

Electromagnetic time interfaces in wire media: Innovations for subwavelength imaging

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In this work, we theoretically study temporal interfaces between media with strong spatial dispersion and dielectrics. In particular, we consider a temporal discontinuity that transforms a wire medium sample, a metamaterial with resonant spatial dispersion, into a uniaxial dielectric. We show that this transition results in a transformation of the deeply subwavelength spatial spectrum of TEM waves propagating in the wire medium at a certain frequency into a spectrum of plane waves at new frequencies, that are all higher than the initial one. The waves at different frequencies propagate in different directions. Their complex amplitudes and propagation directions are uniquely related to the amplitudes of the spatial harmonics of the field, which existed before the transition. We explain how to implement this transition. The revealed effect may result in a promising method of subwavelength imaging.

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I. INTRODUCTION

In recent years, there has been a surge of interest in electromagnetic time-varying systems, e.g., Refs. [1,2], particularly focusing on temporal discontinuities [3]. This research domain considers instantaneous changes at some moments of effective parameters characterizing the electromagnetic response of a system. In the realm of electromagnetic media, these abrupt jumps, such as in the permittivity tensor, give rise to time interfaces between different media. Generation of time-reflected waves, frequency conversion, breaking the power conservation [4–6], and the creation of photon pairs [7,8] are four fundamental features that were considered theoretically [4–9] and checked experimentally [10–13]. Time interfaces in complex media may be used in many applications such as polarization conversion [14], polarization splitting [15], polarization rotation [16], analogous computing [17], inverse prism [18], antireflecting coatings [19,20], temporal aiming [21], wave freezing and thawing [22], wiggler-mode production [9,23,24], transformation of surface waves into radiation [22,25], direction-dependent wave manipulation [16], and so on.

Despite extensive exploration in this field, investigations into time interfaces of media with strong spatial dispersion remain conspicuously absent. One specific example of those media is a metamaterial of thin parallel wires called wire medium (WM), whose electromagnetic properties and potential applications were thoroughly reviewed in Ref. [26]. Among these applications, the most relevant for the present

paper is the radio-frequency, microwave, THz, and infrared WM-based endoscopy, which grants subwavelength resolution of imaging devices.

WM-based endoscopy is a chapter of far-field subwavelength imaging. Far-field subwavelength imaging is a set of imaging methods in which the so-called Abbe diffraction limit (half-wavelength spatial resolution of two ultimately small scatterers separated by a tiny gap) is beaten. It is the subject of a large body of literature that describes linear, nonlinear, direct, and indirect methods that grant superresolution (spatial resolution finer than $\lambda/2$) in the far zone of the imaged object. Recently, this abundant literature has been reviewed in overview articles [27,28].

Among the many known methods of far-field subwavelength imaging via postprocessing of diffraction patterns, it is worth noting recent developments in diffraction grating ultramicroscopy (Łukosz microscopy) [29–31]. These developments utilize additional opportunities granted by time management. In work [32] the authors achieved similar superresolution as in Ref. [31] without a silver nanolayer coupled to the grating. Instead, they utilized the high-frequency tails of a quasimonochromatic laser pulse (excitation in form of the product of a time-harmonic function and a step function). In work [33], superresolution is achieved using a controllable diffraction grating whose period is changed over time. We mention these works in the context of the present research because in both cases a temporal change (of the incident wave amplitude or of the diffraction grating period) is involved to obtain superresolution. In our case, we study the impact of an instantaneous change in the environment in which the signal carrying out subwavelength information propagates.

In WM endoscopes propagating wave packages with high spatial frequencies (large transverse wave numbers) transfer information about tiny features of an object to electromagnetically long distances along the wires. The development of these devices started from work [34]. Conceptual microwave studies (e.g., Refs. [35,36]) lead to creation of wire

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medium endoscopes that offer THz and infrared imaging with deep subwavelength resolution [37–40]. The subwavelength resolution of WM endoscopes operating in different frequency ranges was experimentally demonstrated in papers [34,35,37,39–42]. If these endoscopes are based on silver microwires located in a porous plastic (THz range) or silver nanowires in a fiber glass (infrared range), deeply subwavelength details of electric field distributions can be replicated at substantial distances from the object, much larger than the free-space wavelength λ . The ultimate resolution of such WM endoscopes is close to the period a of the WM. For existing THz and infrared endoscopes $a \sim 0.1\lambda$ but, in principle, this period can be smaller [43]. However, such THz and IR endoscopes suffer from restrictions related to the detection of subwavelength images. To capture the image, one needs detectors with the size a or smaller. In conceptual experiments with THz and IR WM endoscopes, researchers used raster scanning, photoresist layers or photoluminescent markers [39–41]. These methods are not suitable for *in situ* measurements.

Most interesting WM endoscopes magnify subwavelength images. Such endoscopes are WM-filled cables of a conical shape with gradually increasing distances between the wires, which are also conical. Their operation is analogous to that of commonly known hyperlenses, where wave packages spread in space with the propagation path, so that their spatial spectrum shrinks [44–46]. Unfortunately, the spatial resolution of such endoscopes is only slightly subwavelength (of the order of $\lambda/4$) [37,38,47]. This is so because the technology of three-dimensional (3D) tapering of wire media [37] demands a liquid metal such as indium instead of silver. Ohmic losses in such wires restrict not only the spatial resolution, they also restrict the maximal length of such endoscopes by nearly one cm [37]. This means that the length is comparable with the diameter of the conical base, which makes magnifying WM endoscopes nonflexible. These two drawbacks hinder the use of magnifying WM endoscopes in THz endoscopy.

In this paper, we propose an alternative approach to THz subwavelength imaging. We suggest using a nontapered switchable WM endoscope, which does not need to have any cable that delivers subwavelength information to long distances. Instead, a finite-length switchable WM sample sends a wave pulse into surrounding space, and the frequency spectrum and the radiation pattern of the pulse encode the subwavelength features of the object under test. The subwavelength image of an object can be retrieved from the spectroscopic measurement of the angular intensity distribution of this pulse.

