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Scattering by inhomogeneities in parallel plate waveguides

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Abstract—Electromagnetic scattering by a bounded object located inside a parallel plate waveguide is the subject of this paper. The exciting field in the waveguide is either an arbitrary source located at a finite distance from the obstacle or a plane wave generated in the far zone. The analytic treatment of the problem relies on an extension of the null field approach or T-matrix method.

I. INTRODUCTION

Recent theoretical progress in the development of useful scattering identities — sum rules [1] — have initiated several attempts to verify these identities experimentally, see e.g., [2]. These sum rules relate the dynamical behavior of the scattering and absorption behavior of the scatterer to the static properties of the scatterer (polarizability dyadics).

Initial investigations show that the parallel plate waveguide is accessible [2,3]. A detailed investigation of the static properties of an obstacle between two parallel plates has also been reported recently [4].

The approach employs the integral representation of the solution. This integral representation approach to solve the scattering problem was originally introduced by Peter Waterman for finite scatterers, and it has proven to be a very powerful and useful technique to solve a large variety of scattering problems, not only electromagnetic, but also acoustic and elastodynamic problems. In fact, the present geometry is an extension of the results with buried obstacle close to a planar interface — layered or not, see e.g., [5,6].

II. FORMULATION OF THE PROBLEM

A finite scatterer with bounding surface $S_s$ defines the region $V_s$. Two infinite, perfectly conducting planes, $S_+$ and $S_−$, confine the regions $V_c$ and $V_0$, see Figure 1. These planes are parameterized by $z = z_+$ and $z = z_−$, respectively. The regions above $S_+$ and below $S_−$ are denoted by $V_c$ and $V_0$, respectively. The sources of the problem are assumed to be located in $V_i ⊂ V_c$, between the surfaces $S_+$ and $S_−$.

To proceed, the time-harmonic electric and magnetic fields satisfy the free-space Maxwell equations in $V_c$ (we use the time convention $\exp(-iωt)$),

\[
\begin{align*}
\nabla \times \mathbf{E}(r) &= ik_0\eta_0\mathbf{H}(r), \\
\nabla \times \mathbf{H}(r) &= -i k_0 \mathbf{E}(r)
\end{align*}
\]

where $k_0 = ω/ε_0$ and $η_0$ are the wave number and wave impedance in free space, respectively. The boundary conditions on the bounding surfaces are

\[
\begin{align*}
\mathbf{ν} \times \mathbf{E}(\mathbf{r}) &= 0, & r &\in S_+ \cup S_-
\end{align*}
\]

$\mathbf{ν}$ is the field with no obstacle or plates present. With the directions of the unit normals defined as in Figure 1, the solution of (1) and (2) satisfies the surface integral representation [7]

\[
\frac{1}{i k_0} \nabla \times \left( \nabla \times \int \mathbf{G}_c(k_0, |\mathbf{r} - \mathbf{r}'|) \cdot \mathbf{K}(\mathbf{r}') \, dS' \right) \mathbf{ν} \times \mathbf{E}(\mathbf{r}), \quad r \in V_c
\]

where $\mathbf{K} = \mathbf{ν} \times \mathbf{H}$, and the electric Green’s dyadic

\[
\mathbf{G}_c(k_0, |\mathbf{r} - \mathbf{r}'|) = \left( \mathbf{I}_3 + \frac{1}{k_0^2} \nabla \nabla \right) \frac{e^{ik_0|\mathbf{r} - \mathbf{r}'|}}{4π|\mathbf{r} - \mathbf{r}'|}
\]
The integral representation also contains a surface integral evaluated at large lateral distances, but proper radiation conditions at large lateral distances make this integral vanish. This surface integral representation is the starting point in the null-field approach.

III. BASIS FUNCTIONS AND EXPANSIONS

A. Spherical and planar vector waves

To proceed, out-going or radiating spherical vector waves, \( u_{rrm}(kr) = u_{rn}(kr) \), and regular spherical vector waves, \( v_{rrm}(kr) \), are employed. We adopt the definition in [8, 9].

