Modal Propagation and Excitation on a Wire-Medium Slab

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Abstract—A grounded wire-medium slab has recently been shown to support leaky modes with azimuthally independent propagation wavenumbers capable of radiating directive omnidirectional beams. In this paper, the analysis is generalized to wire-medium slabs in air, extending the omnidirectionality properties to even modes, performing a parametric analysis of leaky modes by varying the geometrical parameters of the wire-medium lattice, and showing that, in the long-wavelength regime, surface modes cannot be excited at the interface between the air and wire medium. The electric field excited at the air/wire-medium interface by a horizontal electric dipole parallel to the wires is also studied by deriving the relevant Green’s function for the homogenized slab model. When the near field is dominated by a leaky mode, it is found to be azimuthally independent and almost perfectly linearly polarized. This result, which has not been previously observed in any other leaky-wave structure for a single leaky mode, is validated through full-wave moment-method simulations of an actual wire-medium slab with a finite size.

Index Terms—Leaky modes, metamaterials, near field, planar slab, wire medium.

I. INTRODUCTION

The 1-D wire medium consists of a periodic arrangement of thin perfectly conducting cylinders (wires), infinitely long and parallel, embedded inside a homogeneous host dielectric medium. Such an artificial medium, also known as a rodded medium, has been known since the 1950s to be described in the long-wavelength regime by a scalar permittivity with a plasma-like dispersion behavior for waves having the electric field polarized along the axis of the wires [1].

As can easily be predicted from the directionality of the wire lattice, the 1-D wire medium is uniaxial for waves with arbitrary polarization. A far less trivial property, first pointed out in [2] and [3], is the presence of spatial dispersion even in the large wavelength regime, i.e., the nonlocal nature of the electrical response of the wire medium. This feature, also present in more complex 2-D and 3-D wire arrangements [4], [5], has remarkable consequences on the propagation of plane waves inside an unbounded 1-D wire medium: for instance, in addition to the usual ordinary and extraordinary waves existing in any uniaxial medium (TE and TM with respect to the wire direction, respectively), an additional TEM wave may exist, propagating in the direction of the wires; furthermore, for extraordinary plane waves, the propagation wavenumber does not depend on the direction of propagation of the wave, i.e., the isofrequency surfaces in the wavenumber space are spherical [3]. The presence of spatial dispersion poses a formidable challenge in the solution of boundary-value problems with nonplanar boundaries. Even in the presence of purely planar interfaces, additional boundary conditions may be required when the wires are not purely parallel to the interfaces [6]. Propagation of waves in bounded wire-medium structures has been considered in [7] and [8].

In this paper, we consider the problem of modal excitation and propagation along a canonical planar waveguide, a slab in free space, made of a 1-D wire medium with wires parallel to the air–slab interfaces (see Fig. 1, where the relevant geometrical parameters and the adopted coordinate system are also shown). Recently, certain properties of a grounded configuration of such a slab have been studied, excited by a simple electric dipole source placed inside the wire-medium and oriented parallel to the wires [9]. For frequencies slightly above the plasma frequency, where the relative permittivity is positive but small, it has been shown that the structure is capable of producing narrow-beam radiation patterns. Depending on the frequency, the beam may be either a pencil beam at broadside, or a conical beam at a particular scan angle. It was shown in [10] and [11] that this directive radiation is due to the excitation of a leaky...
mode supported by the grounded slab having a wavenumber independent of its azimuthal angle of propagation.

The aim of this paper is twofold. First, to generalize the analysis of modal propagation at arbitrary angles to the case of a wire-medium slab in air, considering both even and odd modes, performing a parametric analysis of the properties of the leaky modes, and showing that surface (either ordinary or plasmon-like) modes cannot propagate along the slab in the long-wavelength regime. Second, to consider the excitation of a dominant leaky mode by a horizontal dipole source parallel to the wires, and to study the properties of the resulting near field at the air–slab interface, adopting both a homogeneous model and a full-wave method of moments (MoM) approach.

This paper is organized as follows. In Section II, the homogenized model of the wire-medium slab is presented along with the relevant transverse equivalent networks; these allow for a discussion of the properties of surface and leaky modes and of the electric field excited at the air–slab interface by a horizontal electric dipole. In Section III, the adopted full-wave modal analysis based on the MoM in the spatial domain is described. In Section IV, numerical results for the propagation of leaky modes and for the features of the electric field excited by the dipole source are presented. Finally, in Section V, conclusions are drawn.

II. HOMOGENIZED MODEL AND TRANSVERSE EQUIVALENT NETWORK

A. Homogenized Medium

The wire medium shown in Fig. 1 is anisotropic for electromagnetic waves with an arbitrary polarization, requiring the use of a uniaxial permittivity dyadic with optical axis directed parallel to the wire (y) axis when it is homogenized at wavelengths sufficiently larger than the wire spacing. Actually, as shown in [2] and [3], not only anisotropy, but also spatial dispersion needs to be included in a homogenized description of the medium, even for large wavelengths. The effective permittivity dyadic thus depends on both the frequency and wavenumber, and for thin wires [i.e., \( a < \min (d_x, d_y) \)] and in the large-wavelength limit (i.e., \( \max (d_x, d_y) \ll \lambda_0 \), with \( \lambda_0 \) the free-space wavelength), it reads

\[
\varepsilon = \varepsilon_0 \left[ \varepsilon_{xy} \mathbf{u}_x \mathbf{u}_y + \varepsilon_{rl} \left( \mathbf{u}_x \mathbf{u}_x + \mathbf{u}_z \mathbf{u}_z \right) \right]
\]

\[
= \varepsilon_0 \varepsilon_{rl} \left[ \left( 1 - \frac{k^2}{\varepsilon_{rl} k_0^2} \right) \mathbf{u}_y \mathbf{u}_y + \mathbf{u}_x \mathbf{u}_x + \mathbf{u}_z \mathbf{u}_z \right]
\]

