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Surface Wave Propagation on a Rectangular Bonded Wire Mesh Located Over the Ground

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A mode equation is derived for surface wave propagation along a rectangular bonded wire mesh over a lossy half-space. The mode equation is solved numerically for the complex propagation constant of the surface wave. For a sufficiently small mesh size, the attenuation rate of this surface wave is considerably less than that of the Zenneck surface wave for an isolated half-space or homogeneous ground.

INTRODUCTION

Wire mesh screens are employed in numerous shielding and reflecting applications. The relevant plane wave scattering properties have been analyzed both for meshes in free space [*Kontorovich et al.*, 1962; *Astrakhan*, 1968; *Hill and Wait*, 1974; *Hill and Wait*, 1976] and over a lossy earth [*Otteni*, 1973; *Wait and Hill*, 1976]. The closely related problem of surface wave propagation on a wire mesh in free space has also been analyzed [*Hill and Wait*, 1977a; *Hill and Wait*, 1977b] and studied experimentally [*Ulrich and Tacke*, 1973].

INTRODUCTION

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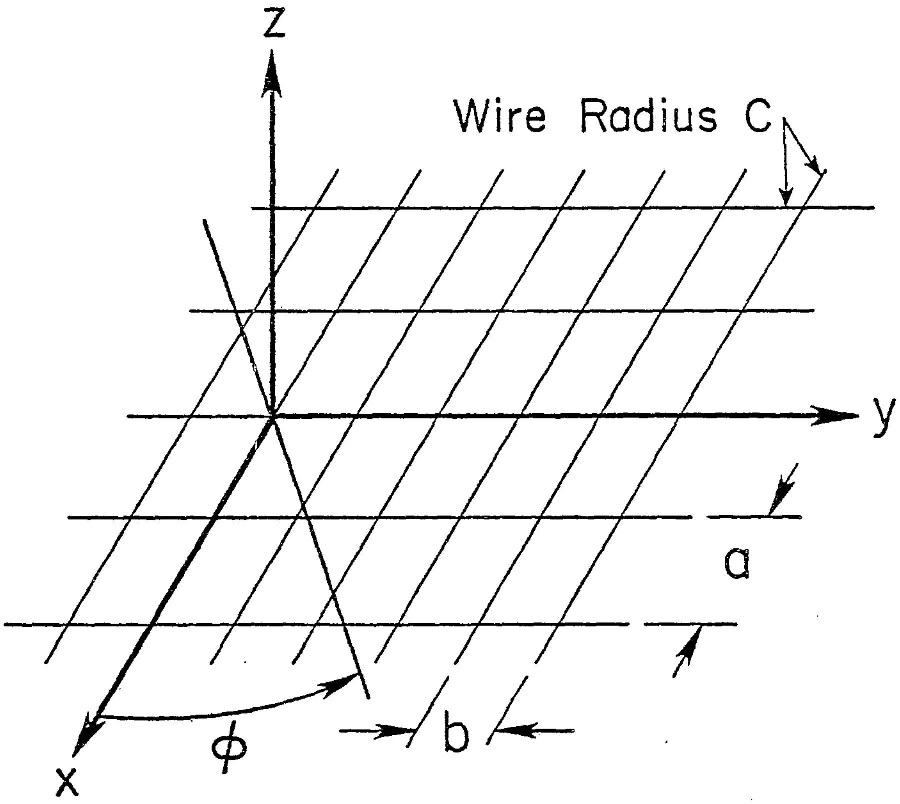
Here we extend our analysis to the case where the rectangular mesh is located at an arbitrary height above a conducting half space. When the height is small, this configuration can be considered as a model to study surface wave propagation over a ground screen. The surface wave contribution to the total field is expected to be significant when the source and observer are located near the ground screen. For example, a conical monopole source over a ground screen has been used to illuminate test objects (also located above the ground screen) with an electromagnetic pulse [*Kehrer and Baum*, 1975].

Our objective here is to calculate the propagation constant of the surface wave. Due to the lossy earth, the wave suffers attenuation and the propagation constant becomes complex.

