I. Introduction

Fig. 1 shows an array of $S$ slots of infinite length on the ground plane $y = 0$, loaded by dielectric semi-cylinders ($\varepsilon_r, \mu_r$) of radii $R_i$ whose axes are located at $y = 0, x = h_i, i = 1, 2, \ldots, S$. Without loss of generality it is taken that $h_1 = 0$. Region 0 ($y > 0$) has the properties of vacuum ($\varepsilon_0, \mu_0$), and the background medium (region $\sigma$), taken to be lossless, has parameters ($\varepsilon_r, \mu_r$). Dielectric losses in the cylinders may be accounted for via complex constitutive parameters. The primary excitation is an arbitrarily polarized plane wave obliquely incident from region 0. This structure may be used as a polarization selector, a radiation suppressor over a given frequency bandwidth, and a multi-slot/multi-dielectric open waveguide.

Previous related studies [1]–[8] only concern certain special (or limiting) cases, with the two-dimensional cases wherein a single slot is illuminated, with a $z$-invariant primary excitation in presence of either dielectric or perfectly conducting cylinders. In such cases TE' and TM' waves decouple and can be treated separately [9]. The case $S = 1$ of the present quasi three-dimensional problem has been treated in [10] via a combined Green’s function and singular integral equation approach. Nevertheless, when $S(S > 1)$ semi-cylinders coexist in region $y < 0$, as in Fig. 1, the method of [10] breaks down due to insuperable difficulties in obtaining the Green’s function of the structure.

To overcome such difficulties, the present communication introduces a hybrid formulation technique, flexible and easily implemented, to the problem. Section II. The main idea is to use series expansions in cylindrical wave functions, in addition to properly selected singular integral

Abstract—We study oblique diffraction of arbitrarily polarized plane-waves by an array of slots of infinite length on a common ground plane, backed by an array of dielectric semi-cylinders. The formulation is based on a combined eigenvectors expansion and integral equation approach. For the diffracted field, series expansions in cylindrical wave functions are used to which several singular integral terms are superimposed that fully account for the presence of each of the slots. The relevant system of singular integral equations is discretized by an exponentially convergent Nyström method. Noticeably, all matrix elements take simple closed-form expressions. Numerical examples and case studies illustrate the convergence of algorithm and bring to light the influence of the dielectric loads on the characteristics of the structure.

Index Terms—Dielectric cylinders, electromagnetic diffraction, integral equations, Nyström method, slot arrays.

I. INTRODUCTION

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Let \((\vec{E}^{inc}, \vec{H}^{inc})\) be the known field excited in region 0 when all slots are absent (slots short-circuited). Then \((\vec{E}^{0}, \vec{H}^{0}) = (\vec{E}^{1d}, \vec{H}^{1d}) - (\vec{E}^{inc}, \vec{H}^{inc})\) defines the scattered field in region 0. The transmitted (total) field in region \(i\) for \(i = 1, 2, \ldots, S\) or \(i = a\) will be denoted by \((\vec{E}^{i}, \vec{H}^{i})\). All fields and current densities will vary as \(\vec{X}(\vec{r}) = \vec{X}(\vec{r})e^{j\omega t}\), just like the incident field, where
\[
k = k_0 \cos \theta_0.
\]

### B. Field Representations

For \(i = 0, 1, 2, \ldots, S\), let
\[
G_i^0(\vec{r}, \vec{r}') = -\frac{1}{4\pi} \left\{ H_0^1(\gamma_i R^+) - \frac{1}{2\pi} \gamma_i \right\}
\]
\[
G_i^0(\vec{r}, \vec{r}') = -\frac{1}{4\pi} \left\{ H_0^1(\gamma_i R^+) + \frac{1}{2\pi} \gamma_i \right\}
\]
where
\[
R^\pm = \left[ (x - x')^2 + (y \pm y')^2 \right]^{1/2},
\]
\[
\gamma_i = (k_0^2 - k_i^2)^{1/2}, \quad -\pi / 2 \leq \arg \gamma_i < \pi / 2
\]
be auxiliary quantities to be used shortly.