where \( \varepsilon_{rl} \) is the relative permittivity of the medium hosting the wires (for the structures considered here, air is assumed as the host medium, i.e., \( \varepsilon_{rl} = \varepsilon_{rt} = 1 \)). \( k_0 \) is the free-space wavenumber, \( k_p \) is the plasma wavenumber, and \( k_y \) is the wavenumber along the wire axis. Here and in the following, boldface symbols denote vectors, while \( \mathbf{u}_x, \mathbf{u}_y \), and \( \mathbf{u}_z \) denote unit vectors along the \( x-, y-, \) and \( z- \) directions, respectively. The plasma wavenumber is given by (see [12])

\[
k_p = \sqrt{\frac{2\pi}{s}} \frac{1}{\sqrt{\ln \frac{s}{2\pi a} + F(\nu)}}
\]

where \( s = \sqrt{d_x d_z}, \nu = d_x/d_z, \) and

\[
F(\nu) = \frac{1}{2} \ln(\nu) + \sum_{n=1}^{\infty} \frac{1}{n} \left[ \cosh(\pi n \nu) - 1 \right] + \frac{\pi \nu}{6}.
\]

In particular, for square lattices, i.e., when \( d_x = d_y (\nu = 1) \), there results \( F(1) \approx 0.527 \), while when \( d_x = 2d_z \) or \( d_x = d_z/2 \), there results \( F(2) = F(0.5) \approx 0.7 \).

B. Modal Dispersion Equations and Surface-Wave Suppression

The field of a leaky mode propagating on a wire-medium slab, as in Fig. 1, with complex wavenumber \( k_{pW} = \beta - j\alpha \) at an angle \( \phi \) has an exponential dependence on the \( x- \) and \( y- \) coordinates with wavenumbers \( k_x = k_{pW} \cos \phi \) and \( k_y = k_{pW} \sin \phi \). The electromagnetic field inside the air region and the wire-medium slab is decomposed into \( \text{TE}_y \) and \( \text{TM}_y \) components, which can be represented through equivalent transmission lines along \( z \) (details are given in the Appendix).

The \( \text{TE}_y \) (ordinary) polarization sees an isotropic equivalent medium equal to the host medium (air). Furthermore, as shown in [10] and [11], at the air/wire-medium interface, no \( \text{TE}_{xy}/\text{TM}_x \) coupling occurs when the host medium is air. Therefore, \( \text{TE}_y \) and \( \text{TM}_y \) polarizations are decoupled and \( \text{TE}_y \) modal solutions are not supported by the wire-medium slab. (For completeness, we note here that nonmodal \( \text{TE}_y \) solutions exist, which are actually \( \text{TE}_{xy} \) plane waves propagating in free space at arbitrary angles \( \phi \) in the \( xy \)-plane with no \( z \) variation).

The study of \( \text{TM}_y \) modes in a slab of thickness \( 2h \) is performed using the transverse equivalent networks shown in Fig. 2 (a) and (b) for odd and even modes; by symmetry, the plane \( z = 0 \) is equivalent to a perfect electric conductor for odd
modes and to a perfect magnetic conductor for even modes. The transverse equivalent network parameters are

\[
k_{0}^{TM} = \sqrt{k_{0}^2 - k_{p}^2}
\]

and

\[
Z_{0}^{TM} = \frac{\gamma_{0}}{k_{0}} \frac{k_{0}^2 - k_{p}^2 \sin^2 \phi}{k_{0}^{TM}}
\]

\[
Z_{1}^{TM} = \frac{\gamma_{0}}{k_{0}} \frac{k_{0}^2 - k_{p}^2 \sin^2 \phi}{k_{0}^{TM}}
\]

where \( \gamma_{0} \) is the free-space characteristic impedance. As concerns the value of the parameter \( h \), i.e., the equivalent thickness of the homogenized slab, we have adopted here the choice \( 2h = Nd_{f} \), where \( N \) is the number of wire layers along the \( z \)-direction; although quite natural [13], this choice has no rigorous justification, but it will be seen in Section IV to yield results in good agreement with the full-wave simulations.

The modal dispersion equations for the extraordinary even and odd \( TM_{y} \) modes can be obtained by enforcing the condition of resonance on the transverse equivalent networks shown in Fig. 2(a) and (b), i.e.,

\[
\frac{1}{Z_{0}^{TM}} + j \frac{1}{Z_{1}^{TM}} \tan \left( k_{0}^{TM}h \right) = 0 \quad \text{(even modes)}
\]

\[
\frac{1}{Z_{0}^{TM}} + j \frac{1}{Z_{1}^{TM}} \tan \left( k_{0}^{TM}h \right) = 0 \quad \text{(odd modes)}.
\]

From (6) and (5), it can be seen that a possible solution is \( k_{y} = k_{0}/\sin \phi \), which corresponds to having \( k_{y} = k_{0} \). The case \( k_{y} = k_{0} \) corresponds to \( TM_{y} \) waves, and represents a special case for which the permittivity model in (1) breaks down for a practical medium since this equation predicts that \( \varepsilon_{yy} = \infty \) for such waves (for an idealized medium that obeys (1), a continuous nonmodal spectrum of \( TM_{y} \) waves may be shown to exist [3]). For \( k_{y} \neq k_{0} \), from (4) and (5), the dispersion equations (6) can be written as

\[
k_{0}^{TM} - jk_{0}^{TM} \cot \left( k_{0}^{TM}h \right) = 0 \quad \text{(even modes)}
\]

\[
k_{1}^{TM} + jk_{1}^{TM} \tan \left( k_{0}^{TM}h \right) = 0 \quad \text{(odd modes)}.
\]

which turn out to be independent of the propagation angle \( \phi \) (and equal to the dispersion equation for \( TE_{z} \) modes of an ordinary isotropic slab with \( \varepsilon_{z} = 1 - k_{z}^2/k_{0}^2 \)). This interesting property arises from the fact that the \( \phi \) dependence in both of the characteristic impedance terms of (5) are the same, and hence, this dependence cancels out when these terms are inserted into the transverse resonance equation.