FORMULATION

The geometry of a rectangular bonded mesh located at a height d above a conducting half-space is illustrated in Figure 1. Arrays of identical perfectly conducting wires parallel to the x axis with spacing b and parallel to the y axis with spacing a are contained in the plane $z = 0$. This configuration is called a bonded rectangular mesh because the contact between the wire junction is bonded and the interwire spacings a and b are not equal in general. Furthermore, the wire radius c is small compared to the spacings a and b, the mesh height d, and the free space wavelength λ. Consequently, only the axial wire currents are important, and the usual thin wire approximations are valid.

The region $z > -d$, external to the wires, is free space with permittivity ε_o and permeability μ_o. The region $z < -d$ is homogeneous with permittivity ε_g, conductivity σ_g, and free space permeability μ_o.



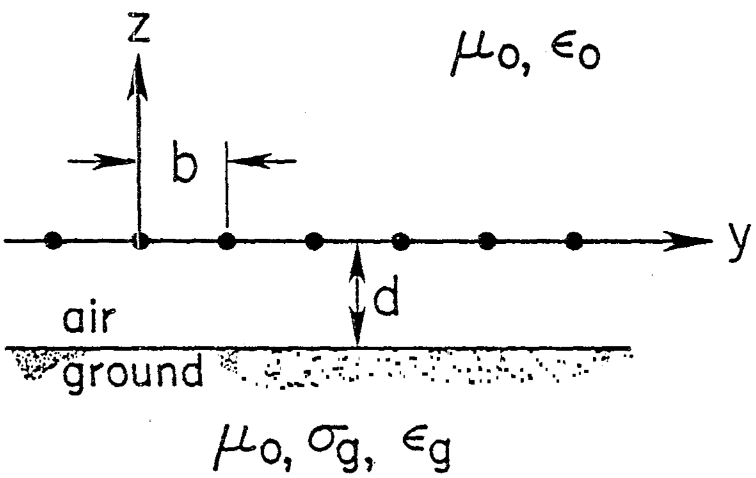


Figure 1. A bonded rectangular wire mesh over a conducting half-space (perspective and side view).

In the following formulation, the mesh is located in free space which means that d is positive. However, a similar formulation is possible for a buried mesh $(d<0)$.

The formulation closely follows that for a rectangular bonded mesh in free space [*Hill and Wait*, 1977b]. Thus, we seek modes which are propagating in x and y but which decay above the mesh $(z>0)$ and in the earth $(z<-d)$. We invoke Floquet's theorem [*Collin*, 1960] in order to express the relevant electromagnetic quantity as an exponential function multiplied by a function which is periodic in x and y. Thus, for a single mode propagating at an angle ϕ to the negative x axis, the current on the qth x-directed wire I_{xq} and the current on the mth y-directed wire I_{ym} can be written:

$$I_{xq} = \exp[\gamma(x \cos\phi + qb \sin\phi)] \sum_m A_m \exp(i2\pi mx/a) \tag{1}$$

and

$$I_{ym} = \exp[\gamma(ma \cos\phi + y \sin\phi)] \sum_q B_q \exp(i2\pi qy/b) \tag{2}$$

where a time factor $\exp(i\omega t)$ is assumed. Here A_m and B_q are the unknown Fourier coefficients, and γ is the propagation constant of the particular mode which we seek. The m and q summations are over all integers including zero from $-\infty$ to ∞.

The calculation of the fields produced by the currents given by (1) and (2) in the presence of a conducting half-space is straightforward [*Wait and Hill*, 1976]. For the present analysis, we employ the following thin wire boundary condition for the assumed perfectly conducting wires:

$$E_x(x,o,c) = E_y(y,o,c) = 0. \tag{3}$$

Although (3) is only applied to the $m = 0$ and $q = 0$ wires, the periodic Floquet form of (1) and (2) assures that the boundary condition will be satisfied at all wires.

