Scattering of Surface Waves at a Dielectric Discontinuity on a Planar Waveguide

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Abstract—**A theoretical treatment is presented of the scattering of a surface-wave mode on a planar surface waveguide at an abrupt axial transition. The latter is due to a dielectric obstacle that covers the line completely up to a given height. The analysis involves the matching of the tangential fields, expressed in terms of complete sets of eigenmodes, on the transition plane. The problem arises in certain obstacle detection schemes currently being proposed for guided transportation, which use the principle of guided radar.**

I. INTRODUCTION

USE of surface-wave lines has been suggested in the literature as a means of providing continuous-access communication for guided ground transportation [1], [2]. Obstacle detection schemes, sometimes referred to as 'guided radar,' have also been discussed (e.g., [2]–[5]).

In a guided radar scheme it is possible for an obstacle to be detected in one of two main ways: by the surface-wave reflections produced directly by the obstacle on the installed surface-wave line; or by the transmission loss of surface-wave signals sent back by transponders on the line at the far end of the guide-way. These types of operation are discussed in detail elsewhere [3]–[5].

In this paper we introduce a theoretical treatment of the scattering problem associated with obstacles that can be characterized as dielectrics, such as rockfalls, landslides, and snow. Typically, these obstacles cover a certain length of the transmission line and hence the incident wave will suffer both reflection loss at the front of the obstacle and transmission loss through the obstacle. Part of the incident power will be radiated in the forward and backward directions. All the scattering parameters at the obstacle are obviously functions of the relative permittivity ϵ_d, the loss tangent tan δ of the obstacle, as well as the general shape of the obstacle and the region of transition to it. As discussed in [4], [5], the transmission type of guided radar is probably more suitable for detecting these obstacles than the reflection type and hence the transmission loss in this case is of primary importance.

There are two extreme cases for the transition region. The first is an abrupt transition for which the reflection

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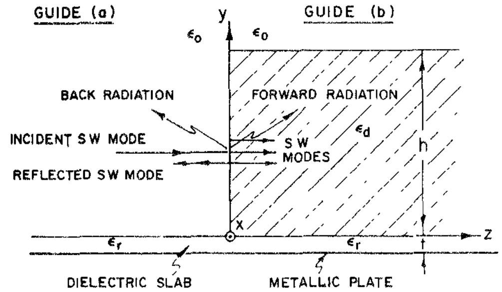

Fig. 1. Dielectric obstacle.

is a maximum. The second is a gradual transition which may be approximated by a matched load. The first case is obviously more favorable to the reflection type of operation of guided radar. As to the transmission type of operation, the transmission loss does not greatly change from one case to the other, as the loss through the obstacle is usually high. For this reason, we shall consider only the abrupt transition from guide a to guide b in Fig. 1, in which the problem has also been simplified to one in which the obstacle is represented as a uniform layer of dielectric on an infinite planar guide. This reduces it to a two-dimensional problem in order, at this stage, to emphasize the more fundamental aspects of guided radar. Guide b in the figure is thus a composite one that consists of the basic planar waveguide covered by the dielectric obstacle. It is assumed that guide a has only one propagating surface-wave mode incident on the transition. The method of solution of the scattering problem contains the following steps.

1) A complete set of eigenmodes, including both the discrete and continuous spectra, is derived in Section II for guides a and b in Fig. 1.

2) The total tangential fields on both sides of the plane $z = 0$ are expressed in terms of their respective sets of eigenmodes and the unknown scattering parameters. These fields are then equated on the whole plane $z = 0$ and the resulting equations are manipulated to derive the scattering parameters (Section III).

The method outlined is obviously applicable to a large class of problems that involve abrupt transitions from one waveguiding structure to another. It is worth mentioning that the present technique is, in essence, similar to that used by Clarricoats and Slinn [6] for transitions in closed waveguides, but is here adapted to cover open waveguides as well, where open waveguides are defined here as those for which the guided fields extend over an infinite cross section.