From the omnidirectionality of modal propagation, important conclusions can be drawn about the existence of bound waves (surface waves) on the wire-medium slab with air as the host medium. Let us consider first propagation in the direction orthogonal to the wire axis. By symmetry, the electromagnetic field of \( TM_{y} \) modes propagating in this direction has nonzero components \( E_{y}, H_{x}, \) and \( H_{z} \) (so that these modes can also be classified as \( TE_{z} \)); from (1), it is then found that such modes see the homogenized medium as an isotropic medium with a plasma-like scalar relative permittivity \( \varepsilon_{r} = 1 - k_{y}^2/k_{0}^2 \). Now, when \( k_{0} < k_{y} \) (i.e., the frequency is lower than the plasma frequency), it follows that \( \varepsilon_{r} < 0 \), i.e., the medium is epsilon negative; therefore, \( TE_{z} \) (i.e., \( TM_{y} \)) surface waves cannot propagate (a proof may be found in [14] for odd modes with the extension to even modes being straightforward). On the other hand, when \( k_{0} > k_{y} \) (i.e., the frequency is higher than the plasma frequency), the medium has a relative permittivity that is positive, but lower than one, so that total reflection cannot occur at the slab–air interface, and again, no surface waves may exist.

Therefore, as long as the homogenization process is valid, surface waves cannot propagate orthogonally to the wires. By virtue of the omnidirectionality of modal propagation established above, it can be concluded that surface waves cannot propagate at any frequency and at any angle along a wire-medium slab in air. The only type of guided mode that may, therefore, exist on the homogenized wire medium slab is a \( TM_{y} \) leaky wave. The properties of such a mode are explored in Section IV-A.

C. Electric Field at the Air–Slab Interface

Assuming now the presence of a horizontal infinitesimal electric dipole parallel to the wires, we wish to calculate the electric field at the air–slab interface, still assuming a homogenized model for the slab. Although a rigorous calculation of such a near field would require a full-wave treatment of the interaction between the aperiodic dipole source and its periodic wire-medium environment, e.g., through the array scanning method [15], the results obtained with the homogenized model may help in gaining physical insight into the radiation mechanism of the structure. Furthermore, as will be shown in Section IV-B, when the near field is dominated by a leaky mode, the homogenized results are in good agreement with those obtained via a rigorous full-wave analysis.

The field excited by a dipole parallel to the metal wires is purely \( TM_{y} \) (since by reciprocity neither \( TE_{y} \) nor \( TM_{y} \) fields can be excited), and its tangential components \( \hat{E}_{x} \) and \( \hat{E}_{y} \) at \( z = h \) can be calculated from the spectral form of Maxwell’s equations (21) and (22) and from (27) as

\[
\hat{E}_{y}(k_{x}, k_{y}) = V_{i}^{TM}(k_{x}, k_{y}) \left[ -j \hat{J}_{y}(k_{x}, k_{y}) \right] = -V_{i}^{TM}(k_{x}, k_{y})
\]

\[
\hat{E}_{x}(k_{x}, k_{y}) = -\frac{k_{x}k_{y}}{k_{0}^{2} - k_{y}^{2}} \hat{J}_{y}(k_{x}, k_{y}) = \frac{k_{x}k_{y}}{k_{0}^{2} - k_{y}^{2}} V_{i}^{TM}(k_{x}, k_{y})
\]

where \( \hat{J}_{y} \) is a 1-A for an elemental electric dipole directed along \( y \). The \( TM_{y} \) equivalent voltage \( V_{i}^{TM} \) at \( z = h \) due to a 1-A equivalent impressed current at \( z = h \) [see Fig. 2(c)] is given by

\[
V_{i}^{TM}(k_{x}, k_{y}) = \frac{j \gamma_{0}}{k_{0}} \frac{k_{0}^{2} - k_{y}^{2} \sin(k_{0}^{TM}h)}{k_{0}^{TM} \cos(k_{0}^{TM}h) + j k_{20}^{TM} \sin(k_{0}^{TM}h)}
\]

It is interesting to note that

\[
V_{i}^{TM}(k_{x}, k_{y}) = \left( 1 - \frac{k_{y}^{2}}{k_{0}^{2}} \right) V_{i}^{TE_{z}}(k_{p})
\]
where $V_i^{\text{TE}_c}$ is the TE$_c$ equivalent voltage due to a 1-A current source for an isotropic slab with relative permittivity $\varepsilon_r = 1 - k_0^2/k_0^2$. Since, as is well known, $V_i^{\text{TE}_c}$ depends only on the radial wavenumber $k_p$, the anisotropy of the structure appears only in the factor $(k_0^2 - k_0^2)$ in the numerator of (9).

The tangential components of the electric field at $z = h$ can be calculated in the spatial domain by inverse Fourier transforming (8), i.e.,

$$
E_y(x, y) = \int_{\mathbb{R}} \int_{\mathbb{R}} -\frac{V_i^{\text{TM}}(k_x, k_y)}{(2\pi)^2} e^{-j(k_x x + k_y y)} \, dk_x \, dk_y,
$$

$$
E_x(x, y) = \int_{\mathbb{R}} \int_{\mathbb{R}} \frac{k_x k_y V_i^{\text{TM}}(k_x, k_y)}{(2\pi)^2} e^{-j(k_x x + k_y y)} \, dk_x \, dk_y.
$$

(11)

By performing the standard change of variables in the transverse plane from rectangular to polar in both the spectral and spatial domains [17], from (11) and (10) we obtain for the electric field excited at the air–slab interface by the considered elementary source

$$
E_y(x, y) = \frac{1}{8\pi} \int_{-\infty}^{+\infty} V_i^{\text{TE}_c}(k_p) \left(2k_p - \frac{k_0^2}{k_p^2}\right) H_0^{(2)}(k_p) \, dk_p
$$

$$
- \frac{1}{8\pi} \cos(2\beta) \int_{-\infty}^{+\infty} V_i^{\text{TE}_c}(k_p) \frac{k_0^2}{k_p} H_2^{(2)}(k_p) \, dk_p
$$

(12)

and

$$
E_x(x, y) = -\frac{1}{8\pi} \sin(2\beta) \int_{-\infty}^{+\infty} V_i^{\text{TE}_c}(k_p) \frac{k_0^2}{k_p} H_2^{(2)}(k_p) \, dk_p.
$$

(13)

The integrands in (12) and (13) have pole singularities in the complex $k_p$ plane corresponding to the leaky modes supported by the structure, and branch-point singularities of square-root type at $k_p = \pm k_0$, which give rise to the space-wave contribution [17]. Moreover, there is a branch point at the origin due to the presence of the Hankel functions; however, with a suitable choice of the relevant branch cut (along the imaginary axis), this branch cut lies along the usual Sommerfeld branch cut for the $k_0$ wavenumber [17], and hence, is of no concern in what follows.