The most appropriate basis functions to deal with the geometry of the parallel plates are planar waves. We define the dimension-less, vector-valued plane waves, \( \varphi_j^\pm (k_t r) \), \( j = 1, 2 \), as:

\[
\begin{align*}
\varphi_1^+(k_t r) &= \frac{\hat{z} \times k_t}{4\pi k_t} e^{i k_t r}, \\
\varphi_2^+(k_t r) &= \frac{\frac{1}{r} k_t k_z + \frac{1}{r} \hat{z} \times k_t}{4\pi k_t} e^{i k_t r},
\end{align*}
\]

where the transverse (tangential) wave vector and the spatial position vector in the plane are

\[ k_t = \hat{x} k_x + \hat{y} k_y, \quad k_x, k_y \in \mathbb{R}, \quad \rho = x \hat{x} + y \hat{y} \]

The length of the transverse wave vector is always a real number, \( \rho z \)., and \( k_z \) is defined by

\[ k_z = \left(k_0^2 - k_t^2\right)^{1/2} = \begin{cases} \\
\sqrt{k_0^2 - k_t^2} & \text{for } k_t < k_0 \\
\sqrt{k_0^2 - k_t^2} & \text{for } k_t > k_0
\end{cases} \]

The plus super-index denotes an exponentially decreasing or evanescent wave as \( z \to \infty \), and similarly for the minus super index as \( z \to -\infty \). The index \( j = 1 \) labels the TE-waves, and \( j = 2 \) labels the TM-waves.

B. Green’s dyadic decompositions

The Green’s dyadic is decomposed in spherical vector waves [8]

\[
G_e(k_0, |r - r'|) = i k_0 \sum_n v_n(k_0 r) u_n(k_0 r') = i k_0 \sum_n u_n(k_0 r) v_n(k_0 r') \quad (4)
\]

where \( r_\prec (r_\succ) \) is the position vector with the smallest (largest) distance to the origin, i.e., if \( r < r' \) then \( r_\prec = r \) and \( r_\succ = r' \). The summation is over the divergence-free vector spherical vector waves, \( \tau = 1, 2 \). Moreover, we need the decomposition of the Green’s dyadic in planar vector waves [8]

\[
G_e(k_0, |r - r'|) = 2i k_0 \sum_{j=1,2} \iint_{\mathbb{R}^2} \varphi_j^+(k_t r) \varphi_j^+(k_t r') \frac{dk_z dk_y}{k_z k_0^2} \quad (5)
\]

where the upper (lower) sign holds for \( V_\tau \) (\( V_\rho \)), and where the dagger corresponds to a change \( k_t \to -k_t \).

C. Transformation between solutions

To connect the spherical vector waves and the planar vector waves, we need a transformation between the two sets of solutions. The result that is relevant in the analysis below are [8, p. 183]

\[
u_n(k_0 r) = 2 \sum_{j=1,2} \int_{\mathbb{R}^2} B_{nj}^\mp (k_t) \varphi_j^\pm (k_t r) \frac{k_0 dk_x dk_y}{k_z k_0^2}, \quad z \geq 0 \quad (6)
\]

where

\[
B_{nj}^\pm (k_t) = i^{\tau - r} C_{\tau m}(\pm 1)^{1+m} \left\{ \mp \delta_{\tau j} \Delta^m_{\tau l}(k_z/k_0) \left\{ \cos m\beta \sin m\beta \right\} \right. \\
\left. - \delta_{\tau j} \pi^m_{\tau l}(k_z/k_0) \left\{ -\sin m\beta \cos m\beta \right\} \right\} \quad (7)
\]

and where \( k_t = k_0 (\hat{x} \cos \beta + \hat{y} \sin \beta) \), \( C_{\tau m} \) are normalization constants [9], and

\[
\Delta^m_{\tau l}(t) = \frac{(1 - t^2)^{1/2}}{\sqrt{l(l+1)}} \quad (8)
\]

IV. INCIDENT ELECTRIC FIELD

The sources of the incident field are assumed to be located between the plates \( S_\pm \) and outside the scatterer \( V_e \). Due to the completeness of the planar vector waves and the spherical vector waves, the incident electric field is assumed to have the following expansions in the three regions, \( V_\pm \) and \( V_e \): \n
\[
E_i(r) = \sum_{j=1,2} \int_{\mathbb{R}^2} a_j^\pm (k_t) \varphi_j^\pm (k_t r) \frac{dk_z dk_y}{k_0 k_z}, \quad r \in V_\pm \quad (9)
\]

where the upper (lower) sign holds for \( V_+ \) (\( V_- \)), and

\[
E_i(r) = \sum_n a_n v_n(k_0 r), \quad r \in V_R \quad (10)
\]

where \( V_R \) is a sphere of radius \( R \) that does not include the circumscribing sphere of the source region. The coefficients \( a_j^\pm (k_t) \) and \( a_n \) are assumed known.