Expressions (1) and (2) for the current are identical to those in the plane wave scattering case except that γ has replaced ikS where k is the free space wave number $(=2\pi/\lambda)$ and S is the sine of the incidence angle. Thus the previous equations for A_m and B_q can be used with the following modifications: 1) set the incident fields equal to zero (source-free problem), 2) set the grid separation h equal to zero (bonded grids in the same plane), and 3) set the wire impedance equal to zero perfectly conducting wires. As a result, equations (24) and (26) from the plane wave case [*Wait and Hill*, 1976] reduce to the following:

$$A_m \frac{(k^2-k_x^2)P_m}{2ikb} + \frac{ik_x}{2ka} \sum_q B_q k_y \Gamma^{-1}\Big[\exp(-\Gamma c) + \left(r_{mq} + \frac{ik\eta\Gamma s_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\Big] = 0 \,, \tag{4}$$

$$B_q \frac{(k^2-k_y^2)Q_q}{2ika} + \frac{ik_y}{2ka} \sum_m A_m k_x \Gamma^{-1}\Big[\exp(-\Gamma c) + \left(R_{mq} - \frac{ik\eta\Gamma S_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\Big] = 0 \,, \tag{5}$$

where

$$P_m = \sum_q [\exp(-\Gamma c) + R_{mq}\exp(-2\Gamma d)]\Gamma^{-1} \,, \tag{6}$$

$$Q_q = \sum_m [\exp(-\Gamma c) + r_{mq}\exp(-2\Gamma d)]\Gamma^{-1} \tag{7}$$

$$\Gamma_{mq}(=\Gamma) = (k_x^2+k_y^2-k^2)^{1/2} \,, \tag{8}$$

$$k_x = (2\pi m/a) + kS\cos\phi\ , \tag{9}$$

$$k_y = (2\pi q/b) + kS\sin\phi\ , \tag{10}$$

and $\eta = (\mu_o/\varepsilon_o)^{1/2}$. S, being a complex quantity, is now defined as $\gamma/(ik)$.

The quantities R_{mq}, r_{mq}, S_{mq} and s_{mq} are functions of the half-space parameters σ_g and ε_g:

$$R_{mq} = \frac{k^2(\Gamma+\hat{\Gamma}K_x)(\Gamma-\hat{\Gamma}\varepsilon_r K_x) + (k_x k_y)^2(1-K_x)^2}{k^2(\Gamma+\hat{\Gamma}K_x)(\Gamma+\hat{\Gamma}\varepsilon_r K_x) - (k_x k_y)^2(1-K_x)^2}\ , \tag{11}$$

$$r_{mq} = \frac{k^2(\Gamma+\hat{\Gamma}K_y)(\Gamma-\hat{\Gamma}\varepsilon_r K_y) + (k_x k_y)^2(1-K_y)^2}{k^2(\Gamma+\hat{\Gamma}K_y)(\Gamma+\hat{\Gamma}\varepsilon_r K_y) - (k_x k_y)^2(1-K_y)^2}\ , \tag{12}$$

$$S_{mq} = \frac{-2ik_x k_y \Gamma(1-K_x)/\eta}{k^2(\Gamma+\hat{\Gamma}K_x)(\Gamma+\hat{\Gamma}\varepsilon_r K_x) - (k_x k_y)^2(1-K_x)^2}\ , \tag{13}$$

$$s_{mq} = \frac{2ikk_x k_y \Gamma(1-K_y)/\eta}{k^2(\Gamma+\hat{\Gamma}K_y)(\Gamma+\hat{\Gamma}\varepsilon_r K_y) - (k_x k_y)^2(1-K_y)^2}\ , \tag{14}$$

where

$$K_x = (k^2-k_x^2)/(k_g^2-k_x^2)\ ,$$

$$K_y = (k^2-k_y^2)/(k_g^2-k_y^2)\ ,$$

$$\hat{\Gamma} = (k_x^2+k_y^2-k_g^2)^{1/2}\ ,$$

$$k_g^2 = k^2\varepsilon_r\ ,$$

$$\varepsilon_r = (\sigma_g+i\omega\varepsilon_g)/(i\omega\varepsilon_o)\ .$$

For two special cases, we note that (11) - (14) simplify considerably. When the half space vanishes (i.e. $\varepsilon_r = 1$), we have $R_{mq} = r_{mq} = S_{mq} = s_{mq} = 0$ which is the free space result [*Hill and Wait*, 1977b]. When the half space is perfectly conducting (i.e. $\sigma_g = \infty$), we have $R_{mq} = r_{mq} = -1$ and

$S_{mq} = s_{mq} = 0$. By image theory, this is the result for a pair of identical meshes [*Hill*, 1977].

