Temporal Wood Anomalies - Smoothing the Path to the Near-Field

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Free space supports propagating waves, whose divergenceless character leads to the well-known diffraction limit. In order to confine waves to subwavelength volumes, advances in fabrication techniques have enabled subwavelength structuring of matter, achieving near-field control of light and other types of waves. The price is often expensive fabrication needs, as well as the introduction of impurities, a major contributor to losses. In this Letter, we propose temporal inhomogeneities, such as a periodic drive in the electromagnetic properties of a surface which supports guided modes, to circumvent the need for subwavelength fabrication in the coupling of propagating waves to evanescent modes across the light line, achieving the temporal counterpart of the Wood anomaly. We show how this concept is valid for any material platform and at any frequency, and propose and model a realistic experiment in graphene, demonstrating that time-modulation of material properties could be a tunable, lower-loss alternative to the subwavelength structuring of matter for near-field wave control.

The idea of trapping light near a surface finds its roots in the curious discovery by Wood that the reflection of light from a metallic grating vanishes at a specific angle of incidence, as diffraction theory alone seemed unable to explain [1]. Simultaneously, theories were developed, which predicted the existence of evanescent modes near metallic surfaces [2–6], although only half a century later the link between these two phenomena was elucidated [3]. This scattering anomaly was thus recognised as the coupling of the incoming photons to surface waves guided along a metal-dielectric interface by the free electron plasma [5–7].

Over the past decades, advances in nano-fabrication have enabled unprecedented progress in the harvesting and control of electromagnetic waves near surfaces or layered materials, with several types of surface excitations available, such as surface plasmon, phonon and exciton polaritons [8–9]. Similarly, in the RF and microwave regimes, the engineering of metasurfaces enables the realisation of effective impedance surfaces which support guided waves with tailored properties [10–11].

By their definition, surface waves are evanescent along at least one spatial dimension. The consequent imaginary out-of-plane momentum component, \(k_z = \sqrt{\omega^2/c^2 - k_x^2}\), forces the in-plane momentum \(k_x\) to be larger than the total momentum \(\omega/c\) of any photon propagating in free space. Here \(\omega\) and \(c\) are the frequency of the wave and the speed of light in the medium. This implies that the coupling of these two waves is conditional to the bridging of this momentum mismatch, which needs to suffice to couple a propagating wave, within the light cone, to an evanescent one, outside of it. In free space this is conventionally done by breaking the translational invariance of the surface either by means of a near-field probe, such as a tip, or, as Wood did, by periodically structuring the surface of the metal, or its surroundings [12]. Alternative schemes have been proposed, although they rely on inherently weak nonlinear processes [13–14].

While spatial design and fabrication has paved the way to the discovery and harnessing of novel near-field phenomena, the limited precision of spatial fabrication methods stands as a major obstacle to the decade-long challenge to reduce losses in nanophotonics. On the other hand, the recent advent of tunable and ultra-thin van der Waals layered materials, such as graphene, and thin-film semiconductors such as indium tin oxide [8–9] has sparked the development of new schemes to tune and modulate the electromagnetic properties of a material in time, empowering the field of metamaterials and metasurfaces with the potential to exploit time as a new degree of freedom for the control of light [10–11]. Similarly, at RF and microwave frequencies, active capacitive elements can be incorporated in a metasurface to achieve temporal and spatiotemporal control of electromagnetic waves [17–19,23], while similar efforts have been made in the acoustics [24], elasticity [25,26] and water-wave [27] communities, each featuring challenges and advantages specific to the physics involved.

In this Letter, we propose the temporal counterpart of the Wood anomaly as a new mechanism to excite surface waves without need for in-plane structuring. By periodically modulating the surface impedance of a thin film, we demonstrate that surface waves can be excited with unit efficiency from the far-field, with no spatial inhomogeneity on their surface, and suggest and model a realistic experiment based on a time-modulated graphene sheet.