II. EIGENMODES

In this section we shall derive both the discrete and continuous spectra of eigenmodes for guide b (Fig. 1), where only TM modes will be considered. The corresponding spectra for guide a can then be obtained from the preceding as a special case in which the dielectric constant ϵ_d is replaced by that of the free space. This derivation follows closely the general outlines given by Shevchenko in his monograph [7], where the usual radiation condition is relaxed into a less stringent one in order to be able to define 'pseudomodes,' which constitute the continuous spectrum. In the following, all lengths will be normalized with respect to $\lambda_0/2\pi$ defined as unity and impedances normalized with respect to 120π. (Hence $\omega\epsilon_0$ and $\omega\mu_0$ will be replaced by unity, where $\lambda_0,\epsilon_0,\mu_0$ are the usual free-space values.)

For TM modes on guide b, the only nonvanishing fields are h_x, e_y, and e_z. We seek a solution characterized by a transverse wavenumber κ in the air region. The magnetic field component $h_x(\triangleq \psi)$ satisfies the appropriate wave equations in the following three regions. 1) $-t \leq y \leq 0$. 2) $0 \leq y \leq h$. 3) $y \geq h$. The solutions for ψ, with the factor $e^{j(\omega t-\beta z)}$ understood, are easily obtained as

$$\psi = \cos S\,(y+t), \qquad -t \leq y \leq 0 \quad (1a)$$

$$= a\cos ry + b\sin ry, \qquad 0 \leq y \leq h \quad (1b)$$

$$= V\exp\left(-j\kappa(y-h)\right) + W\exp\left(j\kappa(y-h)\right), \qquad y \geq h \quad (1c)$$

where

$$S = (\epsilon_r - \beta^2)^{1/2}, \qquad r = (\epsilon_d - \beta^2)^{1/2}, \qquad \text{and } \beta = (1-\kappa^2)^{1/2}$$

and ψ can represent either a discrete mode or a pseudomode, as will be defined later.

From Maxwell's equations, the electric field components e_y and e_z are given by

$$e_y = \beta\psi/\epsilon$$

and

$$e_z = j\psi'/\epsilon \quad (2)$$

where the prime indicates a differentiation with respect to y, ϵ is the relative dielectric constant of the region of interest, and β is the longitudinal wavenumber.

The continuity of h_x and e_z at the interfaces $y = 0$ and $y = h$ provides four equations whose solution results in the coefficients a, b, V, and W as follows.

$$a = \cos St \quad (3a)$$

$$b = -\,\epsilon_d S\sin St/(\epsilon_r r) \quad (3b)$$

and

$$\left.\begin{matrix} V \\ W \end{matrix}\right\} = \tfrac{1}{2}a(\cos rh \mp jr\sin rh/\epsilon_d\kappa) + \tfrac{1}{2}b(\sin rh \pm jr\cos rh/\epsilon_d\kappa). \quad (3c)$$

To complete the derivation of the modal spectra, we have to specify the condition that the fields should satisfy as $y \rightarrow \infty$. The radiation condition [8] could be imposed but it would result in only the discrete spectrum of modes, which is not a complete set. Instead, following Shevchenko [7], we relax the radiation condition into another that requires only the finiteness of the fields as $y \rightarrow \infty$, i.e.,

$$\lim_{y\rightarrow\infty} \psi = \text{finite quantity}. \quad (4)$$

The discrete modes have fields that decay in the air region away from the structure and hence they are given by [see (1c)]

$$W(\kappa) = 0 \text{ with } \operatorname{Im}(\kappa) < 0. \quad (5)$$

Equation (5) along with (3c) defines a discrete and finite set of modes which are referred to as surface-wave modes.

The continuous set of modes, which may be called 'pseudomodes' [7] corresponds to all purely real values of κ. Such modes each satisfy all boundary conditions, including condition (4) as $y \rightarrow \infty$. Although a pseudomode carries an infinite amount of power, a combination of such modes can represent a physical radiation field which indeed satisfies the usual radiation condition, as discussed by Shevchenko [7].