The summations involving $\exp(-\Gamma c)$ in (6) and (7) are slowly convergent as they stand. More rapidly convergent forms have been derived for P_m and Q_q in the free space case [*Hill and Wait*, 1977b] and they can be applied here to yield

$$P_m = \frac{b}{\pi}\left\{-\ln[1 - \exp(-2\pi c/b)] + \Delta_m\right\} + \exp(-\Gamma_{mo}c)\Gamma_{mo}^{-1} + \sum_q R_{mq}\exp(-2\Gamma d)\Gamma^{-1} \ , \tag{15}$$

$$Q_q = \frac{a}{\pi}\left\{-\ln[1 - \exp(-2\pi c/a)] + \delta_q\right\} + \exp(-\Gamma_{oq}c)\Gamma_{oq}^{-1} + \sum_m r_{mq}\exp(-2\Gamma d)\Gamma^{-1} \ , \tag{16}$$

where

$$\Delta_m = \frac{1}{2}\sum_q{}' \left[\frac{2\pi}{b}\ \frac{\exp(-\Gamma c)}{\Gamma} - \frac{\exp(-2\pi|q|c/b)}{|q|}\right] \tag{17}$$

and

$$\delta_q = \frac{1}{2}\sum_m{}' \left[\frac{2\pi}{a}\ \frac{\exp(-\Gamma c)}{\Gamma} - \frac{\exp(-2\pi|m|c/a)}{|m|}\right] \tag{18}$$

The superscripted prime over the summation sign indicates omission of the $q = 0$ (or $m = 0$) term.

The doubly infinite set of linear equations (4) and (5) for A_m and B_q is numerically inefficient in the present form because A_m and B_q decay slowly for large $|m|$ and $|q|$. The difficulty arises because the current expansions (1) and (2) are slowly convergent for the discontinuous current that occurs at the wire junctions in bonded meshes. We can circumvent the convergence problem by modifying the current expansions to allow for a jump discontinuity at the origin. The procedure is nearly identical to that employed for the rectangular bonded mesh in free space [*Hill and Wait*, 1977b]. Thus the Fourier coefficients of the current A_m and B_q are rewritten

$$A_m = A_m' + \Delta(1-\delta_{mo})/(2\pi im) \tag{19}$$

and

$$B_q = B_q' - \Delta(1-\delta_{qo})/(2\pi iq) \tag{20}$$

where

$$\delta_{mo} = \begin{cases} 1, & m = 0 \\ 0, & m \neq 0 \end{cases}$$

A_m' and B_q' are modified current coefficients, and Δ is an unknown current discontinuity in I_{xo} at $x = 0$. By substituting (19) and (20) into (4) and (5), we obtain the following equivalent set of linear equations for the modified coefficients:

$$\begin{aligned} A_m' \frac{(k^2-k_x^2)P_m}{2ikb} &+ \frac{ik_x}{2ka} \sum_q B_q' k_y \Gamma^{-1}\Big[\exp(-\Gamma c) \\ &+ \left(r_{mq} + \frac{ik\eta\Gamma s_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\Big] \\ + \Delta\Bigg\{ \frac{(k^2-k_x^2)P_m}{2kb} \frac{(\delta_{mo}-1)}{2\pi m} &- \frac{k_x}{2ka}\left[\frac{P_m'}{b} + \frac{kS\sin\phi}{2\pi} P_{1m}\right]\Bigg\} = 0 \end{aligned} \tag{21}$$

$$\begin{aligned} B_q' \frac{(k^2-k_y^2)Q_q}{2ika} &+ \frac{ik_y}{2ka} \sum_m A_m' k_x \Gamma^{-1}\Big[\exp(-\Gamma c) \\ &+ \left(R_{mq} - \frac{ik\eta\Gamma S_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\Big] \\ + \Delta\Bigg\{ \frac{(k^2-k_y^2)}{2ka} \frac{(1-\delta_{qo})}{2\pi q} &+ \frac{k_y}{2kb}\left[\frac{Q_q'}{a} + \frac{kS\cos\phi}{2\pi} Q_{1q}\right]\Bigg\} = 0\ , \end{aligned} \tag{22}$$

where

$$P_{1m} = \sum_q{}' \left[\exp(-\Gamma c) + \left(r_{mq} + \frac{ik\eta\Gamma s_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\right]/(q\Gamma) \quad , \tag{23}$$

