 0.9 μm unintentionally doped GaN layers, and a 2 μm high-resistive GaN:C buffer, see Fig. 1(a). Growth of all mentioned epilayers was done on the bulk semi-insulating GaN substrates, grown by the ammonothermal method [82]. In this method, the high resistivity of substrates (typically no less than $10^9 \Omega\text{cm}$) was obtained by compensation of residual oxygen, incorporated during ammonothermal growth, by Mg shallow acceptors.

The structure lithography processing was performed using a commercial 405 nm laser writer system (Microtech, Palermo, Italy). Devices were isolated from each other by shallow 150 nm mesas etched by inductively coupled plasma reactive ion etching (Oxford Instruments, Bristol, UK). To form the drain and source ohmic contacts, Ti/Al/Ni/Au (150/1000/400/500 Å) stacks were deposited on the MOVPE

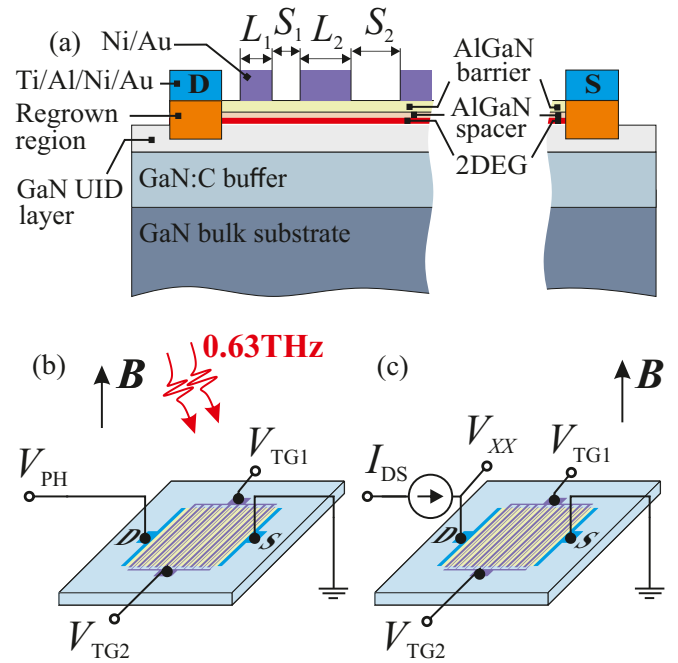


FIG. 1. (a) Cross-sectional view of the sample heterostructure with dual-grating gate (DGG). (b) Scheme of the photoresponse measurements. (c) Scheme of the magnetoresistance measurements $R_{xx} = V_{xx}/I_{DS}$.

regrown heavily doped subcontact regions (for detailed information on the regrowth technique, see Ref. [83]). Source and drain contacts, see Fig. 1, were annealed at 780 °C in a nitrogen atmosphere for 60 s. This procedure yielded reproducible ohmic contacts with resistances in the range of 0.1–0.3 Ωmm . The finishing fabrication step was the deposition of Ni/Au (100/300 Å) to form the DGG superlattice on the top of the AlGaIn/GaN mesas. A schematic view and Nomarski contrast microscope photos of fabricated devices are shown in Figs. 1(a) and 2, respectively. The unit cell of the DGG superlattice consisted of two gates of different lengths ($L_1 = 1.85 \mu\text{m}$ and $L_2 = 3.7 \mu\text{m}$) with different spacings between them ($S_1 = 2.8 \mu\text{m}$ and $S_2 = 5.6 \mu\text{m}$). By that, together with application of different potentials to subgates, we obtained a large lateral asymmetry parameter Ξ [see Eq. (3) below] needed for the enhanced ratchet effect ($j \propto \Xi$) [30]. The cell was repeated 35 times, resulting in a superlattice with a total length of 500 μm . While the ratchet effect requires lateral structure asymmetry, we note that, for real structures, even if they are prepared as “nominally symmetric,” they always have some imperfections, resulting in the irregular lateral asymmetry. In our experiment, we additionally performed measurements on nominally symmetric structures and compared photoresponses in such structures (see the inset in Fig. 3) and strongly asymmetric structures [see Fig. 2(b)]. The results are compared in Fig. 3. As seen, the response in the asymmetric structure is several times larger. The nominally symmetrical structure still provided some signal due to the fabrication imperfections.

All narrow gates in the asymmetric DGG structure were connected forming the multifinger top gate electrode TG1, see Figs. 1(b), 1(c) and 2(a). Similarly connected wide gates

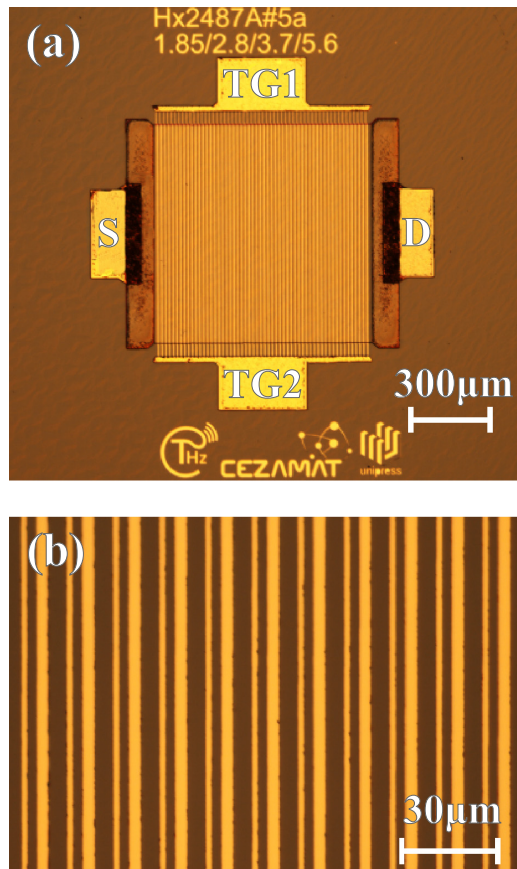


FIG. 2. (a) Nomarski contrast microscope photos of investigated asymmetric dual-grating gate (DGG), where TG1 and TG2 are two multifinger top gate electrodes, and S and D are source and drain electrodes, respectively. (b) Magnified active region of asymmetric DGG.

formed the gate electrode TG2. Independent bias voltages (V_{TG1} , V_{TG2}) could be applied to wide and narrow gates. The width of the whole structure was 0.5 mm, yielding the total active area $A = 0.25 \text{ mm}^2$. The total gate area was $\sim 0.1 \text{ mm}^2$. This is a large area, which is ~ 4 orders of magnitude bigger than that for the “standard” transistor with gate length and width of 0.1 and $100 \mu\text{m}$, respectively. This made the fabrication of the described DGG transistor with a reasonably small gate leakage current very challenging. Figure 4 shows two examples of the transfer current voltage characteristics of the studied devices. Current in the subthreshold region is determined by the gate leakage current (shown as a red dashed line for one of the devices). As seen, the gate leakage current is rather small, significantly smaller than the drain current even at a very low drain voltage of $V_{DS} = 1 \text{ mV}$. Even for those devices with relatively high gate leakage current (#5a in Fig. 4), the drain current and, therefore, electron concentration can be changed several times by the gate voltage.