Our results provide an alternative path towards the coupling of radiation to surface waves, which circumvents the need for near-field probes, subwavelength grating structures and/or advanced fabrication. This concept bears the potential to dramatically reduce losses, as well as costs, in current nanophotonics setups, whilst enabling...
FIG. 1. (a,b) Conventional coupling of propagating waves to surface waves is achieved by breaking momentum conservation via a periodic perturbation in space, inducing horizontal transitions of the radiation field in the dispersion diagram. Conversely, periodic perturbations in time (a,c) break energy (frequency) conservation, enabling (d) far-field excitation of surface waves, without the need for in-plane structuring.

wider tunability for metasurfaces across the electromagnetic spectrum, and in other wave systems.

This Letter is structured as follows: we first generalise the original concept of the Wood anomaly to the temporal dimension, demonstrating analytically and verifying numerically its capabilities for the trapping of radiation into surface waves for the case of a time-modulated, spatially homogeneous, general conducting layer. We then propose a realistic implementation in graphene, and show that perfect coupling can be achieved with the inclusion of a back-reflector, as a result of the interaction between a cavity mode and a graphene plasmon enabled by the temporal grating.

Fig. 1 (panels a and b) depicts the conventional scattering process originally observed by Wood: a TM-polarized, propagating wave with in-plane wavevector $k_i$ and frequency $\omega_{sw}$ = $c k_i / \cos(\theta)$, where $\theta$ is the angle measured from the surface plane, impinges on a periodically structured surface which supports a guided mode with wavevector $k_{sw}$ > $\omega_{sw}/c$. If the surface is patterned by a periodic grating whose wave number $g = 2 \pi / d$, where $d$ is the spatial period, matches the difference between the momentum of the surface wave $k_{sw}$ and that of the photon, $k_i$, then the two modes can be coupled, and a surface wave is launched. This process is described by the red arrow in Fig. 1a.

As a counterpart of this process, if the electromagnetic parameters of a smooth surface are periodically modulated in time with frequency $\Omega$ (Fig. 1b-c), the temporal analogue of a Wood anomaly takes place: an incoming wave with frequency $\omega_i$ and wavevector $k_{sw} = \omega_i \cos(\theta)/c$ is now efficiently coupled to a surface mode if the modulation frequency $\Omega$ matches the difference $\omega_i - \omega_{sw}(k_{sw})$ between the frequency of the incoming photon and that of the surface mode at the incoming wavevector.

Our model system consists of a conductive sheet of infinitesimal thickness, whose surface current density $J$, satisfies the Drude equation:

$$\frac{dJ}{dt} + \gamma J(x,t) = W_D(t) E_z(x,t),$$

where $W_D(t)$ is the Drude weight, which quantifies the density of charge carriers at any given instant of time, and $\gamma$ is a phenomenological dissipation rate. This description is valid for a general Drude sheet of subwavelength thickness. We assume a harmonic modulation of the Drude weight: $W_D(t) = W_{D,0} (1 + 2 \alpha \cos(\Omega t))$.

Since the system is periodic in time, we can assume Floquet solutions of the form $\psi_k(x,t) = e^{i(\omega t - k x)} \sum_j \psi_n e^{i \omega n t}$ for the electric field $E = E_x \hat{x} + E_z \hat{z}$, the magnetic field $H = H_y \hat{y}$, and the current $J = J \hat{x}$. Solving Maxwell’s Equations on the two sides of the current sheet, and imposing continuity of the in-plane electric field $E_x$ and discontinuity of the magnetic field $H_y$ by the surface current $J$, we arrive at the system of equations:

$$\mathbb{D} E_x^{tra} + \mathbb{F} E_x^{tra} = E_x^{inc},$$

where $E_x^{inc}$ and $E_x^{tra}$ are vectors which contain the Fourier amplitudes of the incident and transmitted components of the in-plane electric field, and $\mathbb{D}$ and $\mathbb{F}$ are diagonal and tridiagonal matrices respectively, whose elements read:

$$D_{n,n'} = \frac{1}{2} \left[ 1 + \frac{\varepsilon_1 k_{x,n}^2}{\varepsilon_2 k_{x,n}^2} \right] \delta_{n,n'}$$

$$F_{n,n'} = -\mu_0 \varepsilon_0 k_{z,n}^2 W_{D,0} \left[ \delta_{n,n'} + \alpha (\delta_{n,n' + 1} + \delta_{n,n' - 1}) \right],$$

where $\varepsilon_1$ and $\varepsilon_2$ are the relative permittivities of the lower and upper half spaces [See Fig. 1d)], $c_0$ is the speed of light in vacuum and $k_{z,n} = \sqrt{\varepsilon_j (\omega + n \Omega)^2 / c_0^2 - k_{x,n}^2}$.