1) **Region 0:** Let \(\vec{M}_r(x) = \vec{E}^{1d}(x, 0) \times \hat{y} = M_r^1(x) \hat{x} + M_r^2(x) \hat{z}\) denote the equivalent surface magnetic current density across the \(q-t h\) slot. Then [1], [10], [12]
\[
E_r^0(\vec{r}) = -\frac{j\omega \varepsilon_0}{\gamma_0} \sum_{q=1}^{S} \int M_r^1(x') \frac{\partial G_0^0(x', 0^+; \vec{r})}{\partial y'} dx'
\]
\[
H_r^0(\vec{r}) = \sum_{q=1}^{S} \int M_r^0(x') G_0^0(x', 0^+; \vec{r}) dx'
\]
\[
- \frac{k_0}{\gamma_0} \sum_{q=1}^{S} \int M_r^0(x') \frac{\partial G_0^0(x', 0^+; \vec{r})}{\partial x'} dx'
\]
where \(C_q = \{ x : h_1 - w_i \leq x' \leq h_1 + w_i \}\) denotes the \(x\)-axis interval occupied by the \(q-t h\) slot.

2) **Region \(i\), \(i = 1, 2, \ldots, S\):** For the field in region \(i\) we make use of the representations
\[
E_r^i(\vec{r}) = \sum_{n=1}^{\infty} \frac{a_n J_n(\gamma_i \rho_i) \sin(n \varphi_i)}{\gamma_i} + \frac{j \omega \varepsilon_0}{\gamma_i^2} \int M_r^1(x') \frac{\partial G_i^0(x', 0^+; \vec{r})}{\partial y'} dx'
\]
\[
H_r^i(\vec{r}) = \sum_{n=0}^{\infty} \frac{b_n J_n(\gamma_i \rho_i) \cos(n \varphi_i)}{\gamma_i} - \frac{\varepsilon_0}{\gamma_i^2} \int M_r^1(x') G_i^0(x', 0^+; \vec{r}) dx'
\]
\[
+ \frac{j \omega \varepsilon_0}{\gamma_i^2} \int M_r^1(x') \frac{\partial G_i^0(x', 0^+; \vec{r})}{\partial x'} dx'
\]
where \((\rho_i, \varphi_i)\) are the polar coordinates of the observation point \(\vec{r}\) in the local coordinate system associated with the \(i-t h\) semicylinder.

The integral terms in (9) and (10) represent the \(z\)-components of the field that would have been excited in the semi-infinite region \(y < 0\) if the surface magnetic current density \(-\vec{M}^i\) had been impressed across \(C_i\), assuming that a) all slots are absent (short-circuited) and b) the entire region \(y < 0\) is filled with the medium \((\varepsilon_i, \mu_i)\).
Fig. 2. Translation of the reference system of the $p$-th cylinder.

3) Region $a$: Finally, for the field in region $a$ we make use of the expansions

$$E_i^a(r) = \sum_{n=1}^{\infty} \sum_{m=-\infty}^{\infty} A_n^m H_0^2(\gamma a r) \sin(n \varphi_a)$$

$$H_i^a(r) = \sum_{n=1}^{\infty} \sum_{m=0}^{\infty} B_n^m H_0^2(\gamma a r) \cos(n \varphi_a).$$

Equations (11) and (12) will be supplemented with the following addition-translation formula (parameters are specified in Fig. 2):

$$H_i^{(2)}(\gamma a r) e^{in \varphi_a} = \sum_{n=-\infty}^{\infty} H_{m+n}^{(2)}(\gamma a r) J_n(\gamma a R_j) e^{in \varphi_a}$$

where $\varphi_a = 0$, if $p > q$, otherwise $\varphi_a = \pi$, which enables one to express the elementary cylindrical wave referred to the local coordinate system of the $p$-th cylinder as a series of cylindrical wave functions referred to the local coordinate system of the $q$-th cylinder.