The experimental setup is shown in Figs. 1(b) and 1(c). As a radiation source, a frequency multiplier from Virginia Diodes Inc. (Charlottesville, Virginia, US) with a radiation frequency of $f = 630 \text{ GHz}$ was used to study the ratchet effect. The radiation was guided onto the sample through a steel waveguide and was modulated at a frequency of $\sim 173 \text{ Hz}$. An external magnetic field up to 12 T was applied normally

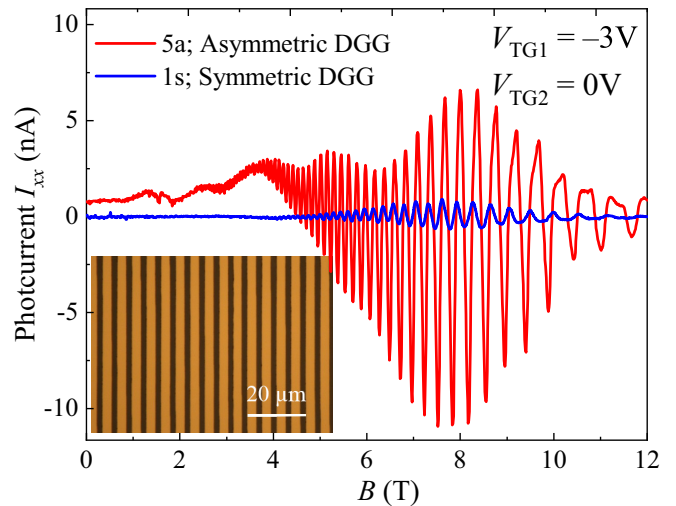


FIG. 3. Ratchet current in nominally symmetric (see inset; gate length $L = 3 \mu\text{m}$ and ungated length $S = 2 \mu\text{m}$) and asymmetric [see Fig. 2(b)] structures at gate voltages $V_{TG1} = -3 \text{ V}$ and $V_{TG2} = 0$.

to the 2DEG plane, as shown in Fig. 1(b). The photoresponse V_{PH} , was measured in a cryostat at the temperature of 4.2 K in the open circuit configuration using the standard lock-in technique. Magnetoresistance was measured by applying a small $< 1 \mu\text{A}$ current to the drain [see Fig. 1(c)].

B. Experimental results

First, we describe the results of the magnetotransport measurements, which are summarized in Figs. 5 and 6. The overall shape of the DGG structures was close to the square. Therefore, in the magnetic field perpendicular to the drain-to-source plane, investigated structures exhibited geometrical magnetoresistance [84]. The full geometrical magnetoresistance is observed either in the disk Corbino geometry or in the samples

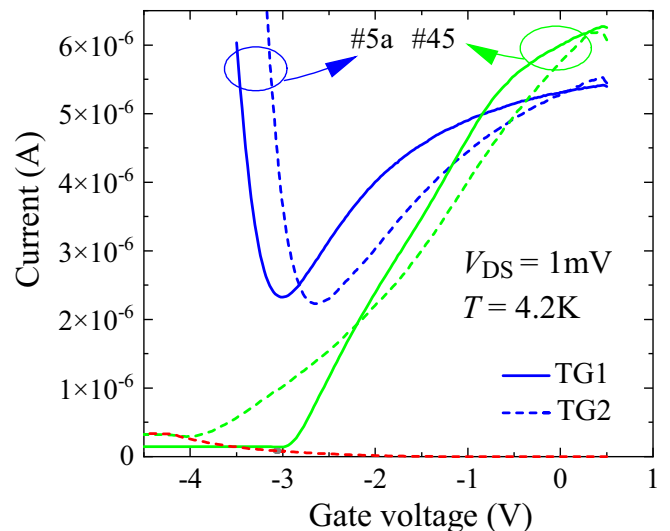


FIG. 4. Transfer current voltage characteristics for two representative devices. Red dashed line shows the gate leakage current for one of the devices.

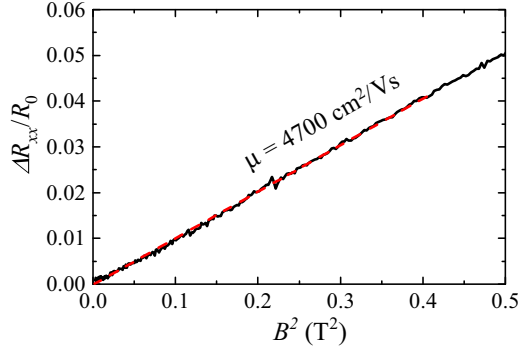


FIG. 5. Magnetoresistance as a function of B^2 in weak magnetic fields for a representative sample. Dashed line shows the linear fit.

with $W \gg L$. For the arbitrarily shaped rectangular samples, the geometrical magneto-resistance can be approximated as [85]

$$\frac{\Delta R}{R_0} \cong (\mu B)^2 \left(1 - 0.54 \frac{L}{W}\right). \quad (1)$$

This allowed us to extract the electron mobility. Figure 5 shows the experimental dependence of the magnetoresistance obtained for a weak magnetic field for one of the samples as a function of B^2 . The estimate yields $\mu = 4700 \text{ cm}^2/\text{Vs}$.

The concentration in the channel can be extracted from the magnetoresistance in the higher magnetic fields. Figure 6 shows the resistance SdH oscillations as a function of the inverse magnetic field $1/B$ measured for both top gate voltages equal to zero. The concentration is given by

$$N = \frac{2e}{h\Delta(1/B)} = \frac{2e\nu}{h}, \quad (2)$$

where $\Delta(1/B)$ and ν are the period and frequency of SdH oscillations, respectively.

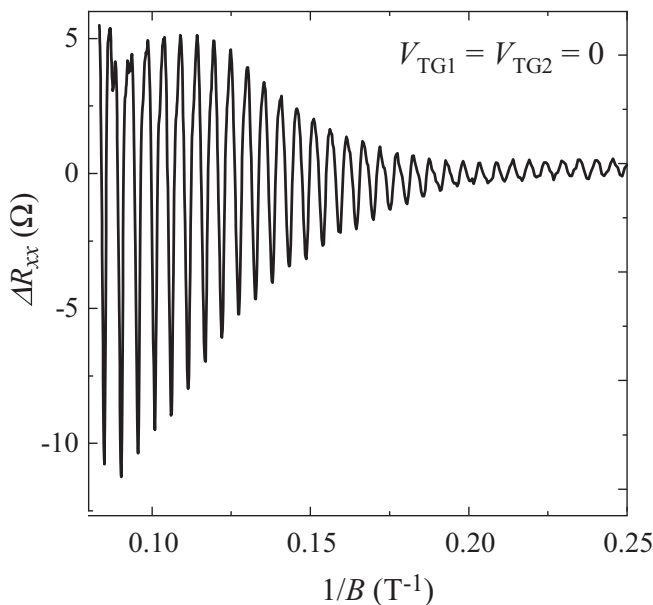


FIG. 6. Resistance as a function of the inverse magnetic field $1/B$ at $V_{TG1} = V_{TG2} = 0 \text{ V}$.