The idea of this method results from the effect that we have theoretically revealed by studying the instantaneous transition of continuous wires into split wires. Such temporal discontinuity transforms a strongly spatially dispersive wire medium into an artificial uniaxial dielectric—a medium with negligible frequency and spatial dispersion. For an infinitely extended WM we deduce the amplitudes of the time-refracted and time-reflected waves at new frequencies and analytically show that at such temporal interface, the full conversion of subwavelength field distributions into propagating waves holds. The analytical expressions for the radiated fields show that the spatial variations of the field distribution that exist before

the switching are encoded in both the frequency spectrum and angular pattern of the propagating wave package.

II. TEMPORAL INTERFACE BETWEEN WIRE MEDIA AND UNIAXIAL DIELECTRICS

In a wire medium (WM) of perfectly conducting wires, eigenmodes at frequencies below the so-called plasma frequency ω_p are transverse electromagnetic (TEM) waves propagating along the wires without dispersion. Both phase and group velocities along the wires are equal to c/n , where $c = 1/\sqrt{\epsilon_0\mu_0}$ is the speed of light in vacuum, and $n = \sqrt{\epsilon_h}$ is the refractive index of the host medium [48]. These waves may be nonuniform, i.e., they may have a nonzero transverse wave number (spatial frequency). For such waves, the wave vector $\mathbf{k}_1$ is tilted with respect to the wires under a nonzero angle $\pi/2 - \psi$, as shown in Fig. 1. The spatial spectrum of such waves is not restricted within the effective-medium model, but the microscopic analysis [49] shows that spatial frequencies higher than π/a (a is the period of the WM) cannot propagate along the wires at frequencies $\omega_1 < \omega_p$.

Following Ref. [48], a lossless WM is described by the spatially dispersive permittivity tensor with the axial component

$$\epsilon_{yy}^{(1)} = \epsilon_h \left(1 - \frac{k_p^2}{k_{10}^2 - k_{1y}^2} \right). \quad (1)$$

In Eq. (1), $\epsilon_h = n^2$ is the relative permittivity of the host medium, $k_p \equiv n\omega_p/c$ denotes the effective plasma wavenumber of the WM given by $k_p^2 = 2\pi\epsilon_h/[a^2 \ln a^2/4r_0(a - r_0)]$ [49] (where r_0 is the radius of the wire cross section), and $k_{10} = n\omega_1/c$ represents the wave number of a wave that would propagate in the host medium without wires at frequency ω_1 . The spatial dispersion is manifested in Eq. (1) by the dependence of the axial permittivity on the wave vector component y_0k_{1y} . Here and below $\mathbf{y}_0$, as well as $\mathbf{x}_0$ and $\mathbf{z}_0$ are the unit vectors of the Cartesian coordinate system). If the wires are thin enough, i.e., $r_0 \leq 0.1a$ (see, e.g., Ref. [49]), the transverse component $\epsilon_{xx}^{(1)} = \epsilon_{zz}^{(1)}$ of the permittivity tensor is practically not affected by the wires and approximately equals ϵ_h . Here, we orient the axis z so that there is no field variation along it and consider the case where the magnetic field $\mathbf{H}_1$ is polarized along z .

Expression (1) allows one to see that the low-frequency ($\omega_1 < \omega_p$) eigenmodes of the WM ($|k_{1x}| < 1/a$) are nonuniform plane waves that have the same group velocity and differ from one another only by their spatial frequencies k_{1x} . In Fig. 1(a), the spatial frequency is shown by a dashed line to stress that this value describes only the phase distribution across the energy propagation direction. Since the phase front of the TEM wave is tilted by the angle $\pi/2 - \psi$ to the plane (x - z), the vector $\mathbf{k}_1$ corresponds to the phase velocity $\omega_1/k_1 = c \sin \psi/n$. The electric and magnetic field vectors are, respectively, $\mathbf{E}_1 = E_1\mathbf{x}_0$, $\mathbf{H}_1 = H_1\mathbf{z}_0 = E_1\mathbf{z}_0/\eta$, where $\eta = \sqrt{\mu_0/\epsilon_0\epsilon_h}$ is the wave impedance of the host medium [48]. For simplicity we assume below that the host medium of our wires is air ($\epsilon_h = 1$).

Let at the moment $t = 0$, simultaneous breaks occur in every wire with a sufficiently small period $d \ll \lambda_1$, where $\lambda_1 = 2\pi c/\omega_1$ (recall that this wavelength is much larger than

