Helicity-Driven Ratchet Effect Enhanced by Plasmons

I. V. Rozhansky,[1,2] V. Yu. Kachorovskii,[1,3] and M. S. Shur[3]
[1]A.F. Ioffe Physical Technical Institute, Russian Academy of Sciences, 194021 St. Petersburg, Russia
[2]St. Petersburg State Polytechnic University, 29 Polytechnicheskaja, St. Petersburg 195251, Russia
[3]Center for Integrated Electronics, Rensselaer Polytechnic Institute, 110, 8th Street, Troy, New York 12180, USA
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We demonstrate that the ratchet effect—a radiation-induced direct current in periodically modulated structures with built-in asymmetry—is dramatically enhanced in the vicinity of the plasmonic resonances and has a nontrivial polarization dependence. For a circular polarization, the current component, perpendicular to the modulation direction, changes sign with the inversion of the radiation helicity. In the high-mobility structures, this component might increase by several orders of magnitude due to the plasmonic effects and exceed the current component in the modulation direction. Our theory also predicts that in the dirty systems, where the plasma resonances are suppressed, the ratchet current is controlled by the Maxwell relaxation.

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Plasmonics—a new branch of electronics—is a rapidly growing area of research with an extremely high potential for practical applications (see, for example, Refs. [1–7] and references therein). General ideas of plasmonics were formulated about 20 years ago and have been attracting much attention in the last decade. The main advantage of plasmonics compared to conventional electronics is a much higher speed of operation due to high values of plasma wave velocity, which is typically at least an order of magnitude larger than the saturated electron drift velocity. Another advantage is very easy tunability by the gate electrodes. These features are very promising for the fabrication of tunable electronic nanodevices capable of operating in a frequency range that is unaccessible by conventional electronic technologies. In particular, plasmonics is one of the candidates most expected to close the famous terahertz (THz) gap in the electromagnetic spectrum.

In this context, plasmonic oscillations in field effect transistors (FETs), which are the basic elements of modern technology, are especially interesting and have recently been the subject of a great number of studies focused both on fundamental and practical aspects (see reviews, Refs. [8,9] and references therein). An initial boost to this activity was given by a prediction that a dc current in the channel of a FET might become unstable with respect to the generation of plasma oscillations [10]. Such oscillations should lead to electromagnetic radiation at the plasma frequency. Furthermore, the nonlinear properties of the electron liquid in the FET channel can be used for rectifying the plasma oscillation induced by incoming electromagnetic wave [11]. The plasma wave velocity s in the FET channel can be tuned by the gate voltage. Its typical value, $s \sim 10^8$ cm/s, corresponds to the typical time scale of 10^{-12} s for the channel length ~ 1 μm. Thus, a FET in the plasma waves regime is expected to provide a tunable coupling to the electromagnetic radiation in the THz frequency range and can serve as a THz emitter or detector.

There is, however, a serious obstacle for such applications. In fact, the coupling with a single FET turns out to be quite weak. Indeed, the typical FET dimensions are 2 or more orders of magnitude smaller than the THz wavelength. Hence, a single device serves neither as an effective source nor as a detector with sufficiently high responsivity. The coupling significantly increases if there is a dc current flowing in the FET channel [12]. However, the dc current leads to the increase of the device noise.

One of the most promising ways to increase coupling is to use periodic structures (FET arrays, grating gate structures, and multigate structures). Such structures attract growing interest as simple examples of plasmonic crystals (PCs) [13–17]. They are much more appealing than single FETs in view of possible plasmonics applications and have already demonstrated excellent performance as THz detectors [18–22], in good agreement with the numerical simulations [23–26]. The first observations of the THz emission from the grating gate structures have been reported [27,28]. However, there are also some difficulties in providing a sufficiently strong dc response in PCs. The point is that nonzero dc photoresponse requires some asymmetry of the system, which would determine the direction of the produced dc current. In a single FET, such asymmetry might be induced by asymmetrical boundary conditions on the source and drain [10]. However, asymmetry of contacts to the whole PC does not provide an effective coupling mechanism. For coupling to be really noticeable, there should be a strong built-in asymmetry inside the unit cell of the PC.

The main purpose of this Letter is to propose an effective mechanism to induce strong asymmetry in the PC. The mechanism is related to the so-called ratchet effect [29–42]. The ratchet dc current $\mathbf{j}$ is induced by the electromagnetic wave impinging on the spatially modulated system provided that the wave amplitude is also modulated with the same wavelength but is phase shifted in space (see examples of such structures in Fig. 1). Although this

beautiful phenomenon has been known for a long time (for review, see Refs. [33,34,38]) its application to plasmonics has not yet been well studied.