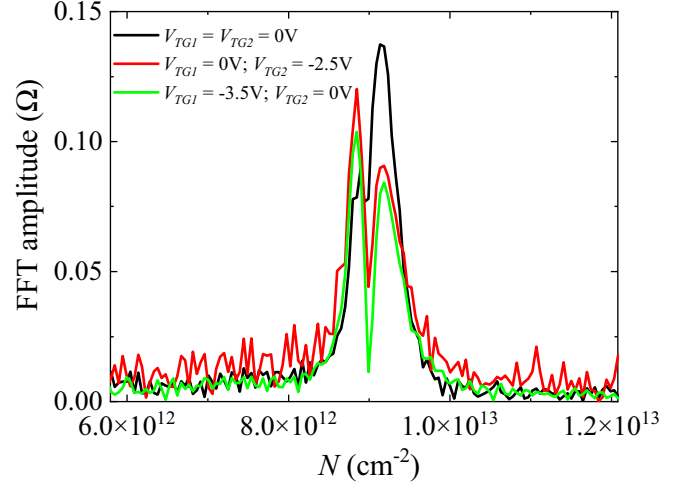


FIG. 7. Fourier transform of Shubnikov–de Haas (SdH) resistivity oscillations at different gate voltages.

Figure 7 shows the result of the Fourier transform of the resistance magneto-oscillations with the frequency taken in the units of the electron concentration $N = 2e\nu/h$. Two peaks correspond to the concentrations $N = 9.1 \times 10^{12} \text{ cm}^{-2}$ and $N = 8.9 \times 10^{12} \text{ cm}^{-2}$.

Let us discuss this result in more detail. In fact, there are two types of regions in the sample—gated and ungated—and in each region, the electrons have their own concentration corresponding to different Fermi energies $E_{F,g}$ and $E_{F,u}$. Due to the SO splitting, the Fourier spectrum of SdH oscillation could show four harmonics, corresponding to $E_{F,g} \pm \Delta_g$ and $E_{F,u} \pm \Delta_u$, where $\Delta_{g,u}$ are band splittings in the gated and ungated regions. However, in experiment, we see only two harmonics. As the concentration under the gate is decreased by variation of the gate potential, the resistance shows a strong (at least several times) increase. At the same time, the experimentally observed spectrum of SdH resistivity oscillations obtained by Fourier transform of these oscillations does not change essentially. Positions of two observed peaks are the same as for zero gate potentials, when we have homogeneous concentration (see Fig. 7). This means that the main contribution to the oscillations comes from the ungated region, in which the concentration is practically independent of the gate potential. We conclude that, in our experiment, we observe SO splitting of ungated electrons, corresponding to energies $E_{F,u} \pm \Delta_u$ (in what follows, we skip subscript u).

Irradiating the unbiased structures, we detected a photo-signal caused by the generation of the ratchet photocurrent. Figure 8(a) shows the photoresponse measured for asymmetric gate voltages applied: $V_{TG1} = -3 \text{ V}$, $V_{TG2} = 0$. The photoresponse current was calculated as $I_{xx} = V_{PH}/R_{xx}$. In low and zero magnetic fields, the response is positive and weakly depends on the magnetic field. When gate voltages were changed to $V_{TG1} = 0 \text{ V}$, $V_{TG2} = -3 \text{ V}$, the magnitude of the response was approximately the same but of the negative sign. The change of the signal sign upon inversion of the lateral asymmetry is a clear indication that the observed photocurrent is caused by the ratchet effect, for review, see Ref. [30]. Indeed, the direction of the current is controlled by

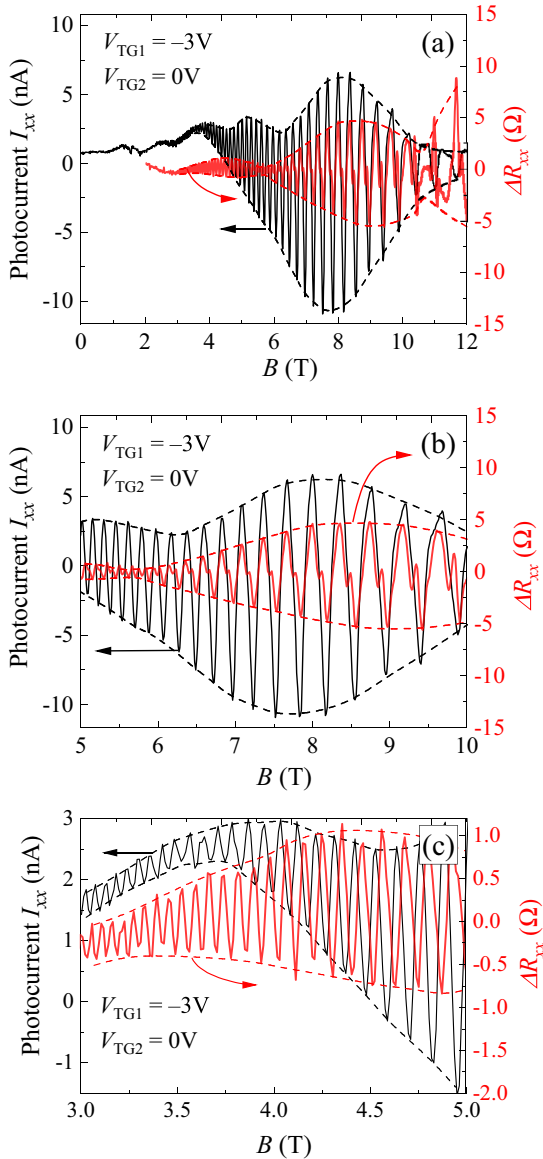


FIG. 8. (a) Photocurrent and resistance Shubnikov–de Haas (SdH) oscillations as a function of magnetic field from $B = 0$ to $B = 12$ T. Close look at (b) intermediate and (c) low magnetic field regions. The dashed lines are guides for the eye.

the lateral asymmetry parameter [30]

$$\Xi = \overline{|\mathbf{E}(x)|^2 \frac{dV(x)}{dx}}, \quad (3)$$

where $V(x)$ and $E(x)$ are spatially modulated by grating gate static potential and electric field amplitude (overline shows average over the modulation period). Exchange of the gate voltages applied to the TG1 and TG2 results in the change of sign of dV/dx and, consequently, in the sign inversion of the ratchet current.

The increase of the magnetic field results in the sign-alternating oscillations with the amplitude by orders of magnitude larger than the signal obtained for zero magnetic field. Moreover, the envelope of the oscillations exhibits beatings as a function of the magnetic field. Comparison of the

observed oscillations with the SdH magnetoresistance oscillations demonstrates that, at high magnetic fields, both photocurrent and resistivity oscillations have the same period and phase. Importantly, SdH effect also show similar beatings of the envelope function. To facilitate the comparison of the phase of the oscillations, we zoom in to the data of panel (a) for the range of fields $B = 6$ – 10 T in panel (b) of Fig. 8.