Remark 1: In terms of $E_i(\tilde{r})$ and $H_i(\tilde{r})$, the transverse (to $z$-axis) components $E_{i\perp}(\tilde{r})$ and $H_{i\perp}(\tilde{r})$ can be obtained everywhere from

$$j(\omega^2 \mu - k_0^2)E_{i\perp} = -k_i \nabla E_i - \omega \mu \times \nabla H_i$$

$$j(\omega^2 \mu - k_0^2)H_{i\perp} = -k_i \nabla H_i + \omega \mu \times \nabla E_i.$$ (14)

Remark 2: It is important here to stress out that, in the context of the proposed field representations, both $E_i(x, 0)$ and $E_{i\perp}(x, 0)$ vanish on the metallic parts of the ground plane and 2) are continuous on the slots, i.e., the boundary conditions for the tangential electric field are satisfied on the entire $y = 0$ plane.

III. INTEGRAL EQUATIONS

A. The First Set

By satisfying for $i = 1, 2, \ldots, S$ the continuity conditions $H_i^{(2)}(x, 0^-) = H_i^{(2)}(x, 0^+)$, we obtain after some manipulations omitted here for brevity we obtain the following (2S) set of coupled singular IEs:

$$\frac{1}{\omega \mu_0} \int_{q_1}^{q_2} \left[ \gamma_0^2 \mathcal{M}_i^0(x) + j k_j \frac{d \mathcal{M}_j^0(x)}{dx} \right] + \frac{\gamma_0^2}{\omega \mu_1} \mathcal{M}_i^0(x)$$

$$\frac{1}{\omega \mu_0} \int_{q_1}^{q_2} \left[ j k_j \frac{d \mathcal{M}_j^0(x)}{dx} + \left( k_0^2 + \frac{d^2}{dx^2} \right) \mathcal{M}_j^0(x) \right]$$

$$+ \frac{1}{\omega \mu_1} \left[ j k_j \frac{d \mathcal{M}_j^0(x)}{dx} + \left( k_0^2 + \frac{d^2}{dx^2} \right) \mathcal{M}_j^0(x) \right]$$

$$+ \frac{j k_j}{\gamma_i} \sum_{n=1}^{\infty} a_n^0 J_n(\gamma_i x) + \frac{j k_j}{\gamma_i} \sum_{m=-\infty}^{\infty} b_m J_m(\gamma_i x)$$

$$= H_i^{(2)}(x, 0^+), \quad x \in C_i; \quad i = 1, 2, \ldots, S.$$ (16)

Fig. 3. Addition theorem for the Hankel function.

where

$$x_i = x - h_i$$

$$\mathcal{M}_i^0(x) = \int_0^{\infty} M_i(x') H_0^{(2)}(\gamma_i x - x') dx'$$

$$\mathcal{M}_i^0(x) = \int_0^{\infty} M_i(x') H_0^{(2)}(\gamma_i x - x') dx'.$$

In (17), the relations $(\partial/\partial y') H_i^{(2)}(\gamma_i R_j) = \pm (\partial/\partial y) H_i^{(2)}(\gamma_i R_j)$ and $(\partial^2/\partial y'^2 - k_0^2 + k_0^2) H_i^{(2)}(\gamma_i R_j) = 0$ have been used, with $R_j$ given in (6).