Defining the matrix on the LHS $T^{-1} = \mathbb{D} + \mathbb{F}$, we can thus calculate the vector of transmitted amplitudes $E_x^{tra} = T E_x^{inc}$ numerically, by matrix inversion.

Furthermore, by truncating the matrices in Eq. 2 to three modes, where we account for the incoming, propagating wave and the two down-converted modes, and assuming plane-wave incidence and symmetric surrounding media $\varepsilon_1 = \varepsilon_2$, we can reduce this problem to the linear system:

$$\begin{pmatrix}
1 + i \xi_{-2} & i \alpha \xi_{-2} & i \alpha \xi_{-1} \\
-i \alpha \xi_{-1} & 1 + i \xi_{-1} & i \alpha \xi_{-1} \\
i \alpha \xi_{0} & i \alpha \xi_{0} & 1 + i \xi_{0}
\end{pmatrix}
\begin{pmatrix}
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\end{pmatrix}
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0 \\
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which can be solved explicitly to give an approximate analytic solution, accurate in the weak-coupling regime, for the scattering amplitude of the (propagating) transmitted wave:

$$\frac{E_{x,0}^{tra}}{E_{x,0}^{inc}} = \left[ \frac{\alpha^2 \xi_{0} \xi_{-1}}{1 + \xi_{-2} + (1 + i \xi_{-1})} \right]^{-1},$$

(5)

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Fig. 2. (a) Dispersion relation for surface waves on a current sheet (or deeply subwavelength conducting film) with Drude dispersion (Eq. 6). (b) Transmittance through the current sheet under a time-modulation of its Drude weight at frequency $\Omega \approx 0.94$, with 30% modulation amplitude ($\alpha = 0.15$) at fixed in-plane wavevector $k_x \approx 1.12$ (vertical red line). A surface mode is excited when the frequency of the incoming waves differs from that of the surface wave by the modulation frequency $\Omega$. Continuous line: analytic theory (3-bands approximation). Dots: exact transfer matrix result. (c-e) FETD simulations of the temporal Wood anomaly predicted in (a-b). (c) A pulse with the carrier frequency $\omega_0 \approx 1.61$ impinges on the time-modulated surface ($\alpha = 0.15$), successfully exciting a surface wave. In (d), the same pulse impinges on the surface without modulation ($\alpha = 0$), and in (e) the modulation is present ($\alpha = 0.15$), but the pulse is detuned in frequency, whilst having the same in-plane momentum.

where $\xi_n = \frac{\mu_0 \varepsilon_{n}^{2} k_{0}^{(1)} W_{D,0}}{2(\omega+n\Omega)(\gamma-i(\omega+n\Omega)^{2})}$.

Fig. 2(a) shows the exact dispersion relation for surface modes on a flat, unmodulated Drude sheet of sub-wavelength thickness [12]:

$$\frac{\varepsilon_1}{\sqrt{k_x^2 - \varepsilon_1 \omega^2}} + \frac{\varepsilon_2}{\sqrt{k_x^2 - \varepsilon_2 \omega^2}} - \frac{1}{\omega^2} = 0,$$