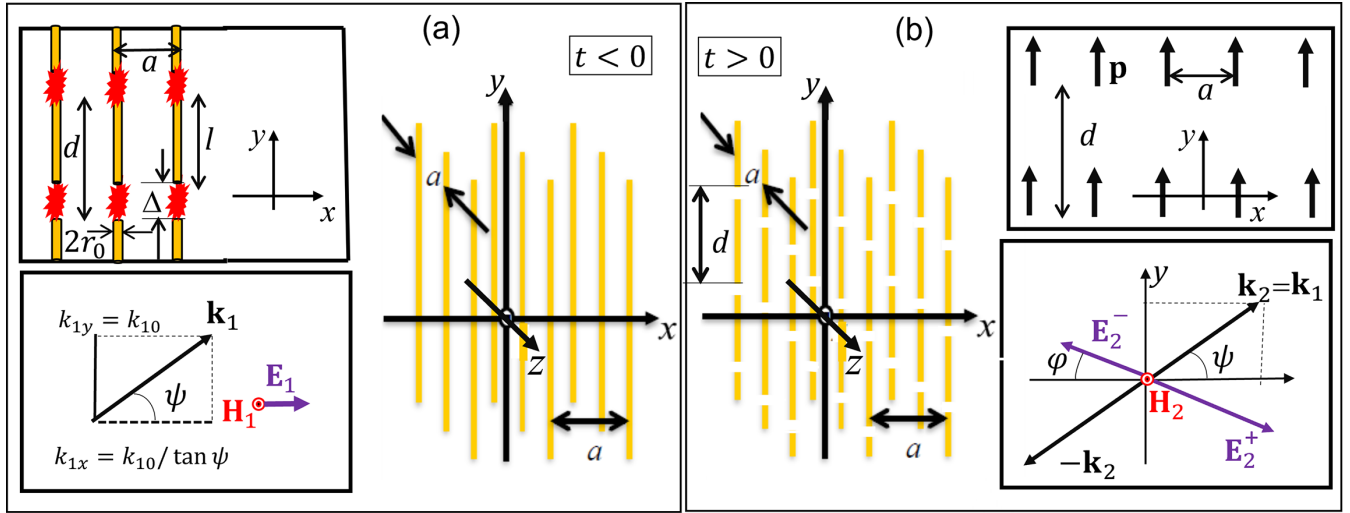


FIG. 1. (a) Wire medium of solid wires in which a non-uniform TEM wave with a nonzero spatial frequency k_{1x} propagates along y . The wave vector $\mathbf{k}_1$ describes the phase velocity tilted to the wires (bottom inset). Continuous wires can be obtained from split wires using, for example, microplasma gas discharge switches (top inset). (b) Wire medium of split wires in which the time-refracted and the time-reflected TM-polarized waves arise. This medium is a uniaxial lattice of electric dipoles (top inset). On the bottom inset the vectors of the time-refracted and time-reflected waves arising in this medium are shown.

a since $\omega_1 \ll \omega_p$). Consequently, the WM of solid wires (WM 1) transforms into a WM of split wires (WM 2). In this transition, waves at a new frequency ω_2 arise in accordance with the general theory of temporal discontinuities in materials [4]. Assume that the splits are of substantial length ($\Delta \gg r_0$, where Δ is the split width and r_0 is the wire radius) and that the period d remains subwavelength also at ω_2 . Hence, even in the case $\Delta \ll d$, WM 2 will be a regular metamaterial of short metallic rods of length $l = d - \Delta$, and its permittivity can be found as [50]:

$$\varepsilon_{yy}^{(2)} = 1 + \frac{\alpha/V}{1 - \alpha C}, \quad (2)$$

where $V = da^2$ is the lattice unit cell volume, $C = 1.2/(\pi d^3) - (8\pi/d^3)K_0(2\pi a/d)$ is the lattice interaction factor (K_0 is the McDonald function), and α is the polarizability of the rod of length l that is defined in Ref. [50] as $\alpha \equiv p/\varepsilon_0 E_{ly}$. Here, $p = p_y$ is the rod dipole moment, and E_{ly} is the y component of the local electric field at the rod center. If $a = d$, the formula for C gives $C = 1/3$, and the expression for $\varepsilon_{yy}^{(2)}$ transforms into the well-known Lorenz-Lorentz formula. In Ref. [50], α is given only for spheroidal metal particles. Formula (5.17) from Ref. [51] gives for a short thin rod $\alpha = 4l^2/(3j\omega_2\varepsilon_0 Z_{sd})$, where Z_{sd} is the complex input impedance of a short dipole [52]. The real part of this impedance corresponds to the imaginary part of α , which describes the scattering by the rod. This imaginary part must be omitted in the expression for effective permittivity. An accurate model of a uniaxial lattice [53] shows that the scattering part of the rod polarizability cancels out with the imaginary part of the interaction constant C (this imaginary part is responsible for the field retardation on the scale of a unit cell). Therefore, the effective permittivity of a lossless lattice is obviously real and in its expression we should replace the complex polarizability α by its real part α' substituting $Z_{sd} = 1/j\omega_2 C_{sd}$. Here, $C_{sd} = \pi\varepsilon_0 l/2 \ln(l/r_0)$ is the capacitance of the short dipole

(see, e.g., in Ref. [52]). In the expression for α' , the frequency ω_2 cancels out, and the expression (2) takes the form

$$\varepsilon_{yy}^{(2)} \approx 1 + \frac{\alpha'}{V(1 - \alpha'C)}, \quad (3)$$

where $\alpha' = l^3/[6 \ln(l/r_0)]$. Equation (3) means that our WM 2 is a dispersion-free uniaxial dielectric.

Let us see what happens if the WM of continuous wires transits to such a dielectric. In accordance with Ref. [4], at $t > 0$, both time refraction and time reflection must occur: two waves with electric field amplitudes $\mathbf{E}_2^\pm$ and magnetic ones $\mathbf{H}_2^\pm$ oscillating at a new frequency ω_2 will propagate in opposite directions, as shown in Fig. 1(b). The time-refracted wave inherits the previous wave vector: $\mathbf{k}_2 = \mathbf{k}_1$. Meanwhile, the tilt angle φ of the electric field vector $\mathbf{E}_2^+$ of the time-refracted wave is different from $\pi/2 - \psi$ since this wave is not TEM. The time-reflected wave propagates oppositely to the time-refracted one and also has the transverse magnetic (TM) polarization, as depicted in Fig. 1(b). To find the amplitudes of both these waves and the new frequency ω_2 , we may use the time boundary conditions for vectors $\mathbf{D}$ and $\mathbf{B}$ [4] that can be written as follows:

$$D_{1x} = \varepsilon_0 E_1 = D_{2x} = \varepsilon_0 (E_2^- - E_2^+) \cos \varphi, \quad (4)$$

$$B_1 = \mu_0 \frac{E_1}{\eta} = B_2 = \mu_0 \frac{(E_2^+ + E_2^-) \cos \varphi}{Z^{TM}}, \quad (5)$$

where the uniaxial impedance Z^{TM} relates the transverse components of the electric and magnetic fields ($E_{2x}^\pm = \pm E_2^\pm \cos \varphi = \pm Z^{TM} H_2^\pm$). The time boundary condition for the axial component of the electric displacement ($D_{1y} = D_{2y}$) is useless because $\varepsilon_{yy}^{(1)}$ is infinite, whereas the electric field component E_{1y} is null.