B. The Second Set

We substitute (9) and (11) into the continuity condition $E_i(R_j, \varphi_j) = E_i^0(R_j, \varphi_j)$, $-\pi \leq \varphi_j \leq 0$, and make use of (4), (13), and the addition theorem for the Hankel functions [13]

$$H_i^{(2)}(\gamma_i R_j) = \sum_{n=-\infty}^{\infty} H_{m+n}^{(2)}(\gamma_i R_j) J_n(\gamma_i x_i)$$

(21)

with the parameters involved in (21) are specified in Fig. 3). Next, we multiply both sides of the resulting equation by $\sin(n \varphi_j)$ for $m = 1, 2, \ldots, \infty$, and integrate from $\varphi_j = -\pi$ to $\varphi_j = \pi$. After some lengthy but otherwise straightforward manipulations, omitted here for brevity, we end up with the following set of IEs:

$$a_m J_m(\gamma_i R_j) - j \gamma_i H_m^{(2)}(\gamma_i R_j) J_m^\prime(m, i)$$

$$- A_m^\prime H_m^{(2)}(\gamma_i R_j) - J_m(\gamma_i R_j) \sum_{n=1}^{\infty} A_n^m F_{m+n}(q, i) = 0$$

$$m = 1, 2, \ldots, \infty; \ i = 1, 2, \ldots, S.$$ (22)

where

$$x_i' = x' - h_i$$

$$J_m^\prime(m, i) = \frac{1}{2} \int_0^{\infty} M_j(x') \left[ J_{m-1}(\gamma_i x') \pm J_{m+1}(\gamma_i x') \right] dx'$$

$$F_{m+n}(q, i) = H_m^{(2)}(\gamma_i D_q) e^{i(m+n) \psi_q}$$

$$\pm (-1)^m H_{m+n}^{(2)}(\gamma_i D_q) e^{-j(m+n) \psi_q}$$

(25)

The prime in the summation over $q$ in (22) is used to remind that the term with $q = i$ is excluded from the sum.

C. The Third Set

In a similar way, from $H_i^{(2)}(R_j, \varphi_j) = H_i^{(2)}(R_j, \varphi_j)$, $\pi \leq \varphi_j \leq 0$, for

$$i = 1, 2, \ldots, S$$

and for $m = 0, 1, \ldots, \infty$, one obtains the IEs:

$$b_m J_m(\gamma_i R_j) + \frac{\gamma_i e_m}{2 \omega \mu_1} H_m^{(2)}(\gamma_i R_j) \left[ j k_j J_i^\prime(m, i) - j k_j J_i^\prime(m, i) \right]$$

$$- B_m^\prime H_m^{(2)}(\gamma_i R_j) - \frac{\varepsilon_m}{2} J_m(\gamma_i R_j) \sum_{n=1}^{\infty} B_n^m F_{m+n}(q, i) = 0$$

(27)

where

$$\varepsilon_m = 2 - \delta_m$$

and

$$J_i^\prime(m, i) = \int_0^{\infty} M_i(x') J_i(m, x') dx'.$$ (28)
D. The Fourth Set

The remaining boundary conditions
\[ \tilde{\varphi}_i = \hat{\tilde{\varphi}}(R_i, \varphi_i) + \hat{\tilde{\varphi}}(R_i, \varphi_i), \quad -\pi \leq \varphi_i \leq 0 \]
treated along the same lines end up with the following set of IEs:
\[
\begin{align*}
\tilde{\mathbf{U}}^M_{m_i} &= \left( \begin{array}{c} A_m \hat{b}_m \frac{1}{2} \\
B_m \hat{b}_m \end{array} \right) + \varepsilon_m \mathbf{T}^M_{m_i} \sum_{q=1}^{S} \left( \sum_{n=1}^{\infty} A_n^q \hat{F}_m^q(q, i) \right) \\
- \mathbf{T}^M_{m_i} \left( \begin{array}{c} \tilde{\varphi}_i \\
\tilde{\varphi}_i \end{array} \right) &= \left( \begin{array}{c} 0 \\
0 \end{array} \right),
\end{align*}
\]
m = 0, 1, \ldots, \infty; \ i = 1, 2, \ldots, S.