For spatially modulated heterostructures, the kinetic theory of the ratchet effect neglecting excitation of plasmons was developed in Refs. [35–40,42]. The theory predicts the Drude peak of the width $1/\tau$ at $\omega = 0$ in the dependence $\boldsymbol{j}(\omega)$ (here τ is the momentum relaxation time and ω is the radiation frequency) and a smooth dependence of $\boldsymbol{j}$ on ω for $\omega \gg 1/\tau$ in agreement with numerical simulations [41].

Below, we demonstrate that the ratchet effect provides a fundamental universal mechanism of strong coupling of the electromagnetic radiation with plasmonic oscillations. The excitation of the plasmonic resonances dramatically increases the rectified dc current. We describe the plasmonic-enhanced ratchet effect in the frame of the hydrodynamic model valid for a collision-dominated regime, $\tau_{ee} \ll \tau$, where τ_{ee} is the electron-electron scattering time. This condition is easily realized for typical experimental parameters (see the discussion of scattering rates and typical plasmonic parameters in Ref. [[10]]). We find that dc current $\boldsymbol{j}(\omega)$ shows sharp resonances for $\omega \approx \omega_q$ provided that the plasmonic quality factor $\omega_q\tau$ is sufficiently large (here ω_q is the frequency of plasma oscillations in the structure spatially modulated with the wave vector q). Remarkably, not only the current component in the modulation direction, j_x (see Fig. 1), is enhanced by the plasmonic resonances, but the transverse component, j_y, is also enhanced and, moreover, for $\omega_q\tau \gg 1$ it becomes much larger than j_x. Another remarkable property is the strong polarization dependence of j_y. In particular, for a circularly polarized wave, j_y changes sign with the inversion of the helicity of polarization. For a single FET, the helicity driven response was measured [43] and explained theoretically [44] by assuming a special type of the boundary conditions. The dependence of the dc current on the helicity in the grating-gate periodic structures was also discussed in Refs. [35–38,40,42] within the approximation ignoring the plasmonic effects. Helicity-driven photocurrents have been measured in different materials [45–48] for purely electronic response only. In this Letter, we demonstrate that plasmonic effects greatly enhance the helicity-dependent part of the response. Exactly at the resonance, the transverse component scales as $j_y \propto (\omega_q\tau)^2$. Hence, in the high mobility structures, the helicity-dependent transverse ratchet might be several orders of magnitude larger then the previously predicted nonresonant ratchet effect. Even more interesting and unexpected is that for a circular polarization and for the external frequency away from the resonance point, j_y has a finite value in the ballistic limit, $\tau \to \infty$: $j_y \propto \omega_q^2/(\omega - \omega_q)^2$. In other words, we predict a giant helicity-driven nondissipative contribution to the ballistic ratchet effect, which diverges at $\omega \to \omega_q$.

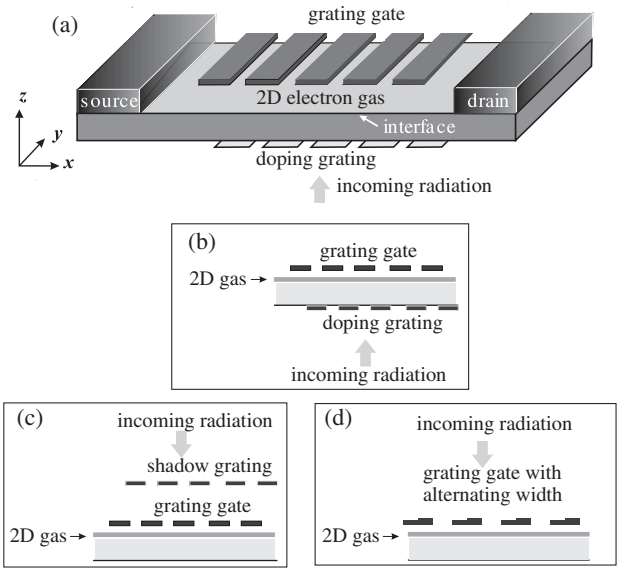

FIG. 1. Design of asymmetrical grating gate structures. Optical modulation can be achieved by fabrication of doping grating from the substrate side (a) [see (b) for side view] or shadow grating from the gate side (c). Also one can use a grating gate that has alternating width and alternating transparency (d).