(6)

where we defined dimensionless units $\tilde{\omega} = \omega/(W_{D,0} Z_0)$ and $\tilde{k}_x = k_x c/(W_{D,0} Z_0)$, where $Z_0 = \sqrt{\mu_0/\varepsilon_0}$ is the impedance of free space, and we assumed, for simplicity, that the sheet is surrounded by vacuum. The orange feature shows the dispersion of a surface mode: fixing a wavevector $k_x \approx 1.12$ (red vertical line), we see that a surface mode exists, with frequency $\tilde{\omega}_{sw} \approx 0.67$; clearly, this mode is inaccessible by any propagating wave, whose dispersion lies within the light cone, unless either in-plane translational symmetry, or, as we set out to show, temporal symmetry, are broken. Assuming a modulation frequency $\Omega \approx 0.94$ and calculating the transmission amplitude for incoming (propagating) waves of different frequencies reveals a clear transmission dip when $\tilde{\omega} \approx \tilde{\omega}_{sw}(\tilde{k}_x) + \Omega \approx 1.61$, showing that incident radiation has coupled to a surface wave. Furthermore, the accuracy of our approximate analytic solution despite up-conversion being neglected highlights how, although the modulation is time-reversal-symmetric, frequency down-conversion strongly dominates the scattering process over high-harmonic generation. This fact highlights the potential of this strategy for the enhancement of parametric down-conversion.

In order to verify the frequency-domain results, we then show finite-element-time-domain (FETD) simulations (COMSOL), fully accounting for material dispersion [See Supplemental Material]. We consider an incident pulse with carrier frequency $\omega_0 \approx 1.61$, incident at an angle $\theta = 44^\circ$, matching in-plane momentum. Fig. 3(c) shows a snapshot of the absolute value of the electric field after the scattering has occurred: The pulse is partly transmitted through and partly reflected from the sheet, leaving behind, however, a surface wave which propagates along the sheet, despite the absence of any spatial inhomogeneity. [Full animations available in the supplementary material]. For completeness, we also show in Fig. 3(d) that no surface wave is excited in the absence of time-modulation under the same illumination, and that (e) the same modulation amplitude $\alpha = 0.15$ does not couple a pulse with detuned carrier frequency $\omega_0 \approx 2.35$ and identical in-plane momentum.

This concept is completely general to any translationally invariant surface which supports a guided mode. Potential photonic implementations are readily possible by using modulated spoof metasurfaces at frequencies in the RF and mm-wave range, where modulation up to speeds of a few GHz are possible [20, 28]. At higher frequencies all-optical approaches may enable this effect to be realised in the terahertz, on graphene, via electro-optic modulation, and potentially in the telecommunication band, by exploiting nonlinearities combined with the wide tunability of epsilon-near-zero films such as indium tin oxide [15, 16, 18, 29, 31]. Here, we propose a terahertz implementation with surface plasmons in single-layer graphene, which is well within current experimental capabilities [13].

Graphene has established itself as a widely tunable material, with reported doping modulation amplitudes of $\approx 38\%$ and $\approx 2.2$ ps response times [30, 32], and a number of schemes have been proposed to exploit temporal control of the carrier density of graphene, both by making use of a time-varying bias [36]. Whilst the quality of graphene has long been known to deteriorate with nanofabrication due to the resulting introduction
FIG. 3. (a) Transmittance (continuous lines) and absorbance (dashed lines) spectra for a graphene sheet whose doping is modulated at an angular frequency $\Omega = 2\pi \times 80$ GHz, with amplitudes of 0 (blue), 10 (red), 20 (purple) and 30% (green). The baseline Fermi level $E_F = 1$ eV, the electron mobility $m = 20\,000$ cm$^2/(V\cdot s)$ and the incident in-plane wavevector $k_x = 15$ rad$\cdot$mm$^{-1}$. The inset shows the transition in the phase plane.

of defects and edges, surface plasmons in state-of-the-art, homogeneous graphene can achieve measured propagation lengths of tens of microns, surviving up to 50 oscillation periods [37]. Could a temporal grating excite graphene plasmons without the need for any surface structure?