As is well known, by deforming the original integration path along the real axis to a vertical pair of paths through the branch point at $k_0$, the components of the total field in (12) and (13) can be decomposed as the sum of a surface-wave field, a leaky-wave field, and a space-wave field [17]

$$
\mathbf{E} = \mathbf{E}^{\text{SW}} + \mathbf{E}^{\text{LW}} + \mathbf{E}^{\text{SpW}}.
$$

(14)

The leaky-wave field is given by the sum of the residue contributions at the complex poles, which are captured by the above-mentioned integration-path deformation; moreover, in the present case, the surface-wave field is absent so that

$$
\mathbf{E} = \sum_n \mathbf{E}_n^{\text{LW}} u_n + \mathbf{E}^{\text{SpW}}
$$

(15)

where $u_n$ is equal to 1 if the $n$th leaky-wave pole is captured, otherwise it is zero. When only one term in the sum is significant, the corresponding leaky mode is said to be dominant. Finally, the space-wave field is given by the integral along the above-mentioned vertical paths through the branch point at $k_0$.

The contribution of the dominant leaky mode to the total field excited by the considered source is evaluated as the residue contribution to the integrals in (12) and (13) at the dominant complex leaky-wave pole $k_p^{\text{LW}} = \beta - j\alpha$ (i.e., the one with the lowest attenuation constant $\alpha$) and it has the expression

$$
E_y^{\text{LW}}(x, y) = -\frac{1}{4} E_0 \left(2k_p^{\text{LW}} - \frac{(k_p^{\text{LW}})^3}{k_0^2}\right) H_0^{(2)}(k_p^{\text{LW}}) \rho
$$

$$
- \frac{1}{4} \cos(2\beta) E_0 \frac{(k_p^{\text{LW}})^3}{k_0^2} H_2^{(2)}(k_p^{\text{LW}}) \rho
$$

(16)

and

$$
E_x^{\text{LW}}(x, y) = -\frac{1}{4} \sin(2\beta) E_0 \frac{(k_p^{\text{LW}})^3}{k_0^2} H_2^{(2)}(k_p^{\text{LW}}) \rho
$$

(17)

where $E_0 = -j\text{Res} \left[ V_i^{\text{TM}}(k_p) \right]_{k_p = k_p^{\text{LW}}}$.

Numerically, the integrations in (12) and (13) are performed by suitably deforming the integration path to avoid the branch-point singularity located at $k_p = k_0$ on the real axis and by adopting standard adaptive quadrature routines. The residue term in (16) and (17) can instead be evaluated by numerical integration of the function $V_i^{\text{TE}_c}(k_p)$ along a small circle enclosing the relevant leaky-wave pole $k_p = k_p^{\text{LW}}$. Finally, the space-wave field is calculated through the numerical integration of the integrand of (12) and (13) along a vertical pair of paths through the branch point at $k_p = k_0$; for large distances from the source, this path becomes a steepest descent path along which the integrand decays exponentially fast [17].

It is noted that in the second addend in (16), as well as the cross-polarized term in (17), the terms in front of the Hankel functions are proportional to $k_p^{\text{LW}}$ for small values of $k_p^{\text{LW}}$. Note that when the leaky mode has a small propagation wavenumber so that $\left| k_p^{\text{LW}} \right| \ll k_0$, the weighting coefficients from the $k_p$ terms are such that the cross-polarized field of the leaky mode is negligible. The wavenumber will be small when the permittivity $\varepsilon_{\text{xyy}}$ is small (corresponding to a weakly attenuated leaky mode that radiates near broadside).

The space-wave field excited by the source is mainly associated with the branch point at $k_p = k_0$, and hence, $k_p$ is not small for this wave. Therefore, the space-wave field has significant cross-polarization. In fact, this field will mainly be polarized perpendicular to the wires since the wires will act to short out the parallel component of the field for this wave. However, if the leaky mode is dominant, the overall field will exhibit a good polarization purity. This will be the case when the permittivity $\varepsilon_{\text{xyy}}$ is small, as will be demonstrated later in the results. These points explain why the radiation patterns in [9]–[11] have a good polarization purity.

III. FULL-WAVE MODAL ANALYSIS

Leaky-wave propagation on the wire-medium slab with air as the host medium is also studied here by means of a rigorous
formulation based on the MoM in the spatial domain. The periodicity of the structure allows for considering one spatial period only (unit cell), thus requiring the use of a periodic (free-space) Green’s function $G_p$. Since the latter is known to be represented by a slowly convergent series, an effective numerical acceleration procedure has been implemented based on the Ewald summation technique [16].

As mentioned in Section II, when the host medium is air, $TE_{y}$ and $TM_{y}$ polarizations do not couple, and they can be studied separately. Furthermore, for very thin wires (as the ones considered here), only the $TM_{y}$ polarization is of interest since, in this limit, the $TE_{y}$ polarization does not interact with the wires and propagates as in free space (and therefore no $TE_{y}$ modes exist). It should also be noted that for the discrete wire-medium structure shown in Fig. 1, $TE_{y}$ multicarrier transmission line modes may exist, which are different than the modal $TM_{y}$ radially propagating solutions that are examined here. However, such modes will not be considered in the following since they cannot be excited by an electric dipole that is parallel to the wires (the type of source considered here). This follows from reciprocity since the dipole is orthogonal to the electric field of these modes.