Here,
\[
\begin{align*}
\mathbf{T}^M_{m_i} &= \left( \begin{array}{c} \frac{\sum_{k=1}^{L} \omega^k}{\gamma_i} H^{(2)}(\gamma_i R_i) \\
-\frac{\sum_{k=1}^{L} \omega^k}{\gamma_i} H^{(1)}(\gamma_i R_i) \end{array} \right), \\
\mathbf{T}^J_{m_i} &= \left( \begin{array}{c} \frac{\sum_{k=1}^{L} \omega^k}{\gamma_i} H^{(2)}(\gamma_i R_i) \\
-\frac{\sum_{k=1}^{L} \omega^k}{\gamma_i} H^{(1)}(\gamma_i R_i) \end{array} \right),
\end{align*}
\]
whereas \( \mathbf{T}^J_{m_i} \) and \( \mathbf{T}^M_{m_i} \) result from \( B_m \) and \( A_m \), respectively, after replacing \( H^{(2)}(\cdot) \) and \( H^{(1)}(\cdot) \) in the right side of (30) and (31) by \( J_m(\cdot) \) and by \( J_m(\cdot) \). The primes in \( B_m^{(2)}(\cdot) \) and \( J_m(\cdot) \) denote derivatives with respect to the argument.

IV. SOLUTION BY THE NYSTROM METHOD

In view of the edge conditions, \( M_i'(x') \) and \( M_i'(x') \) will be sought in the form
\[
M_i'(x') = \frac{F_i'(\tau')}{\sqrt{1 - (\tau')^2}}, \quad M_i'(x') = \sqrt{1 - (\tau')^2} F_i'(\tau')
\]
where
\[
\tau' = \frac{x'-h_i}{w_i}.
\]

A. Discretization of (16) and (17)

Substitute from (32) into (19) and (20). Then, working in the framework of the Nyström method of [14], any singular integrals encountered can be analytically treated as in [11] (see also, [15]), and any regular integrals can be computed via the Gauss-Ghebey rules [13]
\[
\begin{align*}
\int_{-1}^{1} \sqrt{1 - \tau^2} f(\tau) \tau d\tau &= \frac{\pi}{L+1} \sum_{n=1}^{L} \left( 1 - \chi_n^2 \right) f(\chi_n) \\
\int_{-1}^{1} f(\tau) \tau d\tau &= \frac{\pi}{L} \sum_{n=1}^{L} f(t_n)
\end{align*}
\]
where
\[
t_n = \cos \left( \frac{2n-1}{2L} \pi \right), \quad \chi_n = \cos \left( \frac{n\pi}{L+1} \right)
\]
\((L \text{ has to be selected as high as needed to assure any prescribed accuracy, as specified in Section V.})\). After carrying out the relevant integrations we satisfy the resulting two equations, respectively, at the collocation points \( x = x_m' \) and \( x = x_m' \), where
\[
x_m' = h_i + w_i t_m, \quad \chi_m = h_i + w_i \chi_m.
\]
This yields the following discrete counterparts of (16) and (17) for \( m = 1, 2, \ldots, L \) and for \( i = 1, 2, \ldots, S \):
\[
\begin{align*}
\sum_{n=0}^{\infty} b_n^j J_n(\gamma_i w_i \chi_n) &= \sum_{q=1}^{S} \left( \sum_{n=1}^{\infty} F_n^j(t_n) K_{m_i}^{(i)} \right) \\
+ \sum_{q=1}^{S} \frac{L}{\gamma_i} F_n^j(t_n) L_{m_i}^{(i)} &= \frac{\varepsilon_i}{\pi} (x_m') \quad \text{(38)}
\end{align*}
\]
The matrix elements \( K_{m}^{(i)}, L_{m}^{(i)}, F_{m}^{(i)}, \) and \( Q_{m}^{(i)} \) have simple closed form expressions, omitted here for brevity.