The overall behavior of the observed current corresponds to that of the magnetoratchet current most recently detected in CdTe-based quantum wells [74,75] and graphene [63]. Importantly, the oscillations of the magnetoratchet current are in phase with the SdH oscillations. As we discussed in Sec. II, this differs from the photocurrent magneto-oscillations detected in single transistors [80] described by the theoretical model of Lifshits and Dyakonov [78]. We will show in Sec. IV that the theory of the magnetoratchet effect predicts the photoresponse to be proportional to the second derivative of the magnetoresistance and, hence, describes well the experimental findings. Importantly, the theory demonstrates that, in agreement with the experiment (see Fig. 8), the photoresponse in the regime of SdH oscillations is significantly enhanced as compared with that at zero magnetic field. This point is very important in view of possible applications and deserves a special comment. The theoretical limit for the response of the detectors based on the direct rectification is defined by the device built-in nonlinearity. In Schottky diodes and FETs, the maximum current responsivity is $\sim e/2\eta kT$, where e is the elemental charge, and η is the ideality factor of the Schottky barrier or the subthreshold slope of the FET transfer characteristic [86,87]. The responsivity of the real device is usually orders of magnitude smaller due to the parasitic elements and not perfect coupling. However, increasing of the theoretical limit still should be beneficial for the increasing of the responsivity in real devices. Reducing temperature indeed leads to the responsivity increase but only to a certain limit. As shown in Ref. [88], the temperature decrease below 30 K does not lead to the increase of responsivity. This low-temperature saturation is caused by the increase of factor η with temperature decrease, which is known in Schottky diodes and FETs [89]. In the magnetic field under the regime of SdH oscillations, resistance of FET very sharply depends on the gate voltage and provides the opportunity to go beyond the $e/2\eta kT$ limit. We can speculate that about an order of magnitude increase of the responsivity in an external magnetic field in Fig. 8 demonstrates the increase of the physical responsivity beyond the fundamental limit.

Importantly, the theory presented below also describes well the observed oscillations of the envelope amplitude of the photoresponse, which are clearly seen in Fig. 8. It shows that the oscillations are due to the spin-splitting of the conduction band and can be used to extract this important parameter.

IV. THEORY

Above, we experimentally demonstrated that the photocurrent oscillates, and the shape of these oscillations almost follow the SdH resistivity oscillations. In this section, we consider gated 2DEG and demonstrate that the results can be well explained within the hydrodynamic approach.

The effect which we discuss here is present for the system with an arbitrary energy spectrum. However, calculations are dramatically simplified for the parabolic spectrum, so that we limit calculations to this case only. We assume that electron density in the structure is periodically modulated by built-in static potential and study optical DC response to linearly polarized electromagnetic radiation, which is also spatially modulated with the phase shift φ with respect to modulation of the static potential.

A variation of individual gate voltages of the DGG lateral structure allows one to change controllably the sign of $V(x)$ and, consequently, the direction of the ratchet current. Furthermore, the phase of the oscillations of the magnetic ratchet current is sensitive to the orientation of the radiation electric field vector with respect to the DGG structure as well as to the radiation helicity. In the latter case, switching from right- to left-circular polarization results in the phase shift by π , i.e., at constant magnetic field, the helicity-dependent contribution to the current changes the sign.

We note that the observed beatings cannot be related to the microwave-induced resistance oscillations as well as to the phonon- and Hall-induced resistance oscillations [90–95], because they are observed in uniform 2D electron systems and ordinarily require two orders of magnitude larger mobility. Apart from that, all these oscillations are insensitive to the position of the chemical potential with respect to Landau levels and, therefore, should be almost insensitive to the weak periodic potential modulation, in sharp contrast to the SdH oscillations considered in this paper.

Below, we consider magnetic field-induced modification of the zero B -field electronic ratchet effect and develop a theory which is able to describe giant enhancement of the ratchet effect in the magnetic field and beating of the photocurrent oscillations. We also analyze our results theoretically for different relations between ω and $\omega_p(q)$ to find a signature of the plasmonic effects.

A. Model

We model the electric field of the radiation $\mathbf{E}(x, t) = E(x)e^{-i\omega t} + \text{c.c.}$, and the static potential V as follows:

$$E_x(x, t) = [1 + h \cos(qx + \varphi)]E_0 \cos \alpha \cos \omega t, \quad (4)$$

$$E_y(x, t) = [1 + h \cos(qx + \varphi)]E_0 \sin \alpha \cos(\omega t + \theta), \quad (5)$$

$$V(x) = V_0 \cos qx, \quad (6)$$

where $h \ll 1$ is the modulation depth, φ is the phase, which determines the asymmetry of the modulation, and α and θ are constant phases describing the polarization of the radiation. These phases are connected with the standard Stokes parameters (normalized by E_0^2) as follows:

$$\begin{aligned} P_0 &= 1, & P_{L1} &= \sin(2\alpha) \cos \theta, \\ P_{L2} &= \cos(2\alpha), & P_C &= \sin(2\alpha) \sin \theta. \end{aligned} \quad (7)$$

Within this model, the asymmetry parameter [see Eq. (3)] becomes

$$\Xi = \frac{E_0^2 h V_0 q \sin \varphi}{4}. \quad (8)$$

As seen, Ξ is proportional to the sine of the spatial phase shift φ .

Hydrodynamic equations for concentration and velocity are given by

$$\frac{\partial n}{\partial t} + \text{div}[(1+n)\mathbf{v}] = 0, \quad (9)$$

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v}\nabla)\mathbf{v} + \gamma(n)\mathbf{v} + \boldsymbol{\omega}_c \times \mathbf{v} + s^2 \nabla n = \mathbf{a}. \quad (10)$$

Here,

$$n = \frac{N - N_0}{N_0}, \quad (11)$$

where $N = N(x, t)$ is the concentration in the channel and N_0 its equilibrium value,

$$\mathbf{a} = -\frac{e\mathbf{E}}{m} + \frac{e}{m}\nabla V, \quad \mathbf{E} = \begin{bmatrix} E_x \\ E_y \end{bmatrix}, \quad (12)$$

and $\boldsymbol{\omega}_c = e\mathbf{B}/m_{\text{eff}}c$ is the cyclotron frequency in the external magnetic field \mathbf{B} , s is the plasma waves velocity, and $\gamma(n) = 1/\tau_{\text{tr}}(n)$ is momentum relaxation rate. The nonlinearity is encoded in hydrodynamic terms $\partial(n\mathbf{v})/\partial x$, $(\mathbf{v}\nabla)\mathbf{v}$ as well as independence of transport scattering rate on the concentration. Specifically, we use the approach suggested in Ref. [78]. We assume that $\gamma(n)$ is controlled by the local value of the electron concentration n , which in turn, is determined by the local value of the Fermi energy $n(\mathbf{r}) = [E_F(\mathbf{r}) - E_F^0]/E_F^0$ (here, we considered that the 2D density of states is energy independent). Due to the SdH oscillations, scattering rate is an oscillating function of E_F and, consequently, oscillates with n . In the absence of SO coupling, $\gamma(x, t) = \gamma[n(x, t)]$ is given by [78]

$$\gamma(x, t) = \gamma \left\{ 1 - \delta \cos \left[\frac{2\pi E_F(x, t)}{\hbar\omega_c} \right] \right\}, \quad (13)$$

where

$$\delta = \frac{4\chi}{\sinh \chi} \exp \left(-\frac{\pi}{\omega_c \tau_q} \right), \quad (14)$$

is the amplitude of the SdH oscillations,

$$\chi = \chi(\omega_c) = \frac{2\pi^2 T}{\hbar\omega_c},$$

where T is the temperature in the energy units, $E_F(x, t) = E_F[1 + n(x, t)]$ is the local Fermi energy, which is related to concentration in the channel as $N(x, t) = \nu E_F(x, t)$ (here, ν is the density of states), and τ_q is quantum scattering time, which can be strongly renormalized by ee collisions in the hydrodynamic regime. We assume that $2\pi^2 T + \pi \hbar/\tau_q \gg \hbar\omega_c$. Then

$$\delta \ll 1, \quad (15)$$

and the second term in the curly brackets in Eq. (13) is very small. Hence, $\gamma(n)$ is very close to the value of transport scattering rate γ at zero magnetic field.