From Maxwell's equations for $t > 0$ with the substitution $D_{2y}/\varepsilon_{yy}^{(2)} = \varepsilon_0 (E_2^- - E_2^+) \sin \phi$, we obtain $\cot \varphi = \varepsilon_{yy}^{(2)} \tan \psi$,

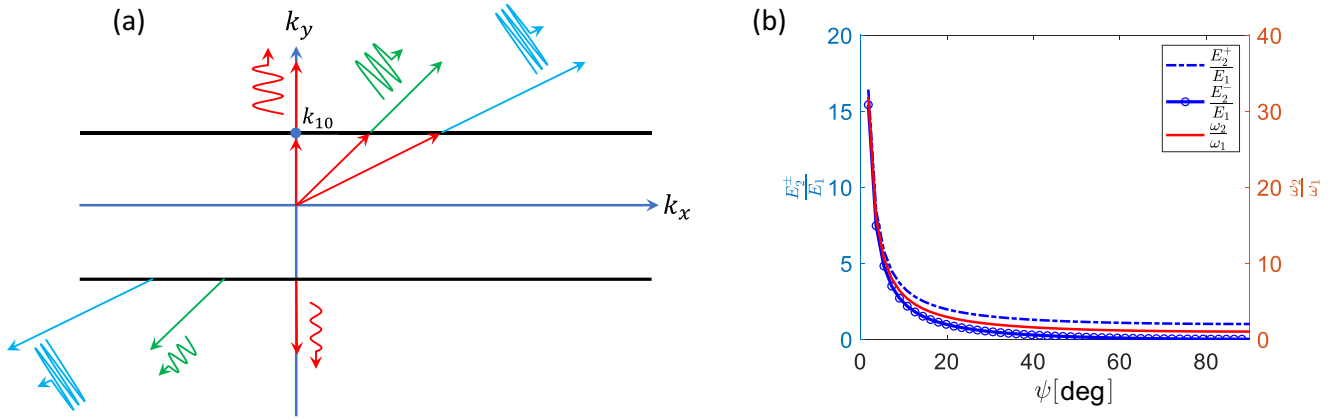


FIG. 2. (a) The isofrequency contour associated with the wire medium made of perfectly conducting solid wires. The red arrows show different wavevectors for waves propagating in this medium at a fixed angular frequency. The red, green, and blue signals show the transformed waves traveling at different angular frequencies and amplitudes in the wire medium made of split wires. (b) The amplitudes of the time-reflected (blue curve with markers) and time-refracted waves (dot-dash blue curve) and the ratio ω_2/ω_1 (red curve) versus the angle ψ .

$Z^{TM} = \eta(1 - k_{2x}^2/k_{20}^2 \varepsilon_{yy}^{(2)})^{1/2}$, and $k_{2x}E_{2y}^\pm - k_{2y}E_{2x}^\pm = \eta k_{20}H_2^\pm$, where it is denoted $k_{20} \equiv \omega_2/c$. Using the continuity of the axial and transverse wave numbers $k_{2y} = k_{1y} = k_{10}$, $k_{2x} = k_{1x} \equiv k_{10} \cot \psi$, we can write the dispersion equation of the disconnected WM in the form $k_{2x}^2/\varepsilon_{yy}^{(2)} + k_{2y}^2 = k_{10}^2(1 + \cot^2 \psi/\varepsilon_{yy}^{(2)}) = k_{20}^2$, which results in the new frequency (recall that $k_{10} = \omega_1/c$ and $k_{20} = \omega_2/c$)

$$\omega_2 = \omega_1 / \cos \varphi = \omega_1 \sqrt{1 + \cot^2 \psi / \varepsilon_{yy}^{(2)}}. \quad (6)$$

In Eq. (6) for $\varepsilon_{yy}^{(2)}$, we can employ the relation expressed in Eq. (3). The wave impedance in WM 2 takes the form $Z^{TM} = \eta / \sqrt{1 + \cot^2 \psi / \varepsilon_{yy}^{(2)}}$. Using this expression with relations (4) and (5), we obtain after some algebra the amplitudes of both time-refracted and time-reflected waves:

$$E_2^\pm = \pm \frac{E_1}{2} \left(1 \pm \frac{\omega_2}{\omega_1} \right). \quad (7)$$

In the case of ultimately high spatial frequencies $k_{1x} \gg k_{10}$, i.e., when $\cot \psi$ approaches $(k_{10}a)^{-1}$, the initial wave is maximally affected by the transformation of WM 1 into WM 2. In this case, both time-refracted and time-reflected waves propagate almost along x , and (7) gives their approximate equivalence: $E_2^- \approx E_2^+ \gg E_1$. The new frequency ω_2 is very high and nearly equal to $\omega_1[\varepsilon_{yy}^{(2)}]^{-1/2}(k_1a)^{-1}$. If the spatial frequency is zero, i.e. the initial wave is a uniform plane wave propagating along the wires, it cannot be affected by this temporal discontinuity. Equations (6) and (7) give the same amplitude E_1 and frequency ω_1 for the time-refracted wave, whereas no time-reflected wave is excited. All theoretical expectations outlined above are illustrated in a schematic representation in Fig. 2(a) and supported by numerical calculations presented in Fig. 2(b). For the latter, we used typical parameters of a wire medium operating in the millimeter wave range, namely $f_1 \equiv \omega_1/2\pi = 100$ GHz, $r_0 = 30$ μm , $a = 10r_0$. For splits, we adopted $d = 2a$ (period) and $\Delta = a/4$ (width). These parameters correspond to the axial permittivity of the split wire medium slightly different from unity $\varepsilon_{yy}^{(2)} = 1.03$. Furthermore, the polarization in the direction orthogonal to the thin wires is negligibly small.

Therefore, in our numerical modeling, replacing this uniaxial material by free space does not noticeably change the result.

