Electromagnetic Wave Scattering from Some Vegetation Samples

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Abstract—For an incident plane wave, the field inside a thin scatterer (disk and needle) is estimated by the generalized Rayleigh-Gans (GRG) approximation. This leads to a scattering amplitude tensor equal to that obtained via the Rayleigh approximation (dipole term) with a modifying function. For a finite length cylinder the inner field is estimated by the corresponding field for the same cylinder of infinite length.

The effects of different approaches in estimating the field inside the scatterer on the backscattering cross section are illustrated numerically for a circular disk, a needle, and a finite length cylinder as a function of the wave number and the incidence angle. Finally, the modeling predictions are compared with measurements.

I. INTRODUCTION

In recent years, scattering from randomly oriented dielectric nonspherical scatterers with shape and dielectric constant similar to the vegetation components has been used to model electromagnetic wave scattering from vegetation canopies [1]-[10]. Circular disks have been used to model leaves in deciduous vegetation [1]-[7], needles to model leaves in coniferous vegetation [7], [8], and finite cylinders to model branches in defoliated vegetation [9]. In such a modeling approach, the far zone scattered field is first calculated in the principal frame of the scatterer. Then by a coordinate transformation [3], [4] the scattered field can be calculated for any orientation that is needed to study electromagnetic wave interaction with vegetation canopies.

The motivation of this study is to justify the previously proposed modeling approaches on the level of vegetation components. Thus, the far zone scattered fields in the principal frame of some dielectric scatterers with cylindrical symmetry suitable for modeling the leaves and branches of trees are summarized. The backscattering cross sections are calculated and compared with measurements.

A general formulation for the scattered field $\mathbf{E}_s$ can be obtained via the Helmholtz integral equation. For a scatterer located at the origin, the Helmholtz integral equation relates the far zone scattered field to the field inside the scatterer $\mathbf{E}_m(r')$ through the relation [10]

$$\mathbf{E}_s = \frac{e^{-ikr}}{r} \frac{k^2}{4\pi} (\mathbf{l} - \delta\mathbf{s}) \int \int \int (\epsilon_r - 1) \mathbf{E}_m(r') e^{i\delta\mathbf{s} \cdot r'} d\mathbf{r}'$$

where $k$ is the host medium wave number, $\mathbf{l}$ is the unit dyad, $\delta$ is the observation direction, and $\epsilon_r$ is the scatterer dielectric constant relative to that of the host medium.

In (1) the main task in the scattered field calculation is to estimate the inner field. The Rayleigh estimation for the inner field is valid for very small scatterers [11], the Rayleigh-Gans estimation is valid for tenuous scatterers, and the generalized Rayleigh-Gans approximation (GRG) is valid for a nontenuous scatterer with at least one of its dimensions small compared to the wavelength [10]. In this study the GRG approximation is used to calculate the scattering amplitude tensor for the disk and the needle shaped scatterers in Sections II-A and II-B, respectively.

In Section III the scattering amplitude tensor for a finite length cylinder is derived by estimating the inner field with the field inside a similar cylinder of infinite length.

In Section IV numerical calculations of the backscattering cross section versus the incidence angle and the wave number are carried out when the field inside the scatterer is estimated via the Rayleigh and the GRG approximation for disks and needles. Also, the difference between estimating the field inside a finite cylinder by the GRG approximation and by the field inside an infinite cylinder with the same properties is illustrated.

Finally, comparisons with backscattering measurements from some vegetation samples are presented.