V. UTILIZING THE SURFACE INTEGRAL REPRESENTATION

The position vector, \( r \), can take four different principle positions, \( r \in V_\pm, r \in V_e, \) and \( r \in V_\rho \). We now explore these possibilities. The decompositions of the Green’s dyadic in spherical and planar vector waves, see (4) and (5) are now used.

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When the position vector is either in $V_+$ or in $V_-$, the lower line of (3) using (4), (5), (6), and (9) yield, since the planar vector waves are linearly independent

$$a_j^\pm(k_i) = 2k_0^2 \frac{k_0}{k_z} \int_{S_+ \cup S_-} \varphi_j^{\pm\dagger}(k_i; r') \cdot K(r') \, dS'$$

$$+ 2k_0^2 \frac{k_0}{k_z} \sum_n B_{nj}^\pm(k_i) \int_{S_n} v_n(k_0 r') \cdot K(r') \, dS'$$

(11)

When the position vector is inside the largest sphere enclosed in $V_+$, the lower line of (3) using (4), (6), and (9) yield

$$a_n = k_0^2 \int_{S_n} u_n(k_0 r') \cdot K(r') \, dS'$$

$$+ 2k_0^3 \sum_{j=1,2} \int_{\mathbb{R}^2} B_{nj}^+(k_i) \int_{S_n} \varphi_j^+(k_i; r') \cdot K(r') \, dS' \frac{dk_x \, dk_y}{k_z^2 k_0^2}$$

$$+ 2k_0^3 \sum_{j=1,2} \int_{\mathbb{R}^2} B_{nj}^-(k_i) \int_{S_n} \varphi_j^-(k_i; r') \cdot K(r') \, dS' \frac{dk_x \, dk_y}{k_z^2 k_0^2}$$

(12)

since the regular spherical vector waves, $v_n(k_0 r')$, are linearly independent.

When the position vector is outside the circumscribing sphere of $S_n$ in $V_+$, the upper line of (3) using (4), (5), (6), and (9) yield

$$E_s(r) = \sum_n f_n u_n(k_0 r')$$

$$+ \sum_{j=1,2} \int_{\mathbb{R}^2} f_j^+(k_i) \varphi_j^+(k_i; r) \frac{dk_x \, dk_y}{k_0^2}$$

$$+ \sum_{j=1,2} \int_{\mathbb{R}^2} f_j^-(k_i) \varphi_j^-(k_i; r) \frac{dk_x \, dk_y}{k_0^2}$$

(13)

where

$$f_n = -k_0^2 \int_{S_n} v_n(k_0 r') \cdot K(r') \, dS'$$

$$f_j^\pm(k_i) = -2k_0^3 \frac{k_0}{k_z} \int_{S_\pm} \varphi_j^{\pm\dagger}(k_i; r') \cdot K(r') \, dS'$$

(14)

VI. EXPANSION AND ELIMINATION OF THE SURFACE FIELDS

Expand the currents on the surfaces in planar vector waves and a complete set of tangential vector functions, $\mathbf{\hat{v}} \times \psi_n$, on $S_n$. We assume

$$K(r) = \sum_{j=1,2} \int_{\mathbb{R}^2} \alpha_j^\pm(k_i) \mathbf{e}_z \times \varphi_j^\pm(k_i; r') \frac{dk_x \, dk_y}{k_0^2}, \quad r \in S_\pm$$

and

$$K(r) = \sum_n \alpha_n \mathbf{\hat{v}}(r) \times \psi_n(r), \quad r \in S_n$$

where the dual index $\overline{j}$ is defined $\overline{1} = 2$ and $\overline{2} = 1$.

Insert these expansions in (11) and (12), use the orthogonality of the planar vector waves and solve for the unknown coefficients $\alpha_j^\pm(k_i)$ and $\alpha_n$. The result is

$$\alpha_j^\pm(k_i) = \mp \frac{2i}{4k_0} a_j^\pm(k_i) + a_j^\pm(k_i)(-1)^j e^{\pm 2ik_z z}$$

$$\times \frac{B_{nj}^\pm(k_i) + B_{nj}^\mp(k_i)(-1)^j e^{\pm 2ik_z z}}{1 - e^{-2ik_z d}}$$