Due to the exponential dependence of the modal field on the $x$- and $y$-coordinates and the invariance of the structure along the $y$-direction, the electromagnetic problem is reduced to a 2-D one in the $xz$-plane. Moreover, for the $TM_{y}$ polarization of interest, the azimuthal component $J_{y}$ of the current along the wires is identically zero and the MoM unknown is thus the longitudinal component $J_{y}$. The electric field integral equation (EFIE) for wave propagation at an arbitrary angle $\phi$ reads (see [11, Appendix])

$$\sqrt{k_0^2-k_p^2} \int_{\partial C} J_y(\mathbf{r}')G_p\left(\sqrt{k_0^2-k_p^2}|\mathbf{r}-\mathbf{r}'|\right) d\mathbf{r}' = 0, \quad \mathbf{r} \in C$$

(18)

where $C$ indicates the union of the boundaries of the wire cross sections in a unit cell, while $\mathbf{r} = x\mathbf{u}_x + z\mathbf{u}_z$ and $\mathbf{r}' = x'\mathbf{u}_x + z'\mathbf{u}_z$ are the observation and source vectors in the $xz$-plane, respectively. Since $k_{y} = k_{\text{LW}} \sin \phi$, from (18) it is immediately seen that the EFIE for wave propagation at an arbitrary angle $\phi$ is equivalent to that at an angle of $\phi = 0^\circ$ (propagation orthogonal to the wires, i.e., no variation of currents or fields with $y$) provided that the free-space wavenumber $k_0$ is replaced by a scaled effective wavenumber $k_{\text{eff}} = \sqrt{k_0^2-k_p^2}$. Finally, the propagation along $x$ determines the phasing $k_x = k_{\text{LW}} \cos \phi$ used in the periodic Green’s function $G_p$.

The EFIE can be discretized using subdomain (e.g., piecewise constant) basis functions and a Galerkin testing procedure. However, because of the simple (circular) cross section of the wires, entire-domain basis functions can also be effectively used, dramatically reducing the computational effort and the size of the MoM matrix. In particular, exponential functions have been used to expand the unknown current density $J_y$ so that for each cylinder

$$J_y(\mathbf{r}) = J_y(\rho, \gamma) = \delta(\rho - a) \frac{\rho}{\rho} \hat{J}_y(\gamma)$$

$$\hat{J}_y(\gamma) = \sum_{m=-\infty}^{+\infty} i_m e^{im\gamma}$$

(19)

where $(\rho, \gamma)$ are polar coordinates in the $xz$-plane centered on the wire axis, $a$ is the wire radius, and $i_m$ are the unknown expansion coefficients. Since the wires are very thin, typically very few basis functions are needed to obtain an accurate representation of the current density; in the simulations, three basis functions (i.e., $m = 0, \pm 1$) have been used, although using only the $m = 0$ basis function would yield sufficiently accurate results. Once the EFIE has been discretized, the complex wavenumber $k_{\text{LW}} = \beta - j\alpha$ of the leaky modes supported by the wire-medium slab has then been determined by searching for the zeros of the determinant (as a function of $k_p$) of the coefficient matrix of the resulting linear system.

As long as the period is much smaller than the wavelength and (1), the leaky mode has a fundamental Floquet wave that is improper (exponentially increasing in the vertical $z$-direction), while all other Floquet waves are proper (exponentially decreasing in the $z$-direction); this means that, in the calculation of the spectral part of the periodic Green’s function in the Ewald representation, the determination of the square root defining the $p$th transverse harmonic

$$k_{z,p} = \sqrt{k_{\text{eff}}^2 - k_p^2}$$

$$= \sqrt{k_0^2 - (k_{\text{LW}}^2 \sin^2 \phi - \left(k_p^2 \cos \phi + 2\pi \rho_d \frac{\rho}{d_e}\right)^2}$$

(20)

has to be correctly chosen. In particular, $\Im\{k_{z,0}\} > 0$ for the fundamental zeroth harmonic and $\Im\{k_{z,p}\} < 0$ for all the other harmonics (i.e., $p = \pm 1, \pm 2, \ldots$) [18].

IV. NUMERICAL RESULTS

A. Parametric Analysis of Leaky Modes

To validate the conclusions reached with the homogenized model of the wire-medium slab, a parametric investigation of modal propagation has been performed through the rigorous moment-method approach described in Section III.

In Fig. 3, dispersion curves for the fundamental even (TM$_1$) and odd (TM$_2$) modes are reported for propagation orthogonal ($\phi = 0^\circ$) and parallel ($\phi = 90^\circ$) to the wires (black lines), along with the corresponding curves obtained by means of the homogenized model (gray lines). (The notation even/odd refers to the variation of the transverse electric field about the center of the slab. At cutoff, the TM$_1$ mode has one half cycle of variation vertically within the substrate, while the TM$_2$ mode has one cycle of variation.) It can be seen that the full-wave results for both the normalized phase and attenuation constants are almost superimposed with the homogenized-model results. In particular, for propagation parallel to the wires, the agreement is independent of frequency in the range shown, whereas it tends to slightly deteriorate as the frequency increases for propagation orthogonal to the wires.

The frequencies for which $\beta = \alpha$ (i.e., 6.8 GHz for the TM$_1$ mode and 7.4 GHz for the TM$_2$ mode) correspond to frequencies of optimum broadside radiation from the leaky modes [19]. At these frequencies, the electrical thickness of the slab $(2\delta/\lambda_z)$ is 0.45 and 0.90, respectively. The slab is essentially operating as a leaky parallel-plate waveguide, where the air region surrounding the slab presents a low impedance relative to the high-impedance slab material (which has a low permittivity).
For an ideal parallel-plate waveguide (with top and bottom metal walls), the electrical thickness of the waveguide would be exactly 0.5 and 1.0 for the TM\textsubscript{1} and TM\textsubscript{2} modes at cutoff, with cutoff corresponding to radiation at broadside (when radiation is allowed to escape from the waveguide so that the waveguide becomes leaky).