III. TEMPORAL JUMPS IN INDUCTIVELY LOADED WIRE MEDIA

Now let us discuss how to transform a continuous WM into a cut-wire lattice in practice. Such transformation between arrays of continuous wires and arrays of disconnected pieces of wires can be done using some kind of switches periodically inserted into the wires. To account for finite impedances of the switches, we use a model of a periodically loaded WM that was developed in Ref. [54]. If the distance between insertions is electrically small, periodical loading of the wires with the lumped impedance Z_{load} and the subwavelength period d can be treated as uniform loading. Then the wire impedance per unit length (p.u.l.) $Z_w = Z_{\text{load}}/d$ should be added to the wire resistance p.u.l. Then, the axial permittivity of WM 2 becomes [54]:

$$\varepsilon_{yy}^{(2)} = 1 - \frac{k_p^2}{k_{10}^2 - j\xi k_{10} - k_{1y}^2}, \quad (8)$$

where (in the case of perfectly electrically conducting rods) the factor ξ arises due to the loads:

$$\xi = \frac{\sqrt{\mu_0 \varepsilon_0} Z_w}{L_{WM}}, \quad L_{WM} = \frac{\mu_0}{2\pi} \ln \left[\frac{a^2}{4r_0(a - r_0)} \right]. \quad (9)$$

Here, L_{WM} is the collective inductance p.u.l. of the unloaded WM (that is, WM 1). At a first glance, for splitting wires in WM 1 one may use such electronic switches as pin-diodes, varactors, or MEMS periodically inserted in small gaps in the wires. In the closed state they usually have some parasitic capacitive impedance Z_{load} , which can be negligibly small compared to $\omega_1 L_{WM}$, and in the open state they have a large Z_{load} . However, our analysis has shown that the open-state impedance is not large enough. In the switched-off state, all these switches have a non-negligible capacitance, so that the ends of the cut pieces of the wires are coupled much stronger than the via free-space gaps.

Therefore, in this paper we propose another approach. Let our WM 1 be loaded by inductive loads of finite length. If we remove these connecting loads, the WM 2 will be the needed array of dipole scatterers, as there will be enough large gaps g between the wire sections. The inductance of these connectors will add to the effective inductance of WM 1, but qualitatively will change nothing. A suitable tool that can enable this approach is gas discharge. A usual discharge such as the static arc is difficult to use in practice as it requires high voltages applied along the wires. However, an electromagnetically assisted discharge over a tiny gap utilized in some nanotechnologies [55] appears to be suitable. If the wire has the thickness of a few μm and the gap Δ has the width in the interval $\Delta = 10\text{--}100\mu\text{m}$, glow microplasma is formed in the gap under practically realistic voltages (dozens of V) and realistic microwave pumping (few mW per gap) [55]. In this scenario, every wire of WM 1 should be an array of metal rods whose ends are connected by microplasma channels, as sketched in the inset of Fig. 1(a).

The impedance properties of the microplasma in such channels were studied in Ref. [56] for sub-GHz frequencies. In Ref. [57], the impedance of a microplasma capacitor with a large (1 mm) gap between the electrodes was studied in the range 0.3–20 GHz. Both these works showed that the channel resistance per unit length drops versus frequency, and at microwaves it is negligibly small compared to the channel inductive impedance. Since the last one linearly grows with frequency up to 20 GHz, one may assume that the channel inductance weakly depends on the frequency in the mm-wave and THz ranges. In our estimations, the microplasma discharge of length $20\mu\text{m}$ has the inductive impedance of the order of $|Z_{\text{load}}| = \omega L_{\text{channel}} \sim 2\text{ Ohm}$ at 100 GHz.

In accordance with Ref. [54], we have $Z_{\text{load}} = j\omega_1 L_{\text{load}}$ and the impedance per unit length $Z_w = j\omega_1 L_{\text{load}}/d$ arises in WM 1, whose axial permittivity becomes equal to

$$\varepsilon_{yy}^{(1)} = 1 - \frac{k_p^2}{k_{10}^2 - j\xi_1 k_{10} - k_{1y}^2}, \quad (10)$$

where the factor ξ_1 is equal to:

$$\xi_1 = j \frac{k_{10} L_{\text{load}}}{d L_{WM}}. \quad (11)$$

Substitution of (11) into (10) transforms this relation to

$$\varepsilon_{yy}^{(1)} = 1 - \frac{k_p^2}{k_{10}^2(1 + \beta) - k_{1y}^2}, \quad \beta \equiv \frac{L_{\text{load}}}{d L_{WM}}. \quad (12)$$

In this model, the eigenmodes of WM 1 are quasi-TEM waves with the wave vector $y_0 k_{1y}$, whose absolute value is equal to $k_{1y} = k_1 \sqrt{1 + \beta}$. It makes the derivations more cumbersome, but the results (6) and (7) remain valid with a minor change: the multiplicative factor $\sqrt{1 + \beta}$ arises in their right-hand sides. This constant factor changes qualitatively nothing in the curves depicted in Fig. 2(b).

IV. APPLICATION OF THE REVEALED EFFECT FOR SUBWAVELENGTH IMAGING

Let us discuss a useful implication of the revealed phenomenon. Here we show a way to an efficient conversion of

the waves with high spatial frequencies into usual propagating plane waves after the time jump in a spatially uniform wire medium.

Following Eq. (6), the waves propagating in WM 1 with large spatial frequencies corresponding to the condition $k_{1x}^2 \gg k_{10}^2 \varepsilon_{yy}^{(2)}$ transform into plane waves propagating in WM 2 with the frequencies $\omega_2 \approx (c/\sqrt{\varepsilon_{yy}^{(2)}})k_{1x}$. So, new frequencies ω_2 of waves that arise in WM 2 are not only uniformly proportional to spatial frequencies k_{1x} but, for large values of ψ correspond to new frequencies $\omega_2 \sim k_{1x}$. Imagine that a complex source (also called object) with strongly subwavelength details produces electromagnetic fields in front of the WM 1 sample. These subwavelength details will produce subwavelength oscillations in waves propagating along the wires of WM 1, and high spatial frequencies $k_{1x} > k_{10}$ will carry out these oscillations so that the spatial spectrum of the propagating pulse repeats the spatial spectrum of the object field.