B. Discretization of (22), (27) and (29)

The discrete counterparts of these IEs can be obtained by simply substituting in each of \( J^\pm_i \) and \( J_i \) by
\[
\begin{align*}
J^\pm_i(m, i) &= \frac{w_i}{L+1} \sum_{n=1}^{L} \left( 1 - \chi_n^2 \right) J_{m+1}(\gamma_i w_i \chi_n) \\
& \pm J_{m+1}(\gamma_i w_i \chi_n) F_i(\chi_n) \\
J_i(m, i) &= \frac{w_i}{L} \sum_{n=1}^{L} J_{m+1}(\gamma_i w_i \chi_n) F_i(\chi_n)
\end{align*}
\]
respectively. Equation (40) and (41) result from (24) and (28) by applying the rules (34) and (35), respectively.

C. Far-Scattered Field

In polar coordinates \((\rho, \varphi)\) the \( z\)-components of the far scattered field in region 0 and of the total far-field in region \( a \) have the following asymptotic expressions.

1) Region 0:
\[
E_i(\rho, \varphi) = \sqrt{\frac{2L}{\pi \rho_0}} e^{-j \rho \varphi} \sin \varphi \mathcal{E}_i(\gamma_0 \cos \varphi)
\]
\[
H_i(\rho, \varphi) = \frac{\gamma_0}{2 \rho^2} \sqrt{\frac{2L}{\rho_0}} e^{-j \rho \varphi} \mathcal{H}_i(\gamma_0 \cos \varphi)
\]
where, for \( v = \gamma_0 \cos \varphi \), see (44) and (45) at the bottom of the page.
2) Region a:

\[
E_1^e(\rho, \varphi) = \sqrt{\frac{2j}{\pi \epsilon_0 \rho}} e^{-j \gamma \rho} \sum_{n=1}^{\infty} \sin(n \varphi) \left[ A_n^1 + \frac{\epsilon_{n+1} - \epsilon_n}{\epsilon_n} \sum_{m=1}^{\infty} A_m^1 G_{m,n}(m,n) \right] \]

\[
H_1^e(\rho, \varphi) = \sqrt{\frac{2j}{\pi \epsilon_0 \rho}} e^{-j \gamma \rho} \sum_{n=0}^{\infty} \cos(n \varphi) \left[ B_1^1 + \frac{\epsilon_{n+1} - \epsilon_n}{\epsilon_n} \sum_{m=0}^{\infty} B_m^1 G_{m,n}(m,n) \right]
\]

where

\[
G_{m,n}(m,n) = J_{m-n}(\gamma a h_0) \pm (-1)^n J_{m+n}(\gamma a h_0)
\]

The transverse (to z-axis) field components in regions 0 and a can be obtained via (14) and (15).

V. NUMERICAL RESULTS AND DISCUSSION

In the following it is assumed, unless otherwise specified, that \( E_0 = 1V/m, (\varepsilon_x, \mu_a) = (\varepsilon_0, \mu_0) \), and that the semi-cylinders have identical geometrical and physical parameters: \( \varepsilon_i = \varepsilon_x, \mu_i = \mu_0, \omega_i = \omega, R_i = R, i = 1, 2, \ldots, S \). Only the symmetric case, \( h_j = (i-1)h (i = 1, 2, \ldots, S) \), wherein the distance between any two consecutive slots is the same, equal to \( h \), will be considered. In obtaining numerical results, the size of the infinite linear algebraic system derived in the preceding section is truncated to the finite value 2LS + 4n_{max}S by retaining n_{max} terms in each of the infinite series in (9)-(12) (and, similarly, in (38), (39)).

The correctness of the results has been tested by following the criterion of energy balance. Moreover, in the special case \( S = 1 \) our results were found to coincide with those of [10].