The electron-electron interaction manifests itself in a nontrivial way even in the opposite limit, $\omega_q\tau \ll 1$, when plasma waves are damped by impurity scattering. One could expect that the response in this case would be given by the Drude peak at $\omega = 0$ with a width $\sim 1/\tau$ in accordance with Refs. [35–38,40,42]. As will be shown, surprisingly, the width of the peak is much narrower and is given by the Maxwell relaxation rate $1/\tau_M = \omega_q^2\tau \ll 1/\tau$.

Let us now specify the model. We will discuss the radiation-induced photocurrent in a structure with a common channel and a large-area grating gate (see Fig. 1). This structure represents a PC created by a modulated gate-to-channel potential. The ratchet dc current arises [35–42] as a result of a combined action of a static spatially periodic in-plane potential

$$V(x) = V_0 \cos(qx) \tag{1}$$

and the electric field of incoming radiation spatially modulated by a grating lattice with the same q [49]:

$$\boldsymbol{E}(t,x) = [1 + \hat{h}\cos(qx + \varphi)]\boldsymbol{e}(t). \tag{2}$$

Here $\boldsymbol{e}(t) = (e_x(t), e_y(t))$ is the in-plane oscillating vector with the components depending on the polarization of the wave, and $\hat{h}$ is the diagonal 2×2 matrix with the diagonal components h_x and h_y. These components describe the modulation depth of e_x and e_y, respectively.

The existence of nonzero average $\langle \boldsymbol{E}(\boldsymbol{E}\nabla V)\rangle_{t,x} \propto \sin\varphi$, implies that the dc current $\boldsymbol{j} = (j_x, j_y)$ controlled by the phase shift φ between $V(x)$ and $\boldsymbol{E}(t,x)$ might appear in the 2D liquid: $\boldsymbol{j} \propto \sin\varphi$. This phase shift provides the required asymmetry, so that the current reverses its direction when φ is incremented by π.

We consider the electron liquid in a 2D channel in the external field, Eq. (2), of general polarization:

$$e_x = E_{0x}\cos\omega t, \qquad e_y = E_{0y}\cos(\omega t+\theta). \tag{3}$$

The case $|E_{0x}| = |E_{0y}|$, $\theta = \pm\pi/2$ corresponds to the circular polarization. In the absence of perturbations ($V = 0$, $\boldsymbol{E} = 0$), the 2D electron concentration $N = N_0$ is controlled by the gate-to-channel voltage U_g:

$$N_0 = \frac{CU_g}{e}. \tag{4}$$

Here $C = \varepsilon/4\pi d$ is the gate-to-channel capacitance per unit area, ε is the dielectric constant, d is the spacer distance, and $e > 0$ is the absolute value of the electron charge. For smooth perturbations with $qd \ll 1$ Eq. (4) is also valid and relates the local concentration in the channel $N = N(x,t)$ and the local gate-to-channel swing. The total electric field in the channel is given by the sum of the external radiation field, static built-in field, and field arising due to the density perturbation: $\boldsymbol{E}_{\rm tot} = \boldsymbol{E} - \nabla V + (e/C)\nabla N$.

The quasiclassical dynamics of electrons in the channel obeys the kinetic equation:

$$\frac{\partial f}{\partial t} + \mathbf{v}\nabla f + \left[\boldsymbol{a} - \frac{e^2}{mC}\nabla N\right]\frac{\partial f}{\partial \mathbf{v}} = \mathrm{St}f, \tag{5}$$

where $\mathbf{a} = -(e/m)(\boldsymbol{E} - \nabla V)$, and $\mathrm{St}f$ is the collision integral including impurity and phonon scattering as well as electron-electron scattering. We study the electron liquid within the hydrodynamic approximation assuming the following hierarchy of the scattering times: $\tau_{ee} \ll \tau \ll \tau_{\rm ph}$, where $\tau_{\rm ph}$ is the electron-phonon scattering time. These inequalities allow one to search a solution as a Fermi-Dirac function in the moving frame $f = 1/[e^{m(\mathbf{v}-\boldsymbol{v})^2/2T-\mu/T}+1]$. This function depends on the local hydrodynamic parameters: velocity $\boldsymbol{v} = \boldsymbol{v}(\mathbf{r},t)$, chemical potential $\mu = \mu(\mathbf{r},t)$, and temperature $T = T(\mathbf{r},t)$. In what follows, we set $\mu \gg T$. This yields $N \approx \nu\mu$, where $\nu = m/\pi\hbar^2$ is the density of states. Having in mind that the electron-electron collisions conserve the particle number, momentum and energy, we multiply Eq. (5) by 1, $m\mathbf{v}$ and $m\mathbf{v}^2/2$ and integrate over $d^2\mathbf{v}$, thus obtaining the system of hydrodynamic equations:

$$\frac{\partial N}{\partial t} + \frac{\partial}{\partial x}(N\boldsymbol{v}) = 0, \tag{6}$$

$$\frac{\partial \boldsymbol{v}}{\partial t} + (\boldsymbol{v}\nabla)\boldsymbol{v} + \frac{\boldsymbol{v}}{\tau} = \boldsymbol{a} - \frac{e^2}{mC}\nabla N - \frac{\nabla W}{mN}, \tag{7}$$

$$\mathcal{C}\left[\frac{\partial T}{\partial t} + \mathrm{div}(T\boldsymbol{v})\right] = N\left(\frac{T_0 - T}{\tau_{\rm ph}} + \frac{mv^2}{\tau}\right), \tag{8}$$

where $W = \int d\epsilon\epsilon\nu[e^{(\epsilon-\mu)/T}+1]^{-1} \approx N^2/2\nu + \nu T^2\pi^2/6$ is the system energy per unit area in the moving frame, T_0 is the lattice temperature, and $\mathcal{C} = \nu T\pi^2/3$ is the heat capacity of the 2D degenerate electrons. Above, we implicitly assumed that τ is energy independent, which is the case for the short-range impurity potential.

Equation (8) is coupled to Eqs. (7) and (6) by the thermoelectrical force $\pi^2\nu\nabla T^2/6mN = \pi^2 T\nabla T/3m\mu$, whose contribution is suppressed in the highly degenerate electron gas. Let us estimate this force in the lowest order in T/μ. To this end, we neglect the lhs of Eq. (8) (which is small compared to its rhs due to the same parameter T/μ), thus arriving at a balance equation between Joule heating and phonon cooling: $mv^2/\tau = (T - T_0)/\tau_{\rm ph}$. Hence, the thermoelectrical force becomes $(\pi^2 T\tau_{\rm ph}/3\mu\tau)\nabla v^2$. Comparing this force with the term $(\boldsymbol{v}\nabla)\boldsymbol{v}$, we conclude that the former is negligible provided that $\mu/T \gg \tau_{\rm ph}/\tau$. Assuming that this inequality is fulfilled, we are left with the system of the hydrodynamic equations for velocity and concentration:

$$\frac{\partial n}{\partial t} + \frac{\partial v_x}{\partial x} = -\frac{\partial (nv_x)}{\partial x}, \tag{9}$$

$$\frac{\partial v_x}{\partial t} + \frac{v_x}{\tau} + s^2\frac{\partial n}{\partial x} = a_x - v_x\frac{\partial v_x}{\partial x}, \tag{10}$$

$$\frac{\partial v_y}{\partial t} + \frac{v_y}{\tau} = a_y - v_x\frac{\partial v_y}{\partial x}, \tag{11}$$

where $n = (N - N_0)/N_0$.

The rhs of Eqs. (9), (10), and (11) include perturbation $\boldsymbol{a}$ as well as nonlinear terms. Assuming that $\boldsymbol{a}$ is small, one can search a solution as a perturbation series over $\boldsymbol{a}$: $n = n^{(0,1)} + n^{(1,0)} + \cdots$, $\boldsymbol{v} = \boldsymbol{v}^{(0,1)} + \boldsymbol{v}^{(1,0)} + \cdots$ Here the two indices denote the order of smallness with regard to $\boldsymbol{e}$ and V_0, respectively. The nonzero dc current $\boldsymbol{j} = -eN_0\langle(1+n)\boldsymbol{v}\rangle_{t,x}$, appears in the third order with respect to $\boldsymbol{a}$ (second order in $\boldsymbol{e}$ and first order in V_0): $\boldsymbol{j} \approx \boldsymbol{j}^{(2,1)}$ (here $\langle\cdots\rangle_{t,x}$ stands for the time and space averaging [51]). Importantly, Eqs. (9) and (10) can be solved independently from the decoupled Eq. (11) [the latter can be solved after the solution of Eqs. (9) and (10) is found]. The details of the calculations are presented in Ref. [52]. Here we estimate one of the terms contributing to the $j_x^{(2,1)}$ in order to clarify the key points of derivation.