The scenario of the temporal discontinuity that we have described here is an alternative approach to subwavelength imaging enabled by a WM endoscope. In our approach, the spatially dispersive WM disappears, when it is replaced by a split WM, which is simply a uniaxial dielectric. The result of this temporal discontinuity is an angular pattern of propagating waves, where the frequencies and intensities encode the spatial structure of the object. As mentioned in the introduction, the object image can be restored from a spectroscopic measurement of the angular pattern produced by the WM sample after the described temporal jump.

To restore an arbitrary wave package propagating in WM 1 at a frequency ω_1 it is enough to retrieve the amplitudes E_1 corresponding to the angles ψ under which the time-refracted and time-reflected waves propagate. First, from (7) we see that $E_1 = E_2^+ - E_2^-$. It means that we can find $|E_1|$ from measured amplitudes of emitted time-refracted and time-reflected waves. Importantly, measuring the phases of time-refracted and time-reflected waves is not necessary. Really, for a given angle ψ after measuring at the corresponding frequency ω_2 we may find $|E_1|$ via the intensity either of the time-refracted (I_{2+}) wave or time-reflected (I_{2-}) wave. It is so because the relations (7) are real-valued. Using (7) one can find both the amplitude and phase of E_1 , if we know the permittivity $\varepsilon_{yy}^{(2)}$. This real-valued (in the lossless approximation) permittivity can be found, e.g., detecting the polarization axis of the electric field E_2^+ or E_2^- (that will deliver the angle ϕ) and using the relation $\varepsilon_{yy}^{(2)} = 1/\tan\phi \tan\psi$ to find the axial permittivity of the split WM. Thus, we will be able to restore the complex spatial spectrum of the initial spatial wave package, which propagated in WM 1 via standard spectroscopic measurements of the intensity angular distribution in the plane (x - y).

V. BEYOND MORGENTHALER'S TIME BOUNDARY CONDITIONS

In general, Morgenthaler's time boundary conditions introduced in Ref. [4] to describe effectively instantaneous transitions between two nondispersive media may be not applicable for media with spatial dispersion. Moreover, the boundary conditions at temporal interfaces depend on the actual mechanism of parameter switching [58]. In the above

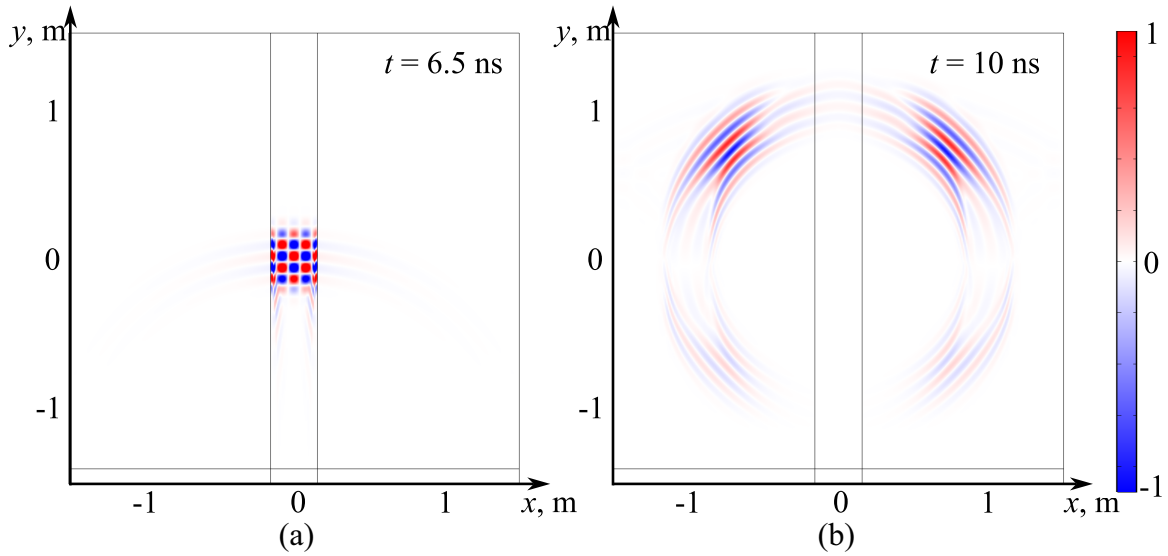


FIG. 3. Simulation results for $k_x = 40 \text{ m}^{-1}$, $k_y = 40 \text{ m}^{-1}$. (a) shows H_z normalized to its maximal value H_0 before the temporal interface (at the moment $t = 6.5 \text{ ns}$). Tails of the guided mode outside of the slab are clearly seen. (b) shows H_z also normalized to H_0 at the time moment $t = 10 \text{ ns}$, after the temporal interface. Two symmetric time-refracted and two symmetric time-reflected waves propagating away from the time-varying sample are clearly seen.

solution, Morgenthaler's conditions are applicable to spatially dispersive media because in our special case [when the electric field is polarized in the $(x-y)$ plane, the same plane in which the phase of the wave package propagating in WM 1 varies] for each value of k_x at a given frequency ω_1 , waves at only one new frequency ω_2 are created. If the electric field is polarized along z , the same temporal interface may result in creation of waves at two or more new frequencies per one spatial frequency. The same refers to the inverse process when WM 2 transits to WM 1.