Equation (13) can be generalized for the case of nonzero SO coupling by using results of Refs. [75,96,97]

$$\gamma(n) = \gamma \left\{ 1 - \frac{4\chi}{\sinh \chi} \exp\left(-\frac{\pi}{\omega_c \tau_q}\right) \times \cos\left[\frac{2\pi E_F(1+n)}{\hbar\omega_c}\right] \cos\left(\frac{2\pi\Delta}{\hbar\omega_c}\right) \right\}. \quad (16)$$

Here, we assumed that there is linear-in-momentum SO splitting of the spectrum $E(k) = \hbar^2 k^2/2m \pm \Delta$, where

$$\Delta = \alpha_{SO} k_F, \quad (17)$$

is characterized by coupling constant α_{SO} . Experimentally measured values of α_{SO} lay between 4–10 meVÅ [54–62]. For such values of α_{SO} and typical values of the concentration, one can assume $\Delta \ll E_F$ and neglect dependence of k_F on n . Equation (16) was derived under the assumption that the quantum scattering time τ_q is the same in two SO split subbands. This assumption is correct only for the model of short-range scattering potential where both transport and quantum scattering rates are momentum independent. For any finite-range potential, the quantum scattering times in two subbands differ because of small difference of the Fermi wave vectors $k_1 = \sqrt{2m(E_F + \Delta)/\hbar}$ and $k_2 = \sqrt{2m(E_F - \Delta)/\hbar}$. Denoting these times as τ_1 and τ_2 , we get instead of Eq. (16)

$$\frac{\gamma(n)}{\gamma} = 1 - \frac{2\chi}{\sinh \chi} \sum_{i=1,2} \exp\left(-\frac{\pi}{\omega_c \tau_i}\right) \cos\left[\frac{2\pi E_i(n)}{\hbar\omega_c}\right], \quad (18)$$

where

$$E_1(n) = E_F(1+n) + \Delta, \quad E_2(n) = E_F(1+n) - \Delta. \quad (19)$$

Detailed microscopical calculation of $\tau_{1,2}$ for a specific model of the scattering potential is out of the scope of this paper. Here, we use $\tau_{1,2}$ as fitting parameters.

Let us now expand $\gamma(n)$ near the Fermi level:

$$\gamma(n) = \gamma(0) + \gamma'(0)n + \gamma''(0)\frac{n^2}{2}, \quad (20)$$

where γ' and γ'' are, respectively, first and second derivatives with respect to n taken at the Fermi level. Since oscillations are very fast, we assume

$$\frac{\gamma'}{\gamma} \propto \frac{\gamma''}{\gamma'} \propto \frac{E_F}{\hbar\omega_c} \gg 1. \quad (21)$$

Due to these inequalities, oscillating contribution to the ratchet current can be very large and substantially exceed zero-field value [63].

Here, we focus on SdH oscillations of the ratchet current, so that we only keep oscillations related to dependence of γ on n and, moreover, skip in Eq. (20) the term proportional to γ' .

We use the same method of calculation as the one developed in Ref. [53]. Specifically, like the impurity-dominated regime [30], we use the perturbative expansion of n and \mathbf{v} and DC current,

$$\mathbf{J}_{dc} = -eN_0 \langle (1+n)\mathbf{v} \rangle_{t,x}, \quad (22)$$

over E_0 and V . Nonzero contribution $\propto E_0^2 V_0$ arises in the order (2,1) [see Eq. (3)].

B. Calculations and results

Let us formulate the key steps of calculations. Due to the large parameter $E_F/\hbar\omega_c \gg 1$, the main contribution to the rectified ratchet current comes from the nonlinear term $\gamma'' \mathbf{v} n^2/2$ in the right-hand side of Eq. (10) [see also Eq. (20)]. We, therefore, neglect all other nonlinear terms in the hydrodynamic equations. Calculating n and \mathbf{v} in linear (with respect to E_0 and V) approximation, substituting the result into a nonlinear term and averaging over time and coordinate, we get $\gamma''(0) \langle \mathbf{v} n^2 \rangle_{x,t}/2 \neq 0$. Next, one can find rectified current \mathbf{J}_{dc} by averaging of Eq. (10) over t and x . This procedure is quite standard, so that we delegate it to the Supplemental Material [98] (similar calculations were performed in Ref. [53] for zero magnetic field). The result reads

$$\frac{\mathbf{J}_{dc}}{J_0} = \frac{\gamma''(0)}{\gamma} \mathbf{R}. \quad (23)$$

Here,

$$J_0 = -\left(\frac{eE_0}{2m}\right)^2 \left(\frac{eV_0 q}{2ms^2}\right) \frac{eN_0 \hbar \sin \varphi}{\gamma^3}, \quad (24)$$

is the frequency and magnetic field-independent parameter with the dimension of the current (physically, J_0 gives the typical value of current for the case when all frequencies are of the same order $\omega \sim \omega_c \sim qs \sim \gamma \sim 1/\tau_{1,2}$), and the dimensionless factor γ''/γ accounts for SdH oscillations and the dimensionless vector

$$\mathbf{R} = \frac{\gamma^4 (P_0 \mathbf{a}_0 + P_{L1} \mathbf{a}_{L1} + P_{L2} \mathbf{a}_{L2} + P_C \mathbf{a}_C)}{|\omega_c^2 - (\omega - i\gamma)^2|^2 (\gamma^2 + \omega^2) (\gamma^2 + \omega_c^2) |D_{\omega q}|^2}, \quad (25)$$

depends on the radiation polarization encoded in the vectors

$$\mathbf{a}_i = \begin{bmatrix} a_{ix} \\ a_{iy} \end{bmatrix},$$

($i = 0, L1, L2, C$) and also contains information about cyclotron and magnetoplasmon resonances which occur for $\omega = \omega_c$ and $\omega = \sqrt{\omega_c^2 + s^2 q^2}$, respectively. The latter resonance appears due to the factor $D_{\omega q}$ in the denominator of Eq. (25). Analytical expressions for \mathbf{a}_i and $D_{\omega q}$ are quite cumbersome and presented in the Supplemental Material [98] [see Eqs. (47)–(51)].

The second derivative of the scattering rate with respect to n is calculated by using Eq. (18):

$$g(\omega_c) = \frac{\gamma''(0)}{\gamma} = \frac{2\chi(\omega_c)}{\sinh[\chi(\omega_c)]} \left(\frac{2\pi E_F}{\hbar\omega_c}\right)^2 \times \sum_{i=1,2} \exp\left(\frac{-\pi}{\omega_c \tau_i}\right) \cos\left[\frac{2\pi E_i(0)}{\hbar\omega_c}\right]. \quad (26)$$

Here, $E_1(0) = E_F + \Delta$, $E_2(0) = E_F - \Delta$ [see Eq. (19)]. The function $g(\omega_c)$ rapidly oscillates due to the factors $\cos[2\pi E_i(0)/\hbar\omega_c]$. For $\omega_c \rightarrow 0$, this function goes to zero due to the Dingle factors $\exp(-\pi/\omega_c \tau_i)$, so that discussed mechanism has nothing to do with the zero-field ratchet effect.