For increasing \( L \) (the number of points of the Gauss-Chebyshev rules used in evaluating the matrix elements) and for n_{max} = L, Fig. 4 shows the logarithm (base 10) of the relative errors \( |M_1^e(0) - M_1^{e_{	ext{asym}}}(0)| / |M_1^e(0)| \) at the center \( x = 0 \) of the first slot when \( S = 2 \). Here \( M_1^{e_{	ext{asym}}}(0) \) and \( M_1^{e_{	ext{asym}}}(0) \) are the values to which \( M_1^e(0) \) and \( M_1^{e_{	ext{asym}}}(0) \) settle down for sufficiently large \( L \) (and n_{max}). Apparently, the convergence is exponential. (Note: Our computations in this example show that \( |M_1^e(0)| \) and \( |M_1^{e_{	ext{asym}}}(0)| \) settle down to their final values 0.4712579942343935 and 0.7371905396420925, respectively, for \( L = n_{max} \geq 22 \); these asymptotic values can be treated as exact values.)

The fast convergence of the algorithm enables one to very accurately evaluate both the near-field and the far-field. As an example, for \( S = 3 \), Fig. 5 shows a typical snapshot of the total magnetic field at \( \omega t = \pi/2 \) both inside and around the dielectric semi-cylinders when \( \lambda_0 = 3mm, 2w = 1.5\lambda_0, R = \lambda_0 / 2, h = \lambda_0 / 2, \varepsilon_0 = \gamma = \pi / 4, f = 10GHz \).

For several values of \( S \), Fig. 7 shows how \( \sigma_{abs} \) varies with the distance \( h \) between the two slots. The dashed straight line about which \( \sigma_{abs} \) oscillates pertains to the case \( S = 1 \) of a single slot. As expected, the amplitude of the observed standing-wave-like pattern is large for small \( h \), due to strong interactions between the cylinders, but it gradually diminishes as \( h \) increases.

For several values of \( S \), Fig. 8 shows how the absorption efficiency is affected by possible uncertainty of the complex permittivity value. As seen, \( \sigma_{abs} \) can be effectively controlled by properly selecting \( S \) and \( \varepsilon_r \).
A computationally efficient Nyström method for analyzing oblique scattering of arbitrarily polarized waves by an array of slots loaded by dielectric semi-cylinders has been presented. Filling the system matrix requires no numerical integration. The algorithm converges exponentially and, thus, extremely accurate results may be obtained both for the near and far fields.

VI. CONCLUSION

A new surface integral equation formulation is presented for characterizing electromagnetic radiation by conformal microstrip arrays on finite curved bodies of arbitrary shapes. The surface equivalence principle is used to reduce the original problem to two equivalent problems, one for the external medium and another for the internal medium. Electric field integral equations are applied to the conducting surfaces, and weighted sums of the field integral equations corresponding to the external and internal dielectric regions with appropriate weighted coefficients are applied to the dielectric interface. The integral equations are solved via the method of moments (MoM) procedure, to which the memory requirement and computational complexity pertinent is reduced by employing the adaptive integral method (AIM). Numerical results are presented to demonstrate the validity and accuracy of the method.

Index Terms—Antenna arrays, antenna radiation patterns, conformal antenna, moment method, patch arrays.

REFERENCES


Analysis of Radiation Characteristics of Conformal Microstrip Arrays Using Adaptive Integral Method

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Abstract—A new surface integral equation formulation is presented for characterizing electromagnetic radiation by conformal microstrip arrays, which is based on the Nyström method for solving singular integral equations. The method is applied to analyze the radiation characteristics of conformal microstrip arrays on finite curved bodies of arbitrary shapes. The surface equivalence principle is used to reduce the original problem to two equivalent problems, one for the external medium and another for the internal medium. Electric field integral equations are applied to the conducting surfaces, and weighted sums of the field integral equations corresponding to the external and internal dielectric regions are applied to the dielectric interface. The integral equations are solved via the method of moments (MoM) procedure, to which the memory requirement and computational complexity is reduced by employing the adaptive integral method (AIM). Numerical results are presented to demonstrate the validity and accuracy of the method.

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