The complete set of eigenmodes satisfies the mutual orthogonality relationships. In particular, for two pseudomodes characterized by κ and κ', the following orthogonality relationship holds [5]

$$\int_{-t}^{\infty} e_y(\kappa)h_x(\kappa')\,dy = 2\pi V(\kappa)W(\kappa)\delta(\kappa-\kappa') \quad (6)$$

where $\delta(\cdot)$ is the Dirac δ-function.

III. THE TRANSITION REGION

A. Basic Equations

The modal spectrum of guide a includes a single surface-wave mode, the incident mode, plus a continuous spectrum of modes. Guide b has a finite number of surface-wave modes, the number being dependent on the obstacle height and the dielectric constant, as well as a continuous spectrum of modes. We assume that the surface-wave mode is incident alone from side a. As a result of scattering at the interface, there will be a reflected surface-wave mode plus a weighted sum of reflected pseudomodes (radiation) in guide a. Forward traveling surface-wave modes and pseudomodes will occur in guide b.

Let the transverse fields (e_y and h_x) belonging to surface-wave modes be denoted by

$$(e_1{}^a,h_1{}^a) \quad \text{and} \quad (e_i{}^b,h_i{}^b), \qquad i = 1,2,\cdots$$

and those belonging to the pseudomodes by

$$(e^a(\kappa),h^a(\kappa)) \quad \text{and} \quad (e^b(\kappa),h^b(\kappa)), \qquad 0 \leq \kappa \leq \infty$$

where the superscripts a and b refer to guides a and b, respectively. Now the continuity of the tangential electric and magnetic fields at the plane $z = 0$ becomes

$$(1 + R)e_1^a + \int_0^\infty d\kappa\Gamma(\kappa)e^a(\kappa) = \sum_{i=1}^{M_s} T_i e_i^b + \int_0^\infty d\kappa T(\kappa)e^b(\kappa)$$

$$(1 - R)h_1^a - \int_0^\infty d\kappa\Gamma(\kappa)h^a(\kappa) = \sum_{i=1}^{M_s} T_i h_i^b + \int_0^\infty d\kappa T(\kappa)h^b(\kappa) \quad (7)$$

where R and T_i are the unknown surface-wave reflection and transmission coefficients, and $\Gamma(\kappa)$ and $T(\kappa)$ are the unknown reflection and transmission coefficients for the pseudomodes. M_s is the number of surface-wave modes in guide b. We notice that all the fields are functions of the transverse dimension, y (over the interval $-t \leq y \leq \infty$), which is dropped here for convenience.

The first step to simplify (7) is to change the continuous summations into discrete ones. This is done by writing the functions $\Gamma(\kappa)$ and $T(\kappa)$ in terms of a known set of functions $f_i(\kappa)$, e.g., normalized Laguerre polynomials, which is complete over the interval $(0,\infty)$. Thus

$$\Gamma(\kappa) = \sum_{j=0}^{\infty} \gamma_j f_j(\kappa)$$

and

$$T(\kappa) = \sum_{j=0}^{\infty} t_j f_j(\kappa). \quad (8)$$

Upon substitution back in (7), these take the form

$$(1 + R)e_1^a + \sum_{j=0}^{\infty} \gamma_j E_j^a = \sum_{i=1}^{M_s} T_i e_i^b + \sum_{j=0}^{\infty} t_j E_j^b$$

$$(1 - R)h_1^a - \sum_{j=0}^{\infty} \gamma_j H_j^a = \sum_{i=1}^{M_s} T_i h_i^b + \sum_{j=0}^{\infty} t_j H_j^b \quad (9)$$

where the fields $\{E_j^{a,b},H_j^{a,b}\}$ are defined as

$$\{E_j^{a,b},H_j^{a,b}\} \triangleq \int_0^\infty d\kappa f_j(\kappa)\{e^{a,b}(\kappa),h^{a,b}(\kappa)\} \qquad j = 0,1,2,\cdots. \quad (10)$$

Since the functions $f_j(\kappa)$ are known and the basic pseudomode functions are also known for a given structure, the preceding set of fields is then completely defined. The field $E_j^{a,b}(\kappa)$ and $H_j^{a,b}(\kappa)$ can be considered basic eigenfunctions that represent the continuous spectra of the guides a and b, respectively. Notice, however, that these eigenfunctions are not necessarily orthogonal. Nevertheless, this should not limit their role as a set of basic eigenfunctions.