The smooth envelope of the function $g(\omega_c)$ reads

$$\begin{aligned} \tilde{g}(\omega_c) &= \frac{2\chi(\omega_c)}{\sinh[\chi(\omega_c)]} \left(\frac{2\pi E_F}{\hbar\omega_c} \right)^2 \\ &\times \left| \exp\left(\frac{-\pi}{\omega_c\tau_1} + \frac{2\pi i\Delta}{\hbar\omega_c} \right) + \exp\left(\frac{-\pi}{\omega_c\tau_2} - \frac{2\pi i\Delta}{\hbar\omega_c} \right) \right|. \end{aligned} \quad (27)$$

The function $g(\omega_c)$ shows rapid SdH oscillations with the beats due to the SO coupling. As seen from the behavior of the envelope function $\tilde{g}(\omega_c)$, the beats are most pronounced for $\tau_1 = \tau_2$, when $\tilde{g}(\omega_c)$ is proportional to $\cos(2\pi\Delta/\hbar\omega_c)$ and therefore vanishes at the values of ω_c^n , obeying $2\pi\Delta/\omega_c^n = \pi/2 + \pi n$. For $\tau_1 \neq \tau_2$, the envelope function is nonzero at these points $\tilde{g}(\omega_c^n) \neq 0$, and beats are less pronounced.

Now we are ready to explain why the response in the SdH oscillation regime is much larger than at zero magnetic field. The enhancement of the response as compared with the case $B = 0$ is due to the factor

$$\delta \left(\frac{2\pi E_F}{\hbar\omega_c} \right)^2 \gg 1. \quad (28)$$

One can check that, for experimental values of parameters, the inequalities in Eqs. (15) and (28) are satisfied simultaneously in a wide interval of magnetic fields $1 < B < 7$ T. It is also important that, due to the coefficient E_F^2 in the $g(\omega_c)$, the response increases with the concentration in contrast to a conventional transistor operating at $B = 0$, where the response is inversely proportional to the concentration at high concentration [81] and saturates at low concentration when a transistor is driven below the threshold [99]. This means that the use of the AlGaIn/GaN system for detectors operating in the SdH oscillation regime is very advantageous because of the extremely high concentration of 2DEG.

Let us discuss the polarization dependence of the response. Importantly, vectors \mathbf{a}_i , responsible for polarization dependence, contain q -independent terms and terms proportional to $\omega_q^2 = s^2 q^2$. The latter describe plasmonic effects. As seen from Eqs. (47)–(49), for small q (or/and small s), vectors \mathbf{a}_{L1} and \mathbf{a}_{L2} are small, $\propto q^2$. In other words, for our case of linearly polarized radiation with polarization directed by angle α , the dependence of the rectified current on α appears only due to the plasmonic effects. For experimental values of the parameters, the value of plasmonic frequency sq was sufficiently small $\sim 0.7 \times 10^{12} \text{ s}^{-1}$, which is much smaller than the radiation frequency (for $f = 0.6$ THz, we get $\omega = 2\pi f \approx 3.8 \times 10^{12} \text{ s}^{-1}$). As follows from this estimate, the plasmonic effects are small and can be neglected.

Then the response does not actually depend on polarization angle α . This justifies our experimental approach, where α is not well controlled. Within this approximation, one can put $q \rightarrow 0$ in Eqs. (47)–(49). Then the analytical expression for current simplifies. In the absence of the circular component of polarization ($P_C = 0$), we get

$$\begin{bmatrix} J_{\text{dc}}^x \\ J_{\text{dc}}^y \end{bmatrix} = \frac{2J_0 g(\omega_c) \gamma^4 \omega_c}{(\gamma^2 + \omega_c^2) |(\omega + i\gamma)^2 - \omega_c^2|^2} \begin{bmatrix} -\omega_c \\ \gamma \end{bmatrix}. \quad (29)$$

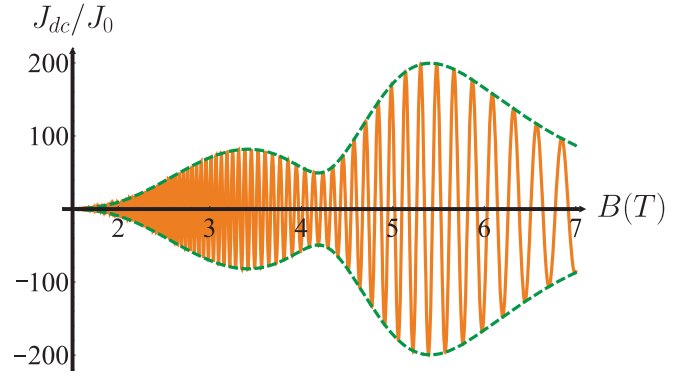


FIG. 9. Theoretically calculated ratchet magneto-oscillations for the following parameters: $\epsilon = 9$, $\alpha_{\text{SO}} = 7.5 \text{ meV\AA}$, $T = 4 \text{ K}$, $m_{\text{eff}} = 0.23m_e$, $d = 2.5 \times 10^{-6} \text{ cm}$, $L = 15 \times 10^{-4} \text{ cm}$, $N_0 = 8 \times 10^{12} \text{ cm}^{-2}$, $\alpha = 0$, $\tau_{\text{tr}} = 1.2 \times 10^{-12} \text{ s}$, $\tau_1 = 1.4 \times 10^{-12} \text{ s}$, $\tau_2 = 10^{-12} \text{ s}$.

This expression simplifies even further in the resonant regime $\omega \approx \omega_c \gg \gamma$:

$$\begin{bmatrix} J_{\text{dc}}^x \\ J_{\text{dc}}^y \end{bmatrix} = \frac{J_0 g(\omega_c) \gamma^4}{2\omega_c^3 [(\omega - \omega_c)^2 + \gamma^2]} \begin{bmatrix} -\omega_c \\ \gamma \end{bmatrix}. \quad (30)$$

This expression shows rapid oscillations, described by the function $g(\omega_c)$, whose envelope represent a sharp CR with the width γ .

Let us now compare theoretical results with experimental observations. In Fig. 9, we plot the x component of the rectified DC (this component was measured in the experiment), calculated with the use of Eq. (23), as a function of the magnetic field. We assumed that the radiation is linearly polarized along the x axis [$\alpha = \theta = 0$, $P_C = P_{L1} = 0$, $P_{L2} = 1$, see Eqs. (4), (5), and (7)] and used experimental values of parameters: $m_{\text{eff}} = 0.23m_e$, $d = 2.5 \times 10^{-6} \text{ cm}$, $L = 15 \times 10^{-4} \text{ cm}$, $\epsilon = 9$, $T = 4 \text{ K}$, $n = 8 \times 10^{12} \text{ cm}^{-2}$, $\tau_{\text{tr}} = \gamma^{-1} = 10^{-12} \text{ s}$, $\omega = 3.8 \times 10^{12} \text{ s}^{-1}$. The best fit was obtained for $\alpha_{\text{SO}} = 7.5 \pm 1.5 \text{ meV\AA}$, in accordance with previous measurements of SO band splitting [54–62]. We used $\tau_{1,2}$ as the fitting parameters, choosing $\tau_{\text{tr}} = 1.2 \tau_1 = 1.4 \tau_2$. We see that exactly this behavior is observed in the experiment (see Fig. 8). Most importantly, we reproduce experimentally observed beats of SdH oscillations using the value of α_{SO} consistent with previous experiments. In Figs. 10–12, we show dependence of the smooth envelope of the current \tilde{J}_x on the magnetic field for different values of $\tau_{1,2}$ and different concentrations. As we explained above, the most pronounced modulation is obtained for $\tau_1 = \tau_2$ (see Fig. 10). Dependence on concentration appears both due to the factor $(E_F/\hbar\omega_c)^2$ in $g(\omega_c)$ and due to the dependence of Δ on k_F .