Furthermore, we need to discuss if the boundary conditions are properly modeling the suggested physical mechanism of the temporal discontinuity. Whereas (4) appears to be valid because the transversal polarization of both WM 1 and WM 2 is negligibly small, the conservation of the magnetic flux across the time interface (condition $B_1 = B_2$) is not so evident. In the effective-medium model, this flux is continuous in time, otherwise the singularity of the electric field arises at $t = 0$. However, (as noticed in work [58]) temporal interface makes sense only as a model of a real transition process. Though its duration is short compared to the field oscillation period, it is finite, and during this process the effective-medium model may fail even for a dense array of wires. Paper [58] did not consider the case of spatially dispersive WM, but introduced several transition models in frequency dispersive media, including the Drude model of plasma. Since wire media for waves propagating across the wires exhibits frequency dispersion of the Drude type, it is reasonable to consider our temporal interface using the approach of Ref. [58]. In that work, the transition process in a medium with the Drude temporal dispersion is modeled as switching in an effective transmission line. The case when the magnetic flux is preserved is schematically shown in Fig. 3(d) of that work. In this case, the inductance of the effective transmission line drops, but the magnetic flux is preserved. For our case it means that Morgenthaler's conditions hold in spite of the

drop of the effective inductance of WM 1—in the end of the transition process the effective inductance of the split WM is equal $L_{\text{SWM}} = L_{\text{WM}}/\alpha$. The coefficient can be easily found as $\alpha = d/(d - \Delta)$ because the mutation of the microplasma discharges in all wires means that every wire piece of length d shortens by Δ . In spite of this inductance drop, the magnetic flux in the effective transmission line is preserved due to an instantaneous injection of electromagnetic energy into the system. This injection is modeled in Ref. [58] by a δ -function source. Physically, this source describes the external forces that produce some work in the system. Thus, Morgenthaler's time boundary conditions in our case describe the active transition scenario in which the discharges in the microgaps of every wire are muted by external forces working against the initial electromagnetic field.

An alternative model of the temporal jump between WM 1 and WM 2 is the passive transition illustrated by Fig. 3(c) of Ref. [58]. In this model, the currents in the remaining parts of the wires are preserved, and the drop of the effective-medium inductance results in the same drop of the magnetic flux. Then we have $B_1 = \alpha B_2$, i.e., the factor α arises in the right-hand side of (5). Repeating the same derivation steps, we obtain the following result for the passive transition from a WM whose wires are connected by microplasma discharges to a split WM:

$$\begin{aligned} \omega_2 &= \omega_1 \sqrt{1 + \cot^2 \psi / \varepsilon_{yy}^{(2)}} \sqrt{1 + \beta}, \\ E_2^\pm &= \pm \frac{E_1}{2} \left(1 \pm \frac{\omega_2}{\alpha \omega_1} \right). \end{aligned} \quad (13)$$

The choice in favor of one of two scenarios, the active transition or the passive one, requires additional studies of the microplasma discharge. However, for the suggested method of imaging the difference between two scenarios is not important. Since $\beta \ll 1$ and $\alpha \approx 1$ the correction terms in (13)

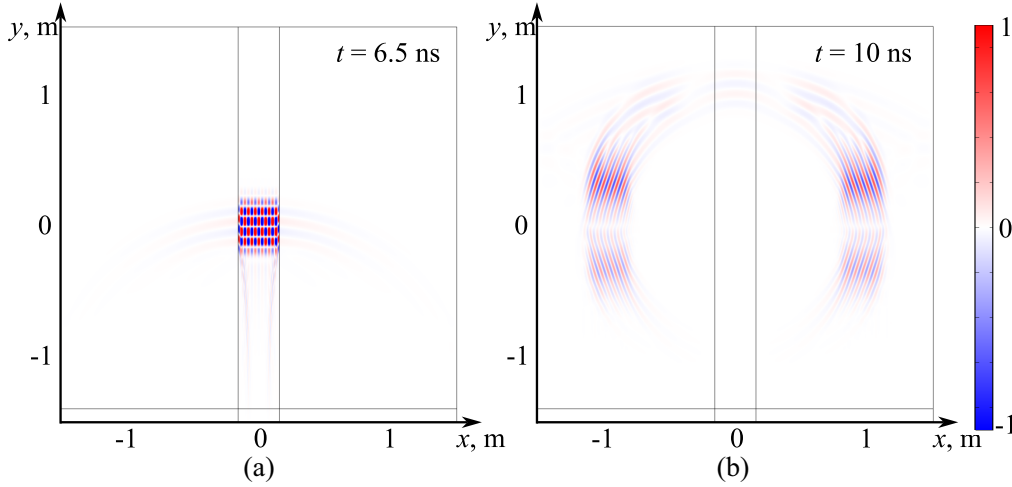


FIG. 4. Simulation results for $k_x = 120 \text{ m}^{-1}$, $k_y = 40 \text{ m}^{-1}$. (a) shows normalized H_z at $t = 6.5 \text{ ns}$. (b) shows normalized H_z at $t = 10 \text{ ns}$. Comparison with Fig. 3 demonstrates the dependence of the propagation direction of the time-refracted waves and the transverse wave numbers of the initial wave.

change qualitatively nothing compared to the formulas (6) and (7).

VI. IMPACT OF ENDOSCOPE FINITE WIDTH AND NUMERICAL RESULTS

The applicability of the model of an infinite effective medium to a finite-width endoscope is fully justified even if this medium is granular and spatially dispersive. In existing WM endoscopes, which found applications from MRI (radio frequencies) to optical sensing (infrared frequencies), the model of WM correctly describes the polarization properties of the endoscope starting from several (5–10) unit cells across the cable width (see, e.g., in Refs. [35,37–39,42,43]). The impact of the finite width may be harmful due to partial internal reflection of the final pulse if the split wires are embedded in a dielectric host. However, this effect can be practically nullified with the use of so-called gradient-index antireflecting coatings adopted to dielectric fibers [59,60]. The only possible impact of the finite width that may disable our claims is the validity of our model of temporal interface in infinite media for a finite-width WM sample. This issue can be studied using full-wave numerical simulations.