One attractive feature of the form of (9) is that they resemble the corresponding equations that would have arisen if our guiding structure were a closed instead of an open waveguide. In the closed waveguide case, the spectrum of modes is entirely discrete and so is the case in (9). This, then, suggests that the procedure of solution used by Clarricoats and Slinn [6] for a closed waveguide can as well be used in the present problem.

To eliminate the dependence on the transverse dimension y in (9) (which is not explicitly displayed) we multiply the first equation by $h_i^b, i = 1,\cdots,M_s$ and H_m^b, $m = 0,1,\cdots$, successively, and then integrate over the range of $y(-t \to \infty)$. We do the same with the second equation except that we multiply by $e_i^b, i = 1,\cdots,M_s$ and $E_m^b, m = 0,1,\cdots$, instead. Before writing the resulting equations, let us define the following inner product

$$\langle v_1,v_2\rangle = \int_{-t}^{\infty} v_1 v_2 \, dy$$

where v_1 and v_2 are functions of y, defined over the interval of integration. We then have the following infinite set of equations

$$(1 + R)\langle e_1^a,h_i^b\rangle + \sum_{j=0}^{\infty} \gamma_j\langle E_j^a,h_i^b\rangle = T_i \qquad i = 1,2,\cdots,M_s$$

$$(1 + R)\langle e_1^a,H_m^b\rangle + \sum_{j=0}^{\infty} \gamma_j\langle E_j^a,H_m^b\rangle = \sum_{j=0}^{\infty} t_j\langle E_j^b,H_m^b\rangle \qquad m = 0,1,\cdots,\infty$$

$$(1 - R)\langle e_i^b,h_1^a\rangle - \sum_{j=0}^{\infty} \gamma_j\langle e_i^b,H_j^a\rangle = T_i \qquad i = 1,2,\cdots,M_s$$

$$(1 - R)\langle E_m^b,h_1^a\rangle - \sum_{j=0}^{\infty} \gamma_j\langle E_m^b,H_j^a\rangle = \sum_{j=0}^{\infty} t_j\langle E_m^b,H_j^b\rangle \qquad m = 0,1,\cdots,\infty \quad (11)$$

where we have used the fact that the modes (e_i^b,h_i^b) are orthogonal to each other and to the modes (E_j^b,H_j^b). The inner product $\langle e_i^b,h_i^b\rangle$ is assumed to be normalized to unity.

An alternative derivation of (11) which may give more insight into these equations can be obtained as follows. The total tangential fields at $z = 0$ on the side of guide a may be expressed in terms of forward and backward traveling modes which belong to guide b. The forward traveling modes are then equated to those on the b side of the boundary while the backward traveling modes are equated to zero. This approach is worked out in detail elsewhere [5] and is shown to result in the same set of equations as (11).

To solve the preceding equations, it is necessary to truncate them. Therefore we limit the number of modes belonging to the continuous spectrum to N in guide a and M in guide b where N and M are finite integers. It will be seen later that only a small number of these modes, about two or three, need to be considered. The functions $f_i(\kappa)$ have here been chosen to be the normalized Laguerre polynomials (e.g., [8]). The first few members of these functions are as follows.

$$f_0(\kappa) = \exp(-\kappa/2)$$
$$f_1(\kappa) = (1-\kappa)\exp(-\kappa/2)$$
$$f_2(\kappa) = (1-2\kappa+\kappa^2/2)\exp(-\kappa/2), \qquad \text{etc.}, \cdots \tag{12}$$

which show that as the order of the function increases, its decay with κ becomes slower. Hence the part of the continuous spectrum corresponding to high values of κ, i.e., the evanescent modes with β imaginary, will be effectively represented by the higher order Laguerre functions. It also follows that the truncation of this set of functions amounts to the neglect of the contribution of the higher evanescent modes, an approximation which is common practice in dealing with closed as well as open waveguides.