$$P_m' = \frac{b}{\pi} \left\{- \ln[1 - \exp(-2\pi c/b)] + \Delta_m\right\} + \sum_q{}' \left(r_{mq} + \frac{ik\eta\Gamma s_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\Gamma^{-1} \quad , \tag{24}$$

$$Q_{1q} = \sum_m{}' \left[\exp(-\Gamma c) + \left(R_{mq} - \frac{ik\eta\Gamma S_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\right]/(m\Gamma) \quad , \tag{25}$$

and

$$Q_q' = \frac{a}{\pi} \left\{- \ln[1 - \exp(-2\pi c/a)] + \delta_q\right\} + \sum_m{}' \left(R_{mq} - \frac{ik\eta\Gamma S_{mq}}{k_x k_y}\right) \exp(-2\Gamma d)\Gamma^{-1} \tag{26}$$

Again the superscipt prime on the summation indicates omission of the $q = 0$ (or $m = 0$) term. Note that by setting Δ equal to zero in (21) and (22), we could retrieve (4) and (5).

Since we have introduced an additional unknown Δ, another equation is required to have an equal number of equations and unknowns (A_m', B_q', and Δ). The following equation can be obtained from charge continuity at the junctions [*Hill and Wait,* 1977b]:

$$-\frac{\Delta}{2\pi}\left(1 + \frac{b}{a}\right) + \sum_m A_m' \left(im\frac{b}{a} + \frac{ikSb}{2\pi}\cos\phi\right) - \sum_q B_q' \left(iq + \frac{ikSb}{2\pi}\sin\phi\right) = 0 \tag{27}$$

Since the doubly infinite set of equations, (21) and (22), are rapidly convergent, they can be truncated with m ranging from -M to M and q ranging from -Q to Q where M and Q are small integers. Thus (21), (22), and (27) yield a set of $T(=2N+2Q+3)$ linear, homogeneous equations in A'_m, B'_q, and Δ:

$$\begin{bmatrix} \\ \\ \text{T X T} \\ \text{coefficient} \\ \text{matrix} \\ \\ \\ \end{bmatrix} \begin{bmatrix} A'_{-M} \\ \vdots \\ A'_o \\ \vdots \\ A'_M \\ \vdots \\ B'_{-Q} \\ \vdots \\ B'_o \\ \vdots \\ B'_Q \\ \vdots \\ \Delta \end{bmatrix} = \begin{bmatrix} 0 \\ \vdots \\ \vdots \\ \vdots \\ \vdots \\ \vdots \\ 0 \end{bmatrix} \qquad (28)$$

A nontrivial solution to (26) exists only if the determinant, which is a function of $S(=\gamma/ik)$, vanishes. Thus the mode equation to be solved for S is:

$$\begin{vmatrix} \text{T X T} \\ \text{coefficient} \\ \text{matrix} \end{vmatrix} = 0 \qquad (29)$$

The above equation has been programmed and solved numerically for S by Newton's method.

NUMERICAL RESULTS

Convergence of the mode equation (29) was examined by increasing M and Q until the value of $S(=\gamma/ik)$ did not change significantly. The most rapid convergence was obtained by making $M = Q$ for a/b ratios from 1 to 3. For the cases considered here, convergence was obtained for $M = Q = 2$ (T=11), and all results shown here were computed for $M = Q = 2$. The required determinant calculation is fairly rapid for the resultant 11×11 matrix. The matrix fill time dominates the determinant calculation time for such cases.

In Fig. 2, we illustrate the ϕ dependence of Re(S) for several a/b ratios. Note that for the mesh in free space, Re(S) is always greater than one (slow wave). For the half space environment where S is complex, this is not always so. The relative dielectric constant $\varepsilon_r = 10 - i1.8$ would correspond to a ground conductivity $\sigma_g = 10^{-2}$ mho/m and relative permittivity $\varepsilon_g/\varepsilon_o = 10$ at a frequency of 100 MHz. The lack of ϕ dependence for $a/b = 1$ is to be expected for square bonded meshes which are electrically small (e.g. $b/\lambda = 0.05$). When a preferred direction of propagation exists, a rectangular mesh (i.e. where $a \neq b$) can be useful, and a 3 to 1 mesh has been used in some EMP simulator applications [*Baum*, 1972; *Kehrer and Baum*, 1975]. Note that Re(S) is closer to unity at $\phi = 0°$ for the rectangular mesh, but that the ϕ dependence is quite strong. A value of $c/b = 10^{-2}$ has been used in all calculations shown here, but the results are only weakly dependent on the wire radius c.