Simulation of a temporal interface in a finite-width WM endoscope is not trivial. An array of metal wires whose parts of finite length δ disappear during a short transition process can hardly be simulated, as well as the sample of the homogenized WM 1 due to its spatial dispersion. However, our main conclusions can be validated. We can use the fact that a slab of a uniaxial dielectric with huge axial and small transverse components of the permittivity tensor supports nonuniform quasi-TEM waves are eigenmodes. The isofrequency contour in such media is an extremely stretched ellipsoid, which is a good approximation of the flat contour of the ideal wire media for moderate values of the transverse wavenumbers. Therefore, our switchable WM endoscope can be emulated by a switchable fiber (3D) or a slab (2D) with sufficiently high axial permittivity (e.g., 1000), whereas the transverse one is equal to that of the ambient. At a certain moment, the huge

axial permittivity drops and becomes equal to that given by formula (3), i.e., $\varepsilon_{yy}^{(2)} = 1.03$. Then a similar conversion of propagating waves in the uniaxial dielectric sample into space waves should occur. Details of the simulation setup and videos of the transition process are reported in the Supplemental Material [61]. These simulations confirmed the validity of our model for endoscopes whose thickness is comparable to λ .

In Fig. 3 a screenshot of the video file available in Ref. [61] is shown. In this case, the spatial spectrum of the quasi-TEM wave package propagating in the ultimately anisotropic slab is formed by two spatial frequencies $k_x = \pm k_y$ (k_y is close to ω_1/c). At $t = 6.5 \text{ ns}$ the slab axial permittivity ε_{yy} changes from $\varepsilon_{yy}^{(1)} = 1000$ to $\varepsilon_{yy}^{(2)} = 1.03$, and, as predicted by the theory in the above sections, the initial wave package transforms into two time-refracted and two time-reflected wave pulses. The propagation directions of these pulses and their carrier frequency practically coincide with the predictions of our theory (the difference is in the fifth character). For the amplitudes of the pulses, the deviation from our theory is about 1.8%, which confirms the validity of the developed model. A similar screenshot, shown in Fig. 4 corresponds to the case when $k_x = \pm 3k_y$. Again, the new frequency and the propagation directions of the created pulses match the predictions of our theory almost perfectly. The disagreement for the amplitudes is noticeable but also small (within 2.5%). Actually, this estimation of disagreement serves also as an estimation of the image restoration uncertainty due to the finite size of the endoscope and the finite duration of the initial pulse. These numerical data simulate a measurement, while the analytical values correspond to the exact initial field distribution. More details about the simulation setup can be found in Ref. [61].

Our simulations show that our model of temporal interfaces developed for infinite WM is applicable for finite-size samples. Note that the finite length of the sample should have no impact because at the signal frequency the endoscope is matched. At the new frequencies ω_2 , the endoscope is mismatched but the matching at these frequencies is not needed because the new waves propagate obliquely to the waveguide axis and leak from the endoscope through its shell. Nonideal

matching of the endoscope at the signal frequency and electromagnetic losses in it may worsen the resolution, but cannot harm the conversion of waves at spatial frequencies k_x into waves at new frequencies ω_2 . This conversion results from the jump of the system properties governed by the conservation of the wave momentum.

VII. DISCUSSION AND CONCLUSIONS

In this article, a temporal interface in a medium with strong spatial dispersion is considered, at an example of wire media. WM whose properties are modulated in time are studied in papers [62,63]. However, in Ref. [62], the wires are capacitively loaded, allowing the low-frequency wave to propagate orthogonally to the wires. For such waves, spatial dispersion is absent. In Ref. [63], the plasma frequency of the WM resonator is modulated so slowly that the medium permittivity practically does not change during the oscillation period. In that case, spatial dispersion does not matter for temporal variations of the resonance frequencies. Meanwhile, in the present work, we suggest switching a WM (so that it becomes a uniaxial dielectric) in which the propagating eigenmode has an obliquely directed wave vector. In this case, strong spatial dispersion is present and makes the axial component of the medium permittivity tensor infinite for TEM modes. This is the reason why the spatial spectrum of the propagating wave package is unambiguously encoded in the frequency spectrum of the resulting pulse. Importantly, the proposed method allows restoration of subwavelength field distributions including both amplitude and phase by measuring only far-field amplitude and frequency distributions. That is, it is possible to, obtain the phase information of the original field without the need to do phase measurements. We hope that this new effect can be useful as it opens a new gate for subwavelength imaging.

Analyzing the ways of practical implementation of such temporal interfaces we came to the idea of WM 1 that is initially inductively loaded by gas discharges through micron gaps. A uniaxial dielectric arises from it when all discharges are simultaneously muted. We considered two

possible scenarios of such temporal discontinuity—when the Morgenthaler time boundary conditions hold and when they need to be modified. In both these scenarios, we theoretically obtain the same conversion of the initial nonuniform quasi-TEM wave propagating along the wires into a TM-polarized plane wave propagating obliquely to the wires. Moreover, parameters of this resulting wave are basically the same in both scenarios.

In the effective-medium model, the correspondence between the frequency spectrum of the resulting TM pulse forming by such waves and the spatial spectrum of the initial package of nonuniform quasi-TEM waves is exact. It means that the image (which results from the postprocessing of spectroscopic measurements of the resulting pulse angular intensity distribution) in the frameworks of this model is perfect. In practice, the spatial resolution will be restricted by the granularity of both initial and split wire media at the level of their period a . Due to this granularity, both effective-medium models become not valid at the spatial frequencies higher than k_p and at the frequencies higher than ω_p . It will result in a restriction of the relevant frequency spectrum of the resulting broadband pulse by a frequency that is close to the plasma frequency. Due to this restriction, the minimal achievable resolution in our method will be limited, though hardly much stronger than by the scale of the granularity a , which is deeply subwavelength. Of course, the impact of conduction losses and measurement errors will also harm the restored image. In the current stage of research we cannot accurately estimate the effect of all these harmful factors.

We hope that this study will motivate further investigations of temporal discontinuities in spatially dispersive media. The revealed opportunity of wireless superresolution imaging is only the first of potential applications, which may result from temporal interfaces in media with strong spatial dispersion.

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