After the truncation, the number of unknowns in (11) becomes $1 + N + M_s + M$, while the number of equations are $2(M_s + M)$. We can always adjust N and M such that the number of equations is equal to or greater than the number of unknowns. In the latter case, a pseudoinverse of the matrix of coefficients is performed.

A comment on the power balance is now in order. A unit power incident from guide a should be accounted for by the scattered power in guides a and b. We notice that (11), in their truncated form, do not automatically guarantee the satisfaction of the power balance condition. This condition then serves as a good check on the obtained solution. Thus we change N and M within certain limits and we choose that solution which best satisfies the power balance, which helps to reject possible unphysical solutions. This completes the basic formulation of the problem.

B. Derivation of the Inner Products in (11)

The inner products in (11) involve the following basic expressions.

$$\int_{-t}^{\infty} e^a(\kappa^a)h^b(\kappa^b)\,dy \quad \text{and} \quad \int_{-t}^{\infty} e^b(\kappa^b)h^a(\kappa^a)\,dy$$

where κ^a and κ^b are the transverse wavenumbers in the air region. κ^a and κ^b can each be either purely real or purely imaginary, depending on whether the mode considered is a pseudomode or a surface-wave mode. In the following we shall derive the first integral in the preceding in terms of the basic scalar fields $\psi^a(\kappa^a)$ and $\psi^b(\kappa^b)$. The second integral can be obtained by a mere exchange of the superscripts a and b.

Using (1) and (2) we have

$$\int_{-t}^{\infty} e^a(\kappa^a)h^b(\kappa^b)\,dy = \beta^a \int_{-t}^{\infty} (\psi^a/\epsilon^a)\psi^b\,dy \tag{13}$$

where, for convenience, the arguments κ^a and κ^b are omitted on the right-hand side (RHS). ϵ^a and ϵ^b are the relative permittivities in guides a and b, respectively, and are given by

$$\begin{aligned} \epsilon^a &= \epsilon_r, & -t \le y \le 0 \\ &= 1, & y \ge 0 \end{aligned} \tag{14a}$$

and

$$\begin{aligned} \epsilon^b &= \epsilon_r, & -t \le y \le 0 \\ &= \epsilon_d, & 0 \le y \le h \\ &= 1, & y \ge h. \end{aligned} \tag{14b}$$

The following basic differential equations apply for ψ^a and ψ^b.

$$(d^2/dy^2)\psi = -S^2\psi, \text{ region 1} \tag{15a}$$
$$= -r^2\psi, \text{ region 2} \tag{15b}$$
$$= -\kappa^2\psi, \text{ region 3} \tag{15c}$$

where S, r, and κ are transverse wavenumbers in the respective regions. We notice that r is different in the two guides, being $(\epsilon_d - \beta^2)^{1/2}$ in guide b and κ^a in guide a. On writing (15a) for both ψ^a and ψ^b, multiplying the resulting two equations by ψ^b and ψ^a, respectively, and then subtracting the second from the first, we obtain the following.

$$\int_{-t}^{0^-} \psi^a\psi^b\,dy = \frac{1}{S^{b^2} - S^{a^2}}\left[\psi^{a\prime}\psi^b - \psi^a\psi^{b\prime}\right]_{-t}^{0^-} \tag{16}$$

and similar expressions hold for integrations over regions 2 and 3. The prime over ψ refers to a differentiation with respect to y. The boundary conditions at $y = -t$, 0, and h are

$$\psi'(-t) = 0$$
$$\psi(0^-) = \psi(0^+)$$
$$\psi'/\epsilon\,|_{0^-} = \psi'/\epsilon\,|_{0^+}$$
$$\psi(h^-) = \psi(h^+)$$

and

$$\psi'/\epsilon\,|_{h^-} = \psi'/\epsilon\,|_{h^+} \tag{17}$$

which apply to both ψ^a and ψ^b.