The degree of omnidirectionality of modal propagation achieved in actual periodic wire-medium slabs can be appreciated in Fig. 4. Here, the relative errors $\epsilon_\beta$ and $\epsilon_\alpha$ for the phase and attenuation constants, defined as $\epsilon_\beta(\phi) = |\beta(\phi) - \beta(90^\circ)/\beta(90^\circ)|$ and $\epsilon_\alpha(\phi) = |\alpha(\phi) - \alpha(90^\circ)/\alpha(90^\circ)|$, are reported as a function of the propagation angle $\phi$ for both the fundamental even and odd modes of a structure, as in Fig. 3. It can be seen that for low frequencies (e.g., at $f = 8$ GHz), the maximum relative error in the entire angular range is less than 1%, indicating a very high degree of omnidirectionality of modal propagation. In the high-frequency range (e.g., at $f = 13$ GHz), the maximum relative error for the attenuation constant $\epsilon_\alpha$ is still also less than 2%, while the attenuation constant starts to lose its isotropic character (although its maximum relative error $\epsilon_\alpha$ is less than 10%).

The effect of changing the spatial periods of the wire lattice has also been investigated, maintaining the condition of thin wires [i.e., $a \ll \min(d_x,d_z)$]. In Fig. 5, a comparison between full-wave and homogenized results is presented for the fundamental TM\textsubscript{2} odd mode of a structure, as in Fig. 3, with the spatial period $d_x = d_z = d$ reduced from 10 to 3 mm, propagating in the principal directions $\phi = 0^\circ$ and $\phi = 90^\circ$. As expected, the dispersion curves are shifted towards a higher frequency range with respect to those in Fig. 3; in fact, one effect of reducing the spatial period in a square lattice is to increase the plasma wavenumber. Moreover, although the spatial period is smaller with respect to that of the structure considered in Fig. 3, the period-wavelength ratio has increased so that the accuracy of the homogenized model is decreased; however, overall it remains very good. The small difference between the solid black curves ($\phi = 0$ and $\phi = 90^\circ$) shows that the degree of omnidirectionality also tends to reduce at high frequencies (above approximately 40 GHz).

It is also interesting to consider the effect of letting the spatial periods along the $x$- and $z$-directions be different, i.e., $d_x \neq d_z$, maintaining fixed the other parameters (i.e., the wire radius $a$)
and the number of wire layers $N$). In Fig. 6, dispersion curves are shown for the fundamental TM\textsubscript{2} odd mode propagating at $\phi = 90^\circ$ on four different wire-medium slabs, two with a square lattice ($d_x = d_y = 10$ mm and $d_x = d_y = 5$ mm) and two with a rectangular lattice ($d_x = 10$ mm, $d_y = 5$ mm and $d_x = 5$ mm and $d_y = 10$ mm). In all cases, the simulations are in very good agreement with the homogenized results (not reported for clarity). In particular, the frequencies at which $\beta / k_0 = \alpha / k_0$ are always well predicted by the simple formula $f_{\beta=\alpha} = \sqrt{f_{p}^2 + c^2/(2h)^2}$ [19], where $h$ is the equivalent thickness of the homogenized model, $c$ is the speed of light, and $f_{p} = k_p c/(2\pi)$ is the plasma frequency calculated according to (2) and (3). The values for $\beta / k_0$ and $\alpha / k_0$ at $f = f_{\beta=\alpha}$ given by the approximate formula $c^3/(8\pi f_{p}^2 h^3)$ [19] are also very accurate. It can be seen that by decreasing the wire spacing $d_w$, the equivalent thickness $h$ decreases, thus increasing the value of $\beta$ (and $\alpha$) at $f = f_{\beta=\alpha}$; on the other hand, decreasing the wire spacing $d_p$ decreases the value of $\beta$ (and $\alpha$) at $f = f_{\beta=\alpha}$. Since these value determine the directivity of the beam radiated at broadside by the structure excited by a simple source (e.g., a dipole), these parameters can provide a further degree of freedom in order to obtain a desired directivity at a given frequency when the wire-medium slab is used in antenna applications.

Finally, when the assumption of thin wires is dropped, the homogenization procedure cannot be based on (1) for the dyadic permittivity. In this case, the modal propagation is not omnidirectional anymore, as can be observed in Fig. 7, where dispersion curves are reported for the fundamental odd mode supported by a structure, as in Fig. 3, but with a wire radius ten times larger, i.e., $a = 1$ mm. A large discrepancy can be observed between the dispersion curves in the two propagation directions $\phi = 0^\circ$ and $\phi = 90^\circ$ (black lines); furthermore, the homogenized results based on (1) (gray lines) are azimuthally omnidirectional and do not provide accurate results, as expected.

It may be interesting to note that, at $f \approx 11.4$ GHz, the phase constant becomes equal to zero for propagation at $\phi = 90^\circ$ so that, at frequencies lower than 11.4 GHz, the phase constant would be negative (it is reported in absolute value in Fig. 7). This means that the pole of the spectral Green’s function of the structure corresponding to the considered leaky mode has crossed the negative vertical axis of the complex $k_p$ plane; hence, it has crossed a Sommerfeld branch cut and has moved to a proper sheet of the involved Riemann surface; physically, this corresponds to a transition from a forward to a backward regime for the leaky mode.

**B. Electric Field at the Air–Slab Interface**

In this section, the electric field excited by a horizontal electric dipole placed inside a grounded wire-medium slab of thickness $h$ is studied, adopting both the approximate homogenized model and a rigorous full-wave simulation of an actual truncated wire-medium structure. The horizontal electric dipole is assumed to be parallel to the wire direction and placed at a distance $h_n = h/2$ from the ground plane, equidistant from the four nearby metal wires. The tangential field is calculated on
the plane $z = h$, i.e., at the air–slab interface in the homogenized model and at a distance $d_z/2$ from the center of the wires in the upper row in the actual wire-medium structure.