To get a better agreement with experiment, we note that, in addition to the beats of the experimentally observed oscillation amplitude, there is a slowly varying background which also shows oscillation (see Fig. 8). We do not discuss the physical origin of this background here. To focus on the beats of the SdH photocurrent oscillations, we removed this background from the experimental dependence and obtained the dependence shown in Fig. 13(a). This dependence can be perfectly fitted by theoretical formulas Eqs. (23) and (25) [for values of parameters used in experiment, one can use

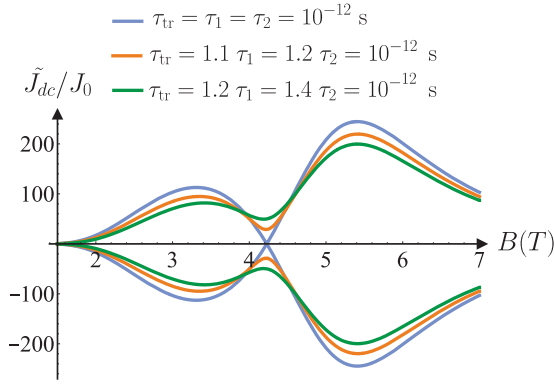


FIG. 10. Envelope of the ratchet current magneto-oscillations for different ratios of quantum times τ_1 and τ_2 (other parameters are the same as in Fig. 9).

simplified Eq. (30)] with the cyclotron mass $m = 0.37m_e$, see Fig. 13(b). The difference between this value and the value $m = 0.23m_e$ for GaN can be explained by tunneling tails of the wave function of the 2D electrons to the AlGaIn barrier with higher mass [100,101].

Evidently, Eq. (29) can be presented as a product of a smooth function describing CR and a rapidly oscillating function, which encodes information about SO splitting. This is illustrated in Fig. 14. At the end of this section, we note that, as follows from our theory, the photocurrent is proportional to the second derivative of the resistivity and, therefore, is in phase with the SdH oscillations. There are some subleading corrections to the photocurrent which are proportional to the first derivative of the resistivity. These corrections can be responsible for a phase shift between the photocurrent and the SdH oscillation seen in Figs. 8(b) and 8(c). The accurate calculation of subleading terms is quite difficult and beyond the scope of this paper.

C. Role of the plasmonic effects

Above, we demonstrated that plasmonic effects can be neglected for our experimental parameters, and as a consequence, the response is insensitive to the direction of

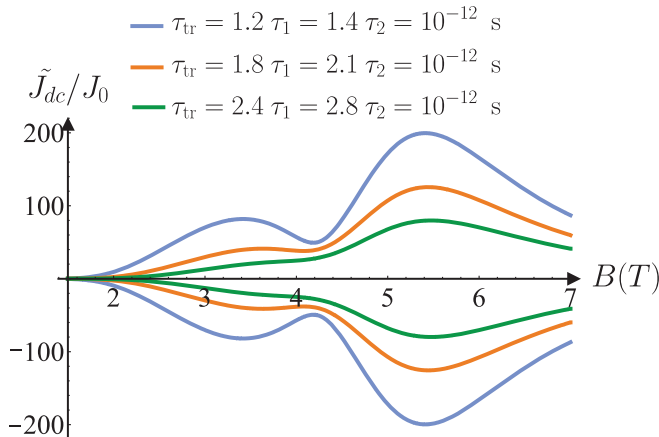


FIG. 11. Envelope of the ratchet current magneto-oscillations for different values of the quantum times τ_1 and τ_2 with the fixed ratio τ_1/τ_2 (other parameters are the same as in Fig. 9).

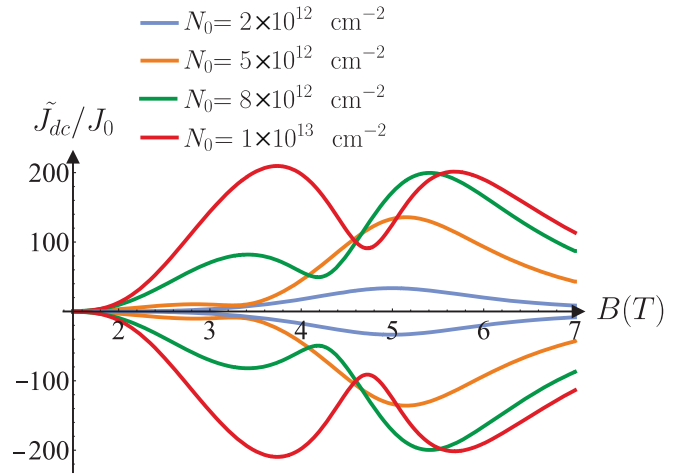


FIG. 12. Envelope of the ratchet current magneto-oscillations for different values of the electron concentration (other parameters are the same as in Fig. 9).

the linear polarization. However, the role of plasmonic effects is not fully understood. The point is that the existing ratchet theory assumes a weak coupling with a diffraction grating. In such a situation, the plasmon wave vector, which determines plasma oscillation frequency, is set by the total lattice period: $q = 2\pi/L$. In the experiment, $L = 13.95 \mu\text{m}$, i.e., is very large, and as a consequence, the plasma frequency corresponding to the full period is small. This frequency does not appear in the experiment, as follows from the theoretical pictures presented above [see Fig. 15(a)]. If we assume that the coupling is not so weak, then the plasmons determined only by the gate region should show up. Then q is determined

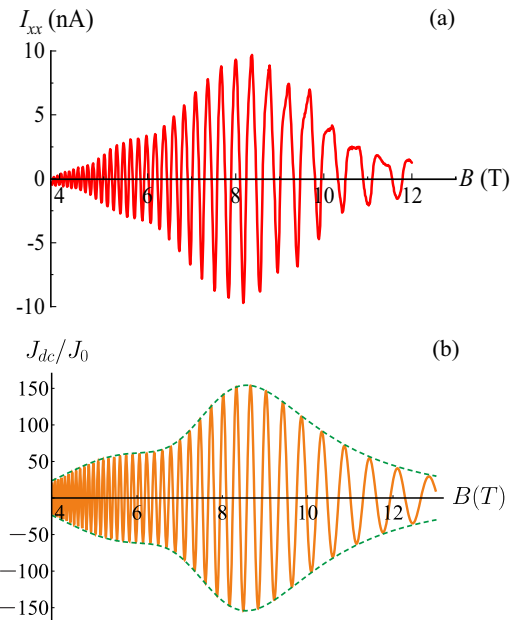


FIG. 13. (a) Experimentally measured photoresponse after removing slow varying background in Fig. 8(a). (b) Theoretical fit for $\alpha_{SO} = 6.8 \times 10^{11}$ eVcm, $n = 9.3 \times 10^{12}$ cm $^{-2}$, $\tau_r = \tau_1 = 1.45 \tau_2$, and $m = 0.37 m_e$.