Now using (16) and similar expressions for regions 2 and 3 along with the boundary conditions (17), we finally arrive, after some manipulations, at an expression for (13) as

$$\int_{-t}^{\infty} e^a(\kappa^a)h^b(\kappa^b)\,dy = 2\pi\beta^a \operatorname{Re}(V^aW^b)\delta(\kappa^a - \kappa^b) + \frac{\beta^a(\epsilon_d - 1)}{r^{b^2} - r^{a^2}}\left[\frac{X_1 - X_2}{\kappa^{b^2} - \kappa^{a^2}} + X_3/\epsilon_d\right] \tag{18}$$

where X_1, X_2, and X_3 are given by

$$X_1 = (1/\epsilon_r)\left[\psi^{a\prime}\psi^b - \psi^a\psi^{b\prime}\right]_{y=0^-}$$
$$X_2 = \left[\psi^{a\prime}\psi^b - \psi^a\psi^{b\prime}\right]_{y=h^+}$$

and

$$X_3 = \left[\psi^a\psi^{b\prime}\right]_{y=t^+} - \left[\psi^a\psi^{b\prime}\right]_{y=h^-}. \tag{19}$$

The inner product with a and b exchanged is also given by (18) if we interchange the superscripts a and b and change X_3/ϵ_d to X_3 and the term $(\epsilon_d - 1)$ into $(1 - \epsilon_d)$.

IV. NUMERICAL RESULTS AND COMMENTS

The truncated set of (11) was solved numerically for the complex scattering parameters $R, T_i, i = 1,2,\cdots,M_s$, $\gamma_i, i = 1,2,\cdots,N$, and $t_i, i = 1,2,\cdots,M$. Some of the inner products in these equations were computed by use of a Laguerre–Gaussian integration subroutine. The integers N and M were increased until the variations in the scattering coefficients became sufficiently small. Values of N and M up to 3 were found adequate for good convergence and in all cases we used $N = 1$ or 2 and $M = 3$. The power balance equation was checked at every choice of N and M and if this was in error by more than 3 percent, that solution was completely rejected.

Fig. 2 shows results for the magnitude of the surface-wave reflection coefficient. Three cases corresponding to different line parameters were treated as follows.

Case 1: $t/\lambda_0 = 0.06$, $\epsilon_r = 2.56$, $v_p/C = 0.972$, $P_a/P_T = 0.925$.

Case 2: $t/\lambda_0 = 0.06$, $\epsilon_r = 9.00$, $v_p/C = 0.882$, $P_a/P_T = 0.888$.

Case 3: $t/\lambda_0 = 0.147$, $\epsilon_r = 2.56$, $v_p/C = 0.837$, $P_a/P_T = 0.473$.

In the preceding, v_p/C is the ratio of the phase velocity of the incident surface-wave mode to that of a plane wave in free space, and P_a/P_T is the proportion of the total incident power in the air region of guide a. The dielectric constant ϵ_d of the obstacle was taken equal to 10 in all cases; a figure that is typical for rocks.

The magnitude of the reflection coefficient $|R|$ tends to increase with the height h of the transition in an oscillatory manner (Fig. 2). These oscillations may be attributed to the fact that the phase velocities of the surface-wave modes in guide b, as well as the number of modes, are continuously varying with the height of the transition. Whenever these modes can make a good match to the incident surface-wave mode from a, the reflection coefficient shows a minimum. Another observation on Fig. 2 is that at sufficiently high h, the reflection coefficient is greater for a greater ratio P_a/P_T, the latter being greatest for Case 1 and lowest for Case 3. Finally, we observe that the first peak of $|R|$ occurs at a lower height h in Case 2 compared with case 1). This may be explained by the fact that the evanescent decay of the incident surface-wave fields in the air is greater in Case 2. It is worth mentioning that the phase of R (relative to that of the incident mode at $z = 0$) was hardly different from 180° in all cases except at very low height (as low as $0.03\lambda_0$) where it deviated by only a few degrees. It may then be suggested that this can always be taken as 180° for the particular parameters of the line and obstacle treated here.