II. SCATTERING FROM A THIN CYLINDRICAL SCATTERER

Consider a plane wave with a dyadic representation

$$\mathbf{E}_i = \sum_{q=1}^{\infty} q_q \hat{q}_q \cdot \mathbf{E}_0 e^{-i\delta\mathbf{s} \cdot \hat{r}}$$

where $\mathbf{E}_0$ is the amplitude of the incident field, $\delta(\theta, \phi)$ is the incident direction, and $\hat{r}_q$ and $\hat{q}_q$ are the polarization
vectors defined as (see Fig. 1)

\[
\begin{align*}
\hat{\mathbf{r}} &= \sin \theta_y (\hat{x} \cos \phi_y + \hat{y} \sin \phi_y) - \hat{z} \cos \theta_y \\
\hat{\mathbf{h}} &= \frac{\hat{z} \mathbf{i}}{|\mathbf{i}|} = \hat{y} \cos \phi_y - \hat{x} \sin \phi_y \\
\hat{\mathbf{v}} &= \hat{\mathbf{h}} \times \hat{\mathbf{r}} = -\cos \theta_y (\hat{x} \cos \phi_y + \hat{y} \sin \phi_y) - \hat{z} \sin \theta_y.
\end{align*}
\]

(3)

For a scatterer with an axis of symmetry aligned with the \(\hat{z}\)-axis and one or two of its dimensions small compared with the wavelength of the incident field so that \(kD(\epsilon_y)^{1/2} \ll 1\), where \(D\) is the smallest dimension, we can estimate the field inside the scatterer as [10]

\[
\mathbf{E}_s(r') = \mathbf{a} \cdot \mathbf{E}_o = \mathbf{a} \cdot \sum_{q=-r,h} \hat{q} \cdot \hat{q} \cdot \mathbf{E}_o e^{-jkr'.z}.
\]

(4)

where \(\mathbf{a}\) is the polarizability tensor [5] defined as

\[
\mathbf{a} = a_r \mathbf{I} + (a_n - a_r) \hat{z} \hat{z}
\]

(5)

with

\[
\begin{align*}
a_r &= \frac{1}{(\epsilon_r - 1) g_r + 1} \\
a_n &= \frac{1}{(\epsilon_n - 1) g_n + 1}
\end{align*}
\]

(6)

and \(g_r\) and \(g_n\) are the demagnetizing factors [12] that will be given for a disk and a needle in Sections II-A and II-B, respectively. The Rayleigh estimation for the inner field can be obtained from (4) by letting the exponential term reduce to unity [10].

By substituting (4) into (1) we can write the scattered field as

\[
\mathbf{E}_o = \frac{k^2}{4\pi} (\epsilon_r - 1) (\mathbf{I} - \mathbf{s} \mathbf{s}) \cdot \mathbf{a}
\]

\[
\cdot \sum_{q=-r,h} \hat{q} \cdot \hat{q} \cdot \mathbf{E}_o e^{-jk\mathbf{r}' \cdot (\mathbf{i} - \mathbf{s})} \int \int e^{-jk\mathbf{r}' \cdot (\mathbf{i} - \mathbf{s})} \int \int e^{-jk\mathbf{r}' \cdot (\mathbf{i} - \mathbf{s})}
\]

\[
= \mathbf{F}(\mathbf{s}, \mathbf{i}) \cdot \mathbf{E}_o e^{-jk\mathbf{r}' \cdot \mathbf{r}}
\]

(7)

where \(\mathbf{F}(\mathbf{s}, \mathbf{i})\) is the scattering amplitude tensor.

By using incident unit vector from (3) and the corresponding unit vector for the scattered direction in cylindrical coordinates \((\rho, \phi, z)\), the scattering amplitude tensor in (7) becomes

\[
\mathbf{F}(\mathbf{s}, \mathbf{i}) = \frac{k^2}{4\pi} (\epsilon_r - 1) (\mathbf{I} - \mathbf{s} \mathbf{s}) \cdot \mathbf{a} \cdot \sum_{q=-r,h} \hat{q} \cdot \hat{q} \cdot \int \int \int \int \exp \left\{ -jk \left\{ \rho' \sin \theta_y \cos (\phi_y - \phi') - \rho' \sin \theta_y \cos (\phi_y - \phi') - \hat{z}' \cos \theta_y + \cos \theta_y \right\} \right\} \rho' \rho' \rho' \rho' d\phi' d\rho' dz' d\theta'.
\]

(8)