(15)

where the distance between the plates is $d = z_+ - z_-$, the T-matrix, $T_{nn'}$, of the scatterer is defined in [10], and

$$\gamma_n = \frac{1}{2i} \sum_{j=1,2} \int_{\mathbb{R}^2} \left( \alpha_j^+(k_i) B_{nj}^{-1}(k_i) \right) \frac{dk_x \, dk_y}{k_0^2}$$

$$- \left( \alpha_j^-(k_i) B_{nj}^+(k_i) \right) \frac{dk_x \, dk_y}{k_0^2}$$

(16)

Insert the formulas of $\alpha_j^\pm(k_i)$ from above in the expression for $\gamma_n$ in (16), and we obtain an infinite set of equations that can be solved for every specified incident field. We write as

$$c_n = d_n + \sum_{n'n'} A_{nn'} T_{nn'} c_{n'n'}$$

(17)

where the array $c_n = a_n + \gamma_n$, and the $d_n$ vector is defined as

$$d_n = - \sum_{j=1,2} \int_{\mathbb{R}^2} \alpha_j^+(k_i)$$

$$\times B_{nj}^+(k_i) + (-1)^j e^{2ik_z z} B_{nj}^{-1}(k_i) \frac{dk_x \, dk_y}{k_0^2}$$

$$- \sum_{j=1,2} \int_{\mathbb{R}^2} a_j^+(k_i)$$

$$\times B_{nj}^{-1}(k_i) + (-1)^j e^{2ik_z z} B_{nj}^+(k_i) \frac{dk_x \, dk_y}{k_0^2} + a_n$$

(18)

and the $A$ matrix is

$$A_{nn'} = -2 \sum_{j=1,2} \int_{\mathbb{R}^2} B_{nj}^{-1}(k_i)$$

$$\times B_{n'j}^+(k_i) + (-1)^j e^{2ik_z z} B_{n'j}^{-1}(k_i) k_0 \frac{dk_x \, dk_y}{k^2}$$

$$- 2 \sum_{j=1,2} \int_{\mathbb{R}^2} B_{nj}^+(k_i)$$

$$\times B_{n'j}^{-1}(k_i) + (-1)^j e^{-2ik_z z} B_{n'j}^+(k_i) k_0 \frac{dk_x \, dk_y}{k^2}$$

(19)

The $A$ matrix is independent of the excitation and the scatterer, and the entries can be computed once and for all and stored for later use.
Finally, we get the expansion coefficients of the scattered field as, see (14)

\[
\begin{align*}
  f_n & = \sum_{n'} T_{\text{int}} c_{n'}', \\
  f_{n}^\pm (k_i) & = \pm \frac{1}{2i} \alpha_j^\pm (k_i)
\end{align*}
\]  

VII. THE PRIMARY AND SECONDARY FIELDS

The total electric field \( E(r) \) between the plates can be decomposed in several ways. Above, we decomposed the total field as \( E(r) = E_i(r) + E_s(r) \). As an alternative to this decomposition of the electric field, we introduce a decomposition in terms of a primary and secondary field, i.e.,

\[ E(r) = E_i(r) + E_s(r) = E_{\text{prim}}^r(r) + E_{\text{sec}}^r(r) \]

The explicit expression of these fields are found in Ref. [9]. The field \( E_{\text{prim}}^r(r) \) is the total electric field in the absence of the scatterer (surfaces \( S_+ \) and \( S_- \) present), and \( E_{\text{sec}}^r(r) \) is the correction due to the presence of the scatterer. All integrals that appear in the calculations can be evaluated exactly by calculus of residues [9], except for the entries of the \( A \) matrix, which can be evaluated exactly in terms of the Lerch function [11]

\[ \Phi(\beta, \nu, \mu) = \sum_{n=0}^{\infty} \frac{\beta^n}{(\mu + n)!} \]

VIII. NUMERICAL EXAMPLES

The first propagating mode is taken as the exciting field, which is a vertically polarized plane wave. In Figure 2 the scattering cross section of a perfectly conducting sphere of radius \( a \) is depicted as a function of \( k_0d \). The explicit data of the example are given in the caption.

IX. CONCLUSIONS

By the use of the integral representation of the scattered field, the solution to the complex electromagnetic scattering problem has been solved. The solution is an extension of the null field method, originally proposed by Peter Waterman, to geometries with two planar interfaces. Similar geometries have been addressed in the past, see e.g., [5], but the present problem shows more complexity. The approach is well suited to numerical implementation, and a numerical example shows the usefulness of the method. Several ways of testing the numerical are available, e.g., the optical theorem (appropriately formulated) and the sum rule. All these verification show excellent agreement.

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