Surface plasmons in graphene are generally strongly confined. However, working sufficiently close to the light line offers an opportunity to apply this concept with realistic modulation rates. In Fig. 3 we calculate the transmittance and absorbance spectra for incoming THz waves with in-plane wavevector $k_x = 15$ rad$\cdot$mm$^{-1}$ incident on a single graphene sheet, whose baseline Fermi level $E_{F,0} = 1$ eV, corresponding to a baseline Drude weight $W_{D,0} = \frac{e^2}{\pi \hbar} E_{F,0} \approx 120$ GHz/$\Omega$, sinusoidally modulated in time at a frequency $\Omega = 2\pi \times 80$ rad$\cdot$GHz. We assume a conservative electron mobility $m = 20\,000$ cm$^2/(V\cdot s)$, corresponding to a loss rate $\gamma = ev_F^2/mE_F \approx 0.45$ THz, where $v_F = 9.5 \times 10^7$ cm/s is the Fermi velocity of the charge carriers. A close-up of the dispersion relation and the coupling process is shown in the inset. From the parameters above, we predict a measurable signal, in transmission, of order ranging between 5% and 30%.

The coupling efficiency of this mechanism is measurable, but not perfect. Many graphene setups include a metallic gate, which is typically used to adjust the doping level of the sample [35]. Can we use such a metal-backing to couple the plasmon resonance in Fig. 3 to a cavity mode, and thus excite graphene plasmons with unit efficiency?

Fig. 4(a) illustrates the scattering problem including a back-reflector consisting of perfect electric conductor (PEC) which is a good approximation for gold at terahertz frequencies. The total reflection coefficient is straightforwardly calculated by summing the multiple-scattering series for the amplitudes computed via the previous single-layer transfer-matrix calculation, which gives [39]

$$R_{\text{tot}} = R_{21}^t + T_{21}^t P (I - R_{12}^t P)^{-1} T_{12}^t,$$

(7)

where the different matrices containing the reflection and transmission coefficients for the Fourier amplitudes are defined in Fig. 4(a), I is the identity matrix, and $P$ is a diagonal matrix with elements $P_{n,n'} = R_{1m}^t e^{2ik_{z,n}d}$ where $R_{1m}^t$ is the reflection coefficient at the metal surface, which is simply $-1$ for PEC, while the exponential accounts for the round-trip phase accumulation (or decay) of the $n^{th}$ Fourier mode inside the cavity, which has width $d$. We assume the same parameters for the graphene and the dielectrics as in Fig. 3.

The resulting reflection spectrum is given in Fig. 4(b) as a colour plot against frequency $f$ and cavity width $d$. The oblique reflection dip shows a Fabry-Perot (FP) mode of the cavity, whereas the vertical one shows the surface plasmon mode (SPP). As the two resonances cross, they hybridize, achieving perfect coupling between the incoming radiation and the graphene surface plasmon.
with no breaking of translational symmetry. A blow-up of the anti-crossing arising from the interaction between the two modes is given in the inset in log-scale, demonstrating that perfect coupling is achieved. Finally, three cut-lines are shown in (c), corresponding to cavity widths \( d = 0.40, 0.44 \) and 0.48 mm, showing how the coupling to the distinct SPP and FP modes (dotted) becomes perfect as they strongly interact when \( d \approx 0.44 \) mm (dashed), and the plasmon becomes dark beyond the anti-crossing (dashed).

In this work we have proposed the temporal counterpart of the Wood anomaly as a general new strategy to couple free-space radiation to surface waves while circumventing the conventional needs for surface fabrication and near-field coupling schemes. Conversely to the momentum coupling introduced by Wood, this scheme relies on the down-conversion of a photon with sufficiently high frequency to match the momentum of a surface mode, whilst being able to propagate in free-space. By exploiting a weak temporal modulation of the impedance of a spatially homogeneous surface, we have shown analytically and numerically how surface waves can be excited on a translationally invariant surface.

Our model accurately accounts for material dispersion, and we have proposed a terahertz implementation in graphene, which is readily accessible with state-of-the-art experimental setups, showing that perfect coupling between radiation and graphene plasmons can be achieved with the addition of a flat back-reflector. To conclude, exploiting the recent introduction of tunable materials, this new coupling scheme may prove a valuable asset in the current struggle to reduce losses in nanophotonics while enabling the highest degree of reconfigurability, providing a new perspective in the use of temporal inhomogeneities to enhance light-matter interactions.

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