From [13] we recall the expansions

\[
\exp \left\{ -jk \rho' \sin \theta_y \cos (\phi_y - \phi') \right\}
\]

\[
= \sum_{n=-\infty}^{\infty} (-j)^n J_n(k \rho' \sin \theta_y) \exp \left\{ jn(\phi_y - \phi') \right\}
\]

\[
\exp \left\{ jk \rho' \sin \theta_y \cos (\phi_y - \phi') \right\}
\]

\[
= \sum_{n=-\infty}^{\infty} j^n J_n(k \rho' \sin \theta_y) \exp \left\{ jn(\phi_y - \phi') \right\}
\]

(9)

where \(J_n(\cdot)\) and \(J_n(\cdot)\) are the cylindrical Bessel functions of the first kind.

Then substituting (9) into (8) and integrating over \(d\phi'\) we obtain

\[
\mathbf{F}(\mathbf{s}, \mathbf{i}) = \frac{k^2}{4\pi} (\epsilon_r - 1) v_0 (\mathbf{I} - \mathbf{s} \mathbf{s}) \cdot \mathbf{a} \cdot \sum_{q=-r,h} \hat{q} \cdot \hat{q} \cdot \mu(\mathbf{s}, \mathbf{i})
\]

(10)

where \(v_0\) is the scatterer volume and \(\mu(\mathbf{s}, \mathbf{i})\) is the modifying function to the Rayleigh scattering amplitude given by

\[
\mu(\mathbf{s}, \mathbf{i}) = \frac{2\pi}{v_0} \sum_{n=-\infty}^{\infty} \exp \left\{ jn(\phi_y - \phi') \right\}
\]

\[
\cdot \int \int \int \int \int \exp \left\{ jk \mathbf{s}' \cdot (\cos \theta_y + \cos \theta_y) \right\} \rho' \rho' \rho' \rho' d\phi' d\rho' dz' d\theta'.
\]

(11)

Note that under the Rayleigh estimation for the inner field, the modifying function reduces to unity.

To express (10) in terms of the scattered field polarization vectors we write

\[
\mathbf{I} - \mathbf{s} \mathbf{s} = \sum_{\rho=-r,h} \hat{\rho} \cdot \hat{\rho}
\]

(12)
where \( \hat{\ell}_i \) and \( \hat{h}_i \) can be obtained from (3) by replacing \( \theta_i \) and \( \phi_i \) with \( \pi - \theta_i \) and \( \phi_i \), respectively. Then by substituting (5) and (12) into (10) the scattering amplitude tensor reduces to

\[
F_{\rho\sigma}(\hat{\ell}, \hat{h}) = \frac{k^2}{4\pi} \left[ 4 \pi \left( \epsilon_r - 1 \right) \left( a_r \beta_r \cdot \hat{\ell}_i \right) + \left( a_N \beta_N \cdot \hat{h}_i \right) (n \cdot a_T) \right] \mu(\hat{\ell}, \hat{h}).
\]  

(13)

By calculating the vector dot product \( (\hat{\ell}_i \cdot \hat{h}_i) \) in (13) using (3) and its correspondence for the scattered field we can write the explicit form of the scattering amplitude tensor as

\[
F_{\rho\sigma}(\hat{\ell}, \hat{h}) = \frac{k^2}{4\pi} \left( \epsilon_r - 1 \right) \left[ a_N \sin \theta_i \sin \theta_i - a_T \cos \theta_i \cos \phi_i \right] \mu(\hat{\ell}, \hat{h}).
\]

\[
F_{\rho\sigma}(\hat{\ell}, \hat{h}) = \frac{k^2}{4\pi} \cos \theta_i \sin \phi_i - a_T \mu(\hat{\ell}, \hat{h}),
\]

\[
F_{\rho\sigma}(\hat{\ell}, \hat{h}) = \frac{k^2}{4\pi} \cos \phi_i - a_T \mu(\hat{\ell}, \hat{h}).
\]

Equation (14) constitutes the main results of this section. It has two quantities that depend on the scatterer geometry, the modifying function and the demagnetizing factors. In the following sections we will discuss the demagnetizing factors and the modifying function for a circular disk and a needle.