The static potential, Eq. (1), leads to density modulation $n^{(0,1)} \propto V_0\cos(qx)$. The homogeneous part of the field, Eq. (2), does not affect concentration but leads to the Drude peak in the velocity: $v_x^{(1,0)} \propto E_{0x}[e^{i\omega t}(i\omega + 1/\tau)^{-1} + \mathrm{H.c.}]$ (we omit here inhomogeneous contribution). Substituting these equations into the nonlinear term $\partial[n^{(0,1)}v_x^{(1,0)}]/\partial x$ in the rhs of Eq. (9), and solving Eqs. (9) and (10) we find that the velocity in the order (1,1) exhibits plasmonic resonance as well as the Drude peak: $v_x^{(1,1)} \propto E_{0x}V_0\cos(qx)\times[e^{i\omega t}(i\omega + 1/\tau)^{-1}(\omega^2 - \omega_q^2 - i\omega/\tau)^{-1} + \mathrm{H.c.}]$. Here $\omega_q = sq$ is the plasma wave frequency. In turn, the nonhomogeneous part of the field, Eq. (2), also excites the plasmonic resonance, thus leading to density correction $n^{(1,0)} \propto E_{0x}h_x\sin(qx+\varphi)[e^{i\omega t}(\omega^2 - \omega_q^2 - i\omega/\tau)^{-1} + \mathrm{H.c.}]$. Combining these equations, we find the nonvanishing correction to the dc current in the order (2,1):

$$j_x^{(2,1)} \propto \langle n^{(10)}v_x^{(1,1)}\rangle_{t,x} \propto \frac{\tau}{1+\omega^2\tau^2}\frac{\sin\varphi}{(\omega^2 - \omega_q^2)^2 + \omega^2/\tau^2}.$$

More detailed calculations [52] yield

$$j_x = j_{0x}\frac{2\omega_q^5\tau}{(1+\omega^2\tau^2)[(\omega^2-\omega_q^2)^2+\omega^2/\tau^2]}, \quad (12)$$

$$j_y = j_{0y}\frac{\omega_q^3[(\omega^2-\omega_q^2)\tau\cos\theta+\omega\sin\theta]}{(\omega^2-\omega_q^2)^2+\omega^2/\tau^2}. \quad (13)$$

Here $j_{0x} = e^4V_0N_0E_{0x}^2h_x\sin\varphi/(4m^3s^3\omega_q^2)$ and $j_{0y} = -e^4V_0N_0E_{0x}E_{0y}h_y\sin\varphi/(4m^3s^3\omega_q^2)$ are frequency- and disorder-independent currents that are proportional to the asymmetry factor $\sin\varphi$ and are sensitive to the polarization of the radiation. We note that the finite value of j_y implies that the electric circuit is closed in the y direction. For a disconnected circuit, the voltage would develop instead, which is analogous to the Hall voltage and, therefore, depends on the geometry of the system.

As seen from Eqs. (12) and (13), there are two different regimes depending on the plasmonic quality factor $\omega_q\tau$. For $\omega_q\tau \gg 1$, the response is peaked both at $\omega = 0$ and at $\omega \simeq \omega_q$ within the frequency window $\sim 1/\tau$. In the vicinity of the plasmonic resonance $\omega \simeq \omega_q$, one can simplify Eqs. (12) and (13):

$$j_x \approx j_{0x}\frac{2\omega_q\tau}{1+4(\omega-\omega_q)^2\tau^2}, \quad (14)$$

$$j_y \approx j_{0y}\frac{\omega_q^2\tau^2[\sin\theta+2(\omega-\omega_q)\tau\cos\theta]}{1+4(\omega-\omega_q)^2\tau^2}. \quad (15)$$

In the opposite nonresonant case, $\omega_q\tau \ll 1$, we find

$$j_x \approx \frac{2\omega_q\tau j_{0x}}{1+\omega^2\tau_M^2}, \qquad j_y \approx \frac{\omega_q\tau(\omega\tau_M\sin\theta-\cos\theta)j_{0y}}{1+\omega^2\tau_M^2}$$

where the width of the response, $1/\tau_M = \omega_q^2\tau$, is determined by the inverse time of the charge spreading at the distance $\sim q^{-1}$ (Maxwell relaxation time). Surprisingly, this width is much smaller, than the inverse momentum relaxation time: $1/\tau_M \ll 1/\tau$. In other words, unexpectedly, the peak turns out to be much narrower than the Drude peak.