In Fig. 8, results obtained with the homogenized model are shown for an infinite grounded wire-medium slab with $d_x = d_y = 20$ mm, $a = 0.5$ mm, and $N = 6$ layers of wires (corresponding to homogenized parameters $h = N d_z = 120$ mm and $f_p \approx 3.88$ GHz) at the frequency $f = 4$ GHz. At this frequency, the dominant leaky mode is the TM$_2$ mode with nearly equal values for the normalized phase and attenuation constants $\beta/h_0 \approx \alpha/k_0 \approx 0.1$ such that it radiates a narrow pencil beam with maximum power density at broadside [11]. As explained in Section II-C, the total field at the air–slab interface has been calculated through a numerical evaluation of the integrals in (12) and (13). In Fig. 8(a), the amplitude of the $E_y$ component (parallel to the wires) is shown. As expected from the discussion in Section II-C, the amplitude of this component is almost perfectly independent of the azimuthal angle $\phi$. On the other hand, the amplitude of the $E_x$ component (orthogonal to the wires) is seen in Fig. 8(b) to be approximately two orders of magnitude lower than the amplitude of $E_y$, thus confirming that the field at the air–slab interface is essentially linearly polarized along the wire direction.

As discussed in Section II-C, the total field on the aperture consists of the leaky-mode field and the space-wave field. For a low-permittivity slab that is optimized for maximum radiation at broadside (as in Fig. 8), the leaky-mode field is expected to dominate the total field of the aperture. To verify this, results are shown in Figs. 9 and 10 for the leaky-mode field and the space-wave field on the aperture for the case shown in Fig. 8 for the co-polarized ($E_x$) and cross-polarized fields ($E_y$), respectively. The leaky-mode field and space-wave field have been numerically evaluated, as explained in Section II-C. The total field on the aperture is also shown, plotted on the same scale as the leaky-mode and space-wave fields for an easy comparison.

Fig. 9 verifies that the leaky-mode field is the dominant contributor to the co-polarized field of the aperture with the amplitude of the co-polarized space-wave field being on the order of 50 dB lower. On the other hand, Fig. 10 shows that the space-wave contribution to the cross-polarized field on the aperture is not negligible, its amplitude being comparable (but a bit weaker)
with that of the leaky-mode contribution (with the cross-polarized field of both the leaky mode and the space-wave field vanishing along the principal axes); this gives rise to some oscillations in the cross-polarized component of the total field.

The co-polarized field of the leaky mode is almost perfectly omnidirectional in azimuth, as expected. The co-polarized field of the space-wave field is instead maximum in the wire direction (i.e., along the $y$-axis), while it decays more rapidly orthogonally to the wires (i.e., along the $x$-direction). The cross-polarized fields of both the leaky-mode field and the space-wave field are maximum in the diagonal planes, and zero along the principal axes.

Fig. 11 shows a plot of the co-polarized and cross-polarized aperture fields of the leaky mode versus distance from the source at a fixed angle of $\phi = 45^\circ$. On this scale, it is obvious that both the co-polarized and cross-polarized leaky-mode fields decay exponentially, as expected.

Fig. 12 investigates the behavior of the space-wave field on the aperture, showing both the co-polarized and the cross-polarized fields together with a reference curve that is a plot of a normalized $1/\rho^2$ function along the $\phi = 45^\circ$ line. It is seen that the space-wave field is mainly cross-polarized near the source, but further away the co-polarized and cross-polarized fields have almost the same amplitude, and both decay as $1/\rho^2$. This decay rate is the same as is observed for sources in the presence of general layered media, such as a dipole over a dielectric half-space.

Due to the different decay rates, the amplitude of the space-wave field will eventually surpass that of the leaky-mode field far away from the source, but this occurs at a distance of approximately 18.5 wavelengths, and is off the scale shown in the figures. This explains why the broadside radiation from such a structure is mainly due to the leaky mode [9]–[11].
The excitation of a grounded wire-medium slab by means of a horizontal electric dipole has been studied. The main features of the field at the air–slab interface, i.e., its character of being almost perfectly linearly polarized along the wire direction and omnidirectional in magnitude, have been ascertained through a MoM simulation of an actual truncated wire-medium structure, and have been shown to be correctly predicted by the approximate homogenized model.

Future investigations will concern the input-impedance properties of more realistic sources placed inside the considered wire-medium slabs, and they will necessarily require the rigorous treatment of the near-field interaction between the source and metal wires through full-wave simulations of the actual periodic material.

**APPENDIX**

**Equivalent Transmission Lines for TE_y and TM_y Fields**

Here, the transverse equivalent networks in the $z$-direction are derived for $\text{TE}_y$ and $\text{TM}_y$ waves inside an electrically uniaxial medium with a spatially and temporally dispersive permittivity, as in (1), starting from the spectral form of Maxwell’s equations and taking into account the presence of impressed electric and magnetic currents.
By Fourier transforming the Maxwell equations with respect to \( x \) and \( y \), we have

\[
\begin{align*}
-jk_y \hat{E}_z - \frac{\partial \hat{B}_y}{\partial z} &= -j\omega \mu \hat{H}_x - \hat{M}_x \\
\frac{\partial \hat{E}_z}{\partial z} + jk_x \hat{E}_y &= -j\omega \varepsilon \hat{H}_y - \hat{M}_y \\
-jk_x \hat{E}_y + jk_y \hat{E}_x &= -j\omega \varepsilon \hat{H}_x + \hat{M}_x \\
-jk_y \hat{H}_z - \frac{\partial \hat{H}_y}{\partial z} &= j\omega \varepsilon \varepsilon_{\text{rt}} \hat{E}_y + \hat{j}_x \\
\frac{\partial \hat{H}_z}{\partial z} + jk_x \hat{H}_y &= j\omega \varepsilon \varepsilon_{\text{rt}} \hat{E}_x + \hat{j}_y
\end{align*}
\]

(21)

In the TE\(_{xy}\) case, \( \hat{E}_y = 0 \) and \( \hat{j}_y = 0 \), and by simple manipulations of (21) and (22), it is possible to express all the remaining field components in terms of \( \hat{E}_x \) and \( \hat{H}_y \) of and of the impressed currents. The second equation of (21) and the first equation of (22) can then be written as transmission line equations

\[
\begin{align*}
\frac{dV_{\text{TE}}}{dz} &= -Z_{\text{TE}}^{s} I_{\text{TE}} + v_{\text{TE}} \\
\frac{dI_{\text{TE}}}{dz} &= -Y_{\text{p}}^{s} V_{\text{TE}} + i_{\text{TE}}
\end{align*}
\]