Fig. 3 gives the percentages of the transmitted power in guide b, both in the form of surface waves and radiation.

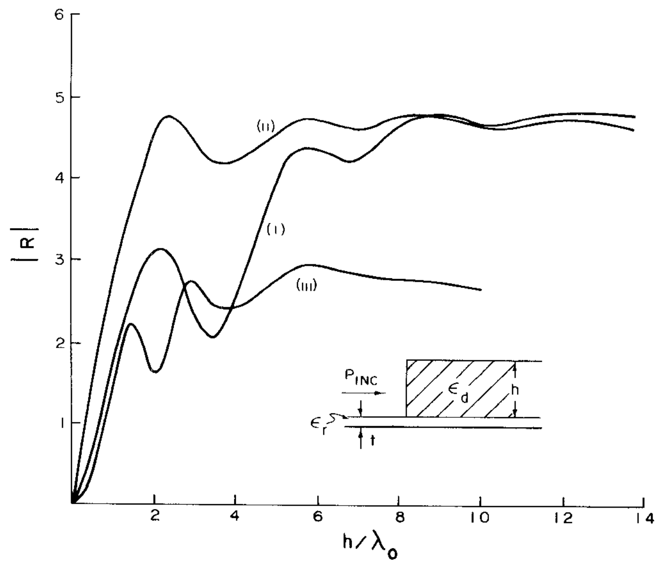

Fig. 2. Reflection coefficient. Case 1: $\epsilon_r = 2.56$, $t/\lambda_0 = 0.06$. Case 2: $\epsilon_r = 9.00$, $t/\lambda_0 = 0.06$. Case 3: $\epsilon_r = 2.56$, $t/\lambda_0 = 0.147$.

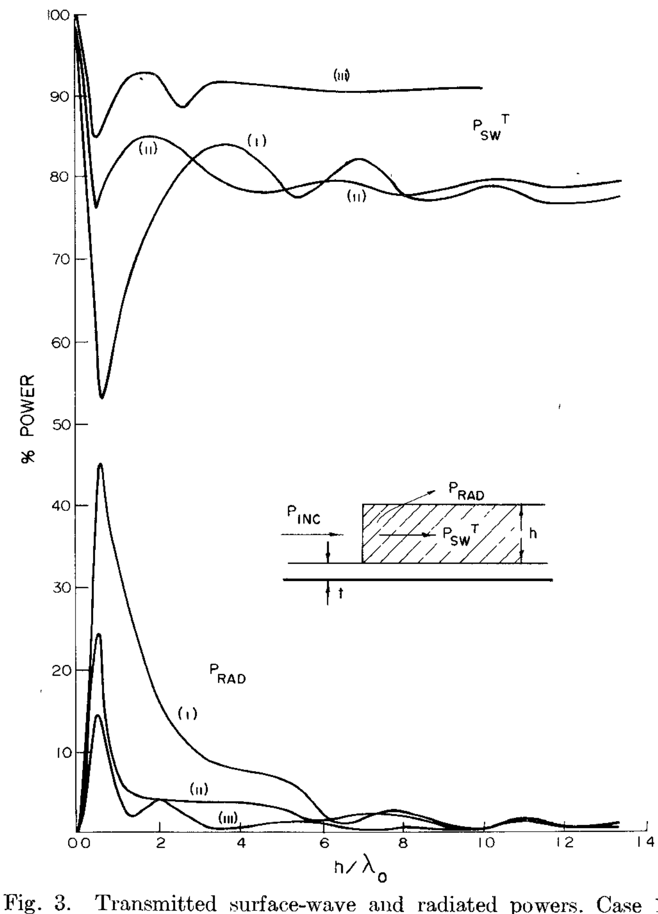

Fig. 3. Transmitted surface-wave and radiated powers. Case 1: $\epsilon_r = 2.56$, $t/\lambda_0 = 0.06$. Case 2: $\epsilon_r = 9.00$, $t/\lambda_0 = 0.06$. Case 3: $\epsilon_r = 2.56$, $t/\lambda_0 = 0.147$.