In the resonant regime $\boldsymbol{j}$ is much larger than in the nonresonant case (due to the largeness of $\omega_q\tau$) and shows a sharp resonant dependence on ω (see Fig. 2). Hence, the plasmon excitation leads to a dramatic enhancement of the ratchet effect. Note that for $\theta = \pm\pi/2$, j_y changes its sign with the sign of θ, i.e., when switching between right and left circular polarizations. Thus, our results predict a strong helicity effect—the circular polarization of the incident light determines the direction of j_y. Remarkably, for clean systems, the transverse component of the current, $j_y^{\max}/j_{0y} \sim (\omega_q\tau)j_x^{\max}/j_{0x}$, might be much larger than the longitudinal one provided that $\omega_q\tau$ is sufficiently large. For $\omega = \omega_q$, we find $j_y \propto (\omega_q\tau)^2$. Therefore, in clean systems, the transverse current might increase by several orders of magnitude. We also see that for $\theta = \pm\pi/2$ the transverse current remains finite in the dissipationless limit $\tau \to \infty$: $j_y \to \pm j_{0y}\omega_q^3\omega/(\omega^2-\omega_q^2)^2$. Hence, we predict the existence of a large helicity-dependent nondissipative contribution to the ratchet effect which infinitely increases as we approach the resonance.

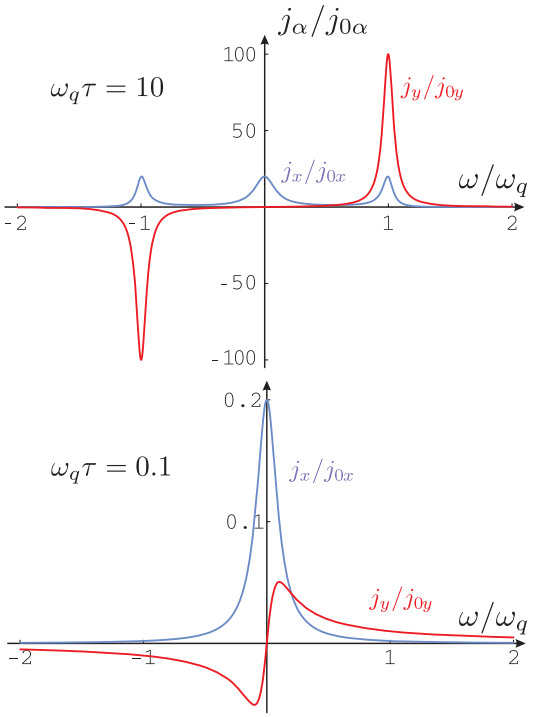

FIG. 2 (color online). Frequency dependence of current components j_α [$\alpha = (x, y)$] in the resonant (upper panel, $\omega_q\tau = 10$) and nonresonant (lower panel, $\omega_q\tau = 0.1$) cases for circular polarization ($\theta = \pi/2$, $E_{0x} = -E_{0y}$, $h_x = h_y$).

To conclude, we predicted a dramatic enhancement of the ratchet effect due to the excitation of plasmonic resonances. We identified a helicity-dependent contribution to the ratchet current and found that this contribution increases with decreasing the static disorder and saturates in the limit $\tau \to \infty$. We also demonstrated that the nonresonant ratchet current is sharply peaked at zero frequency within the width on the order of the inverse Maxwell relaxation time. The results are quite general and can be applied for conventional semiconductor structures based on GaAs and Si as well as for novel materials such as graphene and other van der Waals 2D materials. Hence, our work "builds a bridge" between the fundamental ratchet effect and various applications in the rapidly growing field of plasmonics.

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[49] Note that for realistic systems the density and field perturbations cannot be described by simple harmonic functions and involve infinite number of harmonics (see Ref. [50]). Also, the field amplitude $\boldsymbol{e}$ is much smaller than the amplitude of the external field due to the screening by the gate electrodes. However, the simplified model based on Eqs. (1) and (2) with phenomenological parameters V_0, E_{0x}, and E_{0y} captures the key physics of the problem and is sufficient for clarifying the basic concept of the plasmon-enhanced ratchet.
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[51] In fact, averaged in time, $\boldsymbol{j}$ does not depend on x, so that it is sufficient to make averaging over t only. However, calculations are strongly simplified if we also make x-averaging in all contributing terms.
[52] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.114.246601 for technical details of calculation.