For the same parameters, in Fig. 3 we illustrate the ϕ dependence of Im(S) for three a/b ratios. The actual attenuation rate α is determined from S by

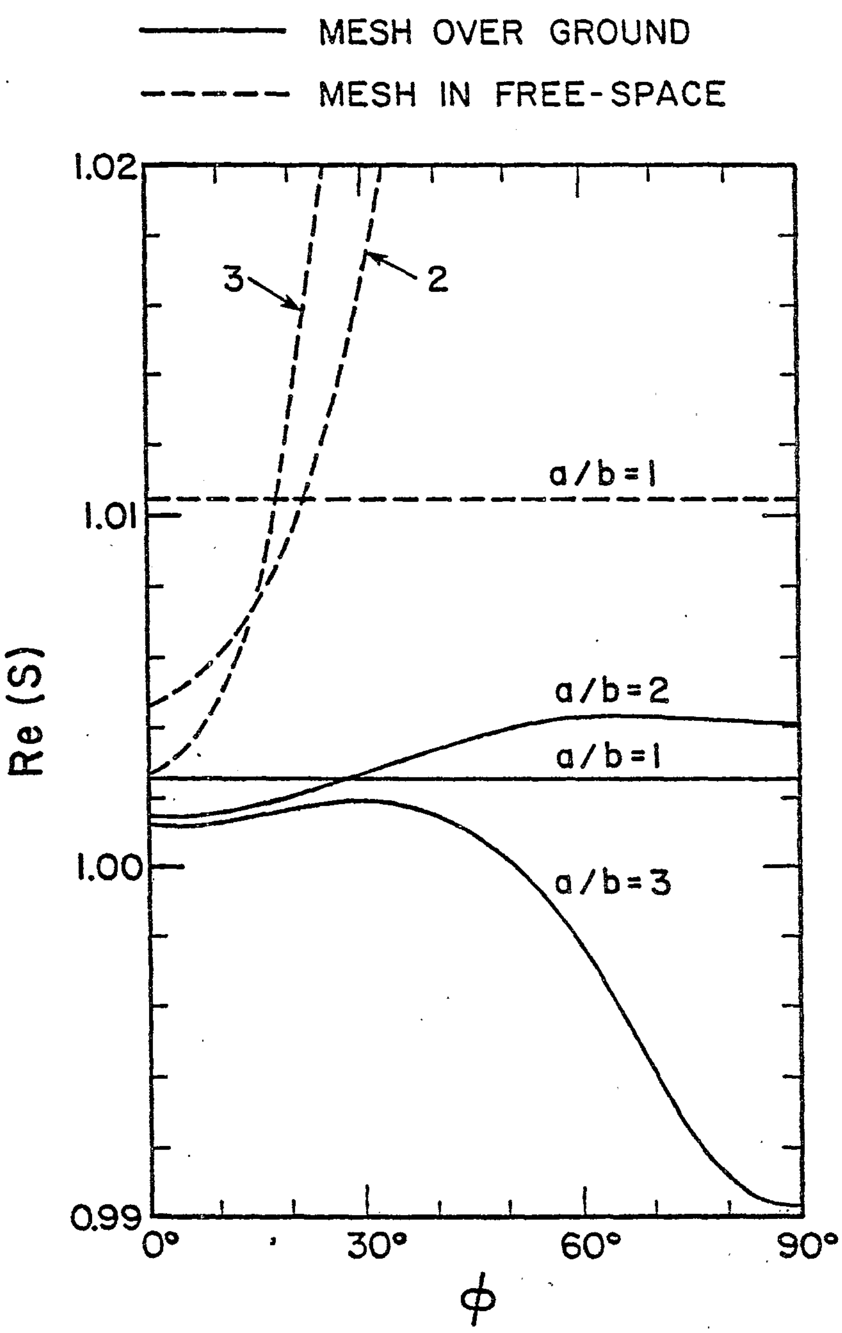


Figure 2. Real part of the normalized propagation constant S as a function of propagation direction ϕ. Parameters: $c/b = 10^{-2}$, $d/b = 10^{-1}$, $b/\lambda = 0.05$, $\varepsilon_r = 10 = i1.8$.