As the transition height $h \to 0$, it is clear that the transition surface-wave power approaches 100 percent of that of the incident wave, with the radiation being zero. On the other hand, for a very high obstacle, the transmitted power is less than 100 percent but it is mostly trapped as surface-wave modes in guide b, with the radiation ap-

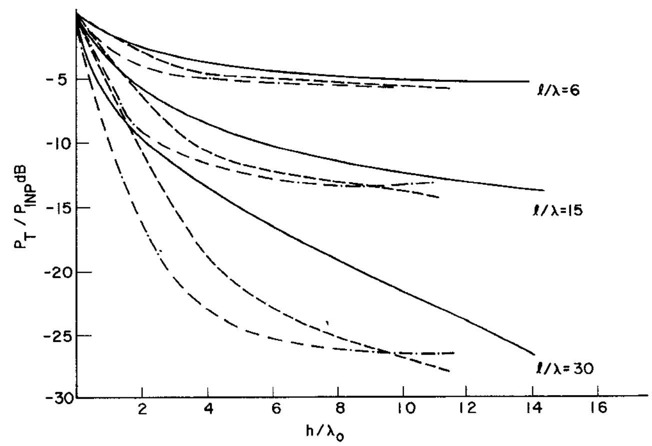

Fig. 4. Transmission loss versus obstacle height. Case 1:—— $t/\lambda = 0.06$, $\epsilon_r = 2.56$, $v_p/c = 0.972$. Case 2: –•–, $t/\lambda_0 = 0.06$, $\epsilon_r = 9.00$, $v_p/c = 0.882$. Case 3: – – –, $t/\lambda_0 = 0.147$, $\epsilon_r = 2.56$, $v_p/c = 0.837$.

proaching zero again. Thus there is some height h at which the radiation component of power is a maximum and, correspondingly, the surface-wave part is a minimum. In all the three cases shown in Fig. 3 this height corresponds to only one surface-wave mode propagating in guide b. We further observe that the peak of the radiated power is less for a slower incident mode and hence it is least for Case 3. For all the cases considered here, we have found that the reflected radiation power in guide a was negligibly small. However, no general conclusions regarding this point will be made since this may be due to the particular parameters chosen for the present problem.

In Fig. 4 we show the power transmission loss as a function of h at various $z = l$ planes in guide b. The attenuation factors for the modes were computed by simple perturbation. Such computation can produce meaningful results only if the loss tangent ($\tan \delta$) of the obstacle material is taken to be very small. We have thus fixed $\tan \delta$ at 0.01, which still lies within the range of the loss tangents of typical rocks and snow at frequencies in the UHF range. An important consequence of Fig. 4 is that a significant discrimination can be observed between a dangerously high obstacle and an insignificantly low one, as there exists a large difference in their transmission losses over a distance of, say, 30 free-space wavelengths. This is an important consideration for the transmission type of guided radar [4] where, for example, a high pile of rocks would have to be distinguished from a small layer of dirt on the line.

V. CONCLUSIONS

The problem treated here belongs to the general class of problems dealing with abrupt axial discontinuities in open, or surface waveguides. The total fields on both sides of the discontinuity are expressed in terms of appropriate complete sets of eigenmodes, that include both the discrete and the continuous spectra, and the boundary conditions at the discontinuity plane are enforced. The continuous summation over pseudomodes is simplified by expanding it as a discrete sum that involves the use of normalized Laguerre polynomials.

The surface-waveguiding structure considered is that of a simple plane dielectric layer on a metallic ground sheet and the discontinuity is due to a finite height of a dielectric material that completely covers the guide. The surface-wave reflection and transmission coefficients, as well as the radiated power, are obtained for a variety of line parameters and heights of the obstacle.

This problem was motivated by the development of obstacle detection schemes, using guided radar to detect landslides, rocks, or snow on the tracks of guided transportation systems.

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