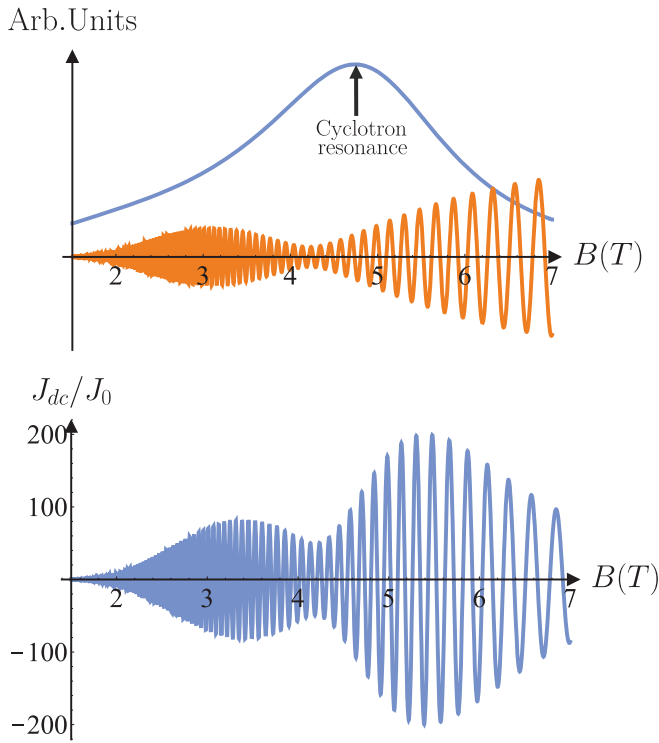


FIG. 14. Direct current (DC) response for small q [shown in bottom panel, described by Eq. (29)] is given by the product of smooth function R_x [blue curve in upper panel], which shows cyclotron resonance (CR) and a rapidly oscillating function $g(\omega_c)$ [orange curve in upper panel, described by Eq. (26)], which contains beats of Shubnikov–de Haas (SdH) oscillations.

only by the gate length $L_g < L$, and it should manifest itself, as it is shown in Fig. 15. Namely, plasma wave effects should lead to the plasmonic splitting of the CR.

Hence, our theory fully considers plasmonic effects including finite plasmonic lifetimes related to the momentum relaxation rate γ . The role of the plasmonic effects can be enhanced either by decreasing the period of the structure, which implies increasing of q , or by decreasing the transport scattering rate γ , as illustrated in Fig. 15, where CR and magnetoplasmon resonance in the smooth function R_x [see Eq. (25)] are shown for different values of L and γ . Although, for large values of γ [Fig. 15(a)], plasmonic effects are fully negligible for small q (large L), and the function R_x shows only CR at $\omega_c = \omega$. With decreasing L , there appears a weak plasmonic resonance at $\omega_c = \sqrt{\omega^2 - s^2 q^2}$. For smaller γ [Figs. 15(b) and 15(c)], both CR and magnetoplasmonic resonance become sharper. For very small γ [Fig. 15(c)], plasmonic resonance appears even for very small q .

V. CONCLUSIONS

To conclude, we presented observation of the magnetic ratchet effect in a GaN-based structure superimposed with a lateral superlattice formed by a DGG structure. We showed that terahertz excitation results in the giant magnetoo-oscillation of the ratchet current in the regime of SdH oscillations. The amplitude of the oscillations is greatly enhanced as compared with the ratchet effect at zero magnetic

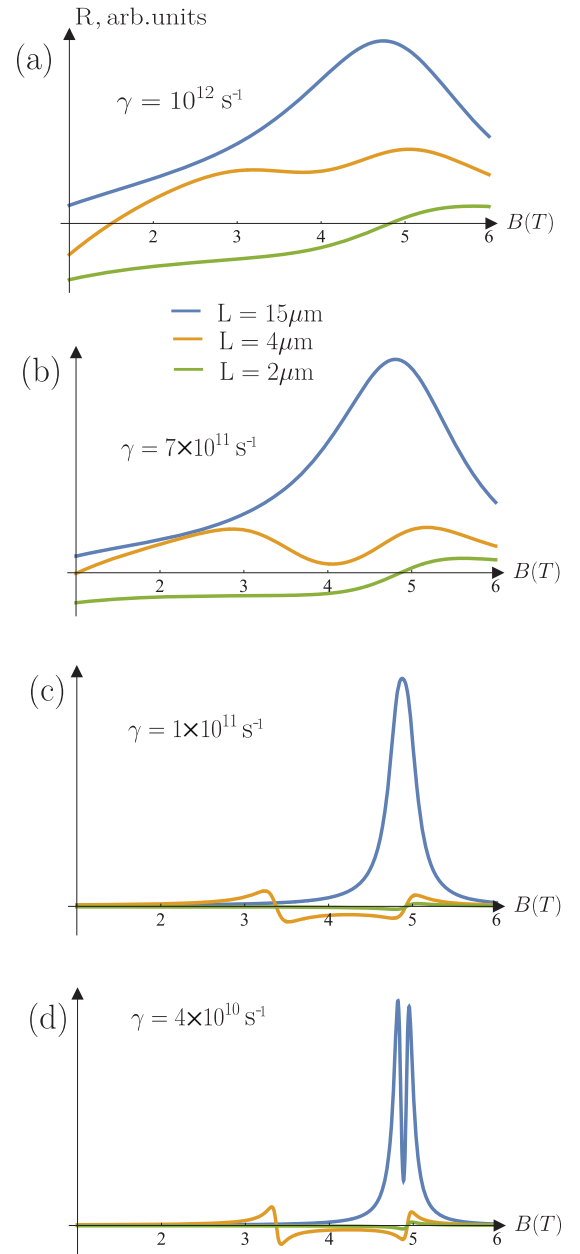


FIG. 15. Cyclotron (CR) and magnetoplasmon resonances in the smooth function R_x [see Eq. (25)] for different values of wave vector $q = 2\pi/L$ (determined by size of the unit cell L) and transport scattering rate γ . (a) For large values of γ plasmonic effects are fully negligible for small q (large L) and the function R_x shows only CR at $\omega_c = \omega$ (this case corresponds to our experimental situation). With decreasing L , there appears a weak plasmonic resonance at the value of ω_c given by $\sqrt{\omega^2 - s^2 q^2}$. (b) and (c) For smaller γ , both cyclotron and magnetoplasmonic resonances become sharper. (d) For very small γ , plasmonic resonance appears even for very small q .

field. We demonstrate that the photocurrent oscillates as the second derivative of the longitudinal resistance and, therefore, almost follows the SdH resistivity oscillations multiplied by a smooth envelope. This envelope encodes information about CR. One of the most important experimental results is the demonstration of beats of the ratchet current oscillations. We interpret these beats theoretically, assuming that they come

from SO splitting of the conduction band. The value of SO splitting extracted from the comparison of the experiment and theory is in good agreement with independent measurements of SO band splitting. We also discuss conditions required for the observation of magnetoplasmon resonances.

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