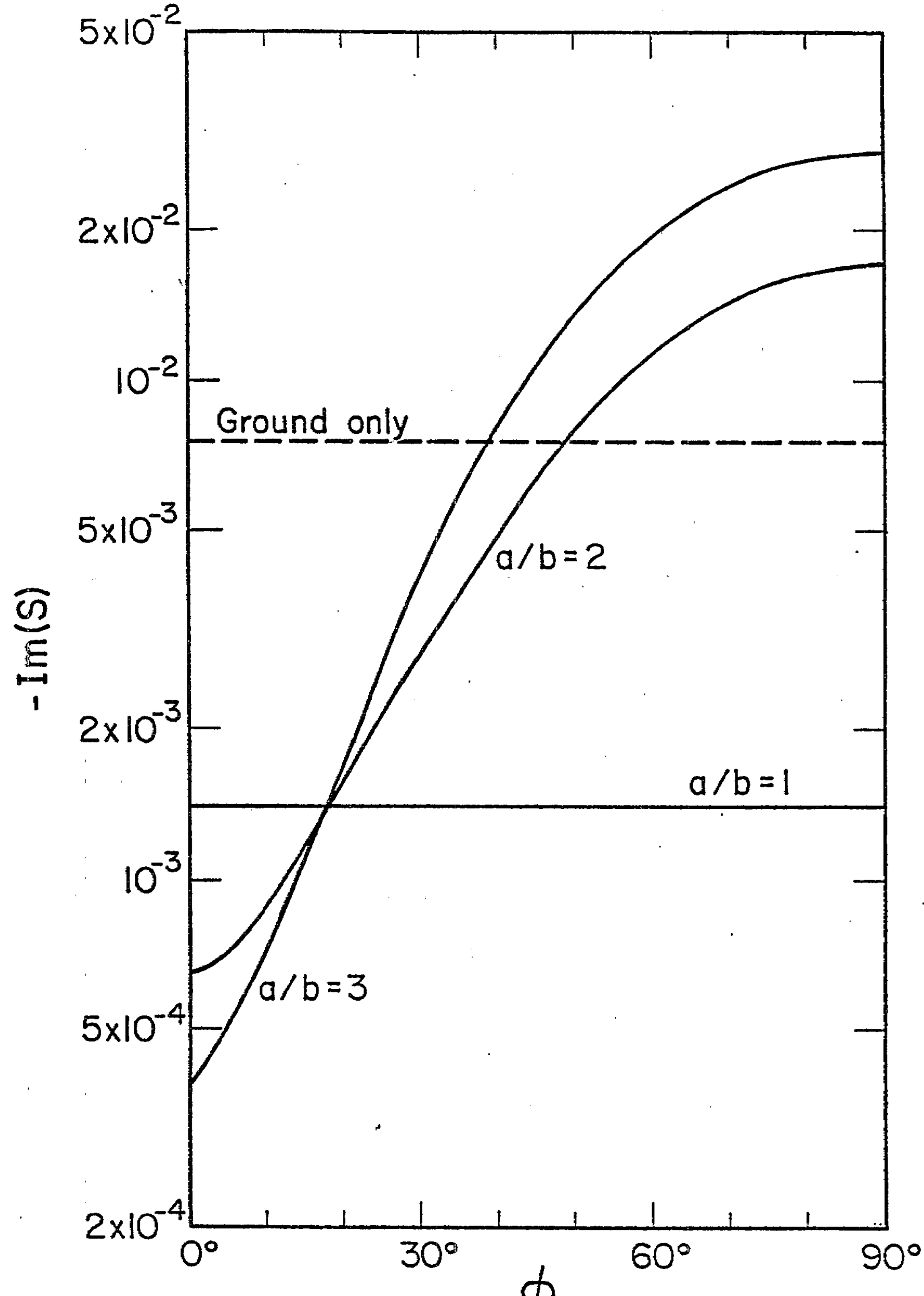


Figure 3. Imaginary part of the normalized propagation constant S as a function of direction ϕ.