(23)

where

\[
\begin{align*}
V_{\text{TE}} &= \hat{E}_x \\
I_{\text{TE}} &= \hat{H}_y \\
Z_{\text{s}} &= j\omega \mu \left( 1 - \frac{k_x^2}{k_0^2 \varepsilon_{\text{rt}}} - \frac{k_y^2}{k_0^2 \varepsilon_{\text{rt}}} \right) \\
Y_{\text{p}} &= \frac{j}{\omega \mu} \left( \frac{k_x^2}{k_0^2 \varepsilon_{\text{rt}}} - \frac{k_y^2}{k_0^2 \varepsilon_{\text{rt}}} \right) \\
v_{\text{TE}} &= \frac{j}{\omega \mu} k_x \frac{k_x^2}{k_0^2 \varepsilon_{\text{rt}}} - \frac{k_y}{k_0} \hat{M}_x + \hat{M}_y \\
i_{\text{TE}} &= -\frac{k_y}{\omega \mu} \hat{M}_y - \hat{j}_x.
\end{align*}
\]

(24)

The propagation constant and characteristic impedance of the TE\(_{xy}\) equivalent transmission line are then

\[
\begin{align*}
k_z^2 &= -Z_{\text{TE}}^{s} Y_{\text{p}}^{s} = k_0^2 \varepsilon_{\text{rt}} - \frac{k_x^2}{k_0^2} - \frac{k_y^2}{k_0^2} \\
Z_{\text{C}}^{s} &= \sqrt{-Z_{\text{TE}}^{s} Y_{\text{p}}^{s}} = k_0 \sqrt{\varepsilon_{\text{rt}} \frac{k_z^2}{k_0^2}} - \frac{k_x^2}{k_0^2} - \frac{k_y^2}{k_0^2}.
\end{align*}
\]

(25)

Note that, by letting \( \beta_0 = \sqrt{\varepsilon_{\text{rt}} k_0} \), from the first equation of (25), it is inferred that

\[
\beta_0 = k_0 \sqrt{\varepsilon_{\text{rt}}}. 
\]

(26)

Hence, the normalized phase constant \( \beta_0/k_0 \) (i.e., the effective index \( n_0 \)) of a uniform TE\(_{xy}\) plane wave (i.e., an ordinary wave) inside the considered medium is equal to \( \sqrt{\varepsilon_{\text{rt}}} \) regardless of its direction of propagation, as expected.

In the TM\(_{xy}\) case, \( \hat{H}_y = 0 \) and \( \hat{M}_y = 0 \); reasoning as in the TE\(_{xy}\) case, the first equation of (21), and the second equation of (22) can then be written as transmission line equations, as in (23), but with a TM superscript where

\[
\begin{align*}
V_{\text{TM}} &= \hat{E}_y \\
I_{\text{TM}} &= -\hat{H}_x \\
Z_{\text{s}} &= \frac{j}{\omega \varepsilon_{\text{rt}}} \left( \frac{k_x^2}{k_0^2} - \frac{k_y^2}{k_0^2} \right) \\
Y_{\text{p}} &= j\omega \varepsilon_0 \left( \varepsilon_{\text{ry}} - \frac{k_x^2}{k_0^2 \varepsilon_{\text{rt}}} - \frac{k_y^2}{k_0^2 \varepsilon_{\text{rt}}} \right) \\
v_{\text{TM}} &= \frac{k_x}{\omega \varepsilon_{\text{rt}}} \hat{J}_z + \hat{M}_x \\
i_{\text{TM}} &= \frac{k_y}{k_0^2 \varepsilon_{\text{rt}}} - \frac{k_y}{k_0^2 \varepsilon_{\text{rt}}} \hat{J}_y - \frac{\omega \varepsilon_{\text{ry}} k_y}{k_0^2 \varepsilon_{\text{rt}}} \hat{M}_x - \hat{J}_y.
\end{align*}
\]

(27)

The propagation constant and characteristic impedance of the TM\(_{xy}\) equivalent transmission line are then

\[
\begin{align*}
k_z^2 &= -Z_{\text{s}}^{\text{TM}} Y_{\text{p}}^{\text{TM}} = k_0^2 \varepsilon_{\text{ry}} + \left( 1 - \frac{k_x^2}{k_0^2 \varepsilon_{\text{rt}}} \right) \left( \frac{k_y}{k_0} \right)^2 - k_x^2 - k_y^2 \\
Z_{\text{C}}^{\text{TM}} &= \frac{n_0}{k_0} \sqrt{\varepsilon_{\text{ry}} - \frac{k_x^2}{k_0^2} \varepsilon_{\text{rt}}} - \frac{k_y^2}{k_0^2}.
\end{align*}
\]

(28)

Again, by letting \( \beta_0 = \sqrt{k_x^2 + k_y^2 + k_z^2} \), one obtains from the first equation of (28) and (1)

\[
\beta_0 = k_0 \sqrt{\varepsilon_{\text{ry}} + \left( 1 - \frac{k_y^2}{k_0^2 \varepsilon_{\text{rt}}} \right) \left( \frac{k_y}{k_0} \right)^2} = k_0 \sqrt{\varepsilon_{\text{ry}} - \frac{k_x^2}{k_0^2} \varepsilon_{\text{rt}}}.
\]

(29)

From (29), it is concluded that the normalized phase constant \( \beta_0/k_0 \) (i.e., the effective index \( n_0 \)) of a uniform TM\(_{xy}\) plane wave (i.e., an extraordinary wave) inside the considered medium is equal to \( \sqrt{\varepsilon_{\text{ry}} - \frac{k_x^2}{k_0^2} \varepsilon_{\text{rt}}} \), regardless of its direction of propagation. The fact that \( \beta_0 \) is independent of the direction of propagation is a consequence of the spatially dispersive model adopted for the effective medium and expressed by (1), as observed in [3].

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