$$\alpha = \mathrm{Re}(\gamma) = -k\ \mathrm{Im}(S)\ \text{(nepers/m)}\ . \tag{30}$$

No results are shown for the free space comparison because the free space mesh is lossless (S real). As expected, the square mesh is again isotropic, and the rectangular mesh is highly anisotropic. Also, the attenuation rate is lowest when propagating in the direction of the more closely spaced wires ($\phi=0°$). For comparison, the result for the Zenneck surface wave of the isolated half space is shown. It is computed from [*Wait*, 1962a; *Banos*, 1966]:

$$S = [\varepsilon_r/(\varepsilon_r+1)]^{1/2} \tag{31}$$

For $\varepsilon_r = 10 - i1.8$, S is approximately 0.9547 - i0.0076. Note the large reduction in attenuation at $\phi = 0°$ due to the mesh.

In Table 1, we illustrate the dependence of S on b/λ and d/b for $\phi = 0°$. As before, $\varepsilon_r = 10 - i1.8$ and $c/b = 10^{-2}$. Note that the attenuation rate decreases as d/b is increased from 0.1 to 0.3. This is expected because ground screen performance generally improves as the screen is slightly elevated [*Wait*, 1962b]. The results for $d/b = \infty$ are those of the mesh in free space [*Hill and Wait*, 1977b]. The variation of S is not necessarily monotonic as d/b is increased from 0.1 to ∞. Also note the large increase in both the real and imaginary parts of S as b/λ is increased from 0.05 to 0.1, particularly for the square mesh.

CONCLUDING REMARKS

A general mode equation has been derived for propagation along a bonded rectangular mesh over a lossy earth. The mode equation has been solved numerically for the dominant surface wave mode. This mode suffers attenuation due to the losses in the earth, and for most cases is a slow wave

TABLE 1

Normalized Complex Propagation Constant, S (ε_r = 10 - i1.8, c/b = 10^{-2}, ϕ = 0°)

a/b	b/λ	d/b	S
1.0	0.05	0.1	1.002485 - i0.001396
"	"	0.3	1.002247 - i0.001161
"	"	1.0	1.002290 - i0.001200
"	"	100.0	1.010513 - i0.000002
"	"	∞	1.010515 - i0.0
"	0.1	0.1	1.510056 - i0.518362
"	"	0.3	1.278652 - i0.338335
"	"	1.0	1.114107 - i0.150493
"	"	100.0	1.038768 - i0.0
"	"	∞	1.038768 - i0.0
3.0	0.05	0.1	1.001114 - i0.000399
"	"	0.3	1.000731 - i0.000197
"	"	1.0	1.000673 - i0.000170
"	"	100.0	1.002683 - i0.000026
"	"	∞	1.002731 - i0.0
"	0.1	0.1	1.003560 - i0.003103
"	"	0.3	1.002708 - i0.001680
"	"	1.0	1.002608 - i0.001471
"	"	100.0	1.010828 - i0.0
"	"	∞	1.010828 - i0.0

(i.e. Re(S) > 1). This is in contrast to the fast surface wave for the lossy earth without a ground screen. For a sufficiently small mesh size, the attenuation rate of the surface wave is considerably less than for the isolated half-space. However, for mesh sizes of 0.1λ or greater, the attenuation reduction is not significant and, in fact, the attenuation rate can even increase. For the rectangular mesh $(a \neq b)$, the propagation constant is highly dependent on the direction of propagation, and minimum attenuation is obtained for propagation along the more closely spaced wires.

Several extensions to this work would seem worthwhile. The approximate method of averaged boundary conditions which has been applied to single meshes [*Astrakhan*, 1968] and a pair of meshes [*Kontorovich et al.*, 1964] in free space could be extended to the lossy half-space geometry. *Casey* [1976] has recently applied this method to a square bonded mesh located at an air-dielectric interface. Since the unbonded mesh has superior reflecting properties [*Kontorovich et al.*, 1962; *Hill and Wait*, 1976], it could also be analyzed for surface wave propagation in the presence of a lossy half-space. Finally a thorough numerical search for all the modes could be made for the mesh-earth structure. We have examined the expected surface wave mode, but others may be possible. Solution of a source problem (such as vertical dipole excitation) would be useful in assessing the various pole (surface wave, etc.) and branch cut (continuous spectrum) contributions to the total field.

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