A. The Demagnetizing Factors and the Modifying Function for a Circular Disk

For a circular disk with a thickness \( 2h \) and radius \( a \), we can write the demagnetizing factors as [12]

\[
g_T = \frac{1}{2 \left( m^2 - 1 \right)} \left[ \frac{m^2}{\sqrt{m^2 - 1}} \sin^{-1} \left( \frac{\sqrt{m^2 - 1}}{m} \right) - 1 \right]
\]

\[
g_N = \frac{m^2}{m^2 - 1} \left[ 1 - \frac{1}{\sqrt{m^2 - 1}} \sin^{-1} \left( \frac{\sqrt{m^2 - 1}}{m} \right) \right]
\]

\[
m = \frac{a}{h}.
\]  

(15)

For the GRG approximation to be valid we should have [10] \( 2k(\epsilon_r)^{1/2} \ll 1 \) and \( a \ll h \). Under these conditions the modifying function in (11) can be approximated as

\[
\mu(\hat{\ell}, \hat{h}) = \frac{4\pi h}{v_0} \sum_{n=-\infty}^{\infty} \exp \left[-j k h \cos \theta_i (\cos \phi_i + \cos \phi_i) \right] J_n(k \rho' \sin \theta_i) J_n(k \rho' \sin \theta_i) \rho' d\rho'.
\]  

(16)

Then by using the addition theorem for the cylindrical functions [14], the modifying function in (16) reduces to

\[
\mu(\hat{\ell}, \hat{h}) = \frac{4\pi h}{v_0} \int_0^a J_n(\rho') \rho' d\rho'
\]

\[
\int_0^a J_n(\rho') \rho' d\rho' = \frac{2J_n^2(Q_{\rho}, a)}{Q_{\rho} a}
\]  

(17)

(18)

where

\[
Q_{\rho} = k \sqrt{\sin^2 \theta_i + \sin^2 \theta_i - 2 \sin \theta_i \sin \theta_i \cos (\phi_i - \phi_i)}.
\]  

(19)

It is clear from (18) and (19) that in the forward direction \( (\theta_i = \theta_i, \phi_i = \phi_i) \), the value of the modifying function reduces to unity. In the backward direction \( (\theta_i = \theta_i, \phi_i = \phi_i + \pi) \), it has its maximum value at normal incidence \( (\theta_i = 0) \).

B. The Demagnetizing Factors and the Modifying Function for a Needle

For a needle with radius \( "a" \) and length \( 2h \) we can write the demagnetizing factors as [12]

\[
g_T = \frac{b(b^2 - 1)}{2} \left[ \frac{1}{b^2 - 1} + \frac{1}{2} \log \left( \frac{b - 1}{b + 1} \right) \right]
\]

\[
g_N = -(b^2 - 1) \left[ \frac{1}{2} b \log \left( \frac{b - 1}{b + 1} \right) + 1 \right]
\]

\[
b = \sqrt{1 - \frac{a^2}{h^2}}.
\]  

(20)

For the GRG approximation to be applicable we should have [10] \( 2k(\epsilon_r)^{1/2} \ll 1 \) and \( a \ll h \). Under these conditions the modifying function in (11) can be approximated as

\[
\mu(\hat{\ell}, \hat{h}) = \frac{2\pi}{v_0} \int_0^a \left[ \rho' \rho' d\rho' \right] \exp \left[-j k h \cos \theta_i (\cos \phi_i + \cos \phi_i) \right] J_n(k \rho' \sin \theta_i) J_n(k \rho' \sin \theta_i) \rho' d\rho'.
\]  

(21)

Similar to the disk case, in the forward direction the modifying function is unity. In the backward direction its maximum value occurs only at normal incidence \( (\theta_i = 90^\circ) \).

III. Scattering from a Finite Length Lossy Dielectric Cylinder

Consider a finite length lossy dielectric circular cylinder with radius \( '"a" \), length \( 2h \) and relative dielectric constant \( \epsilon_r \). The cylinder axis is aligned with the z-axis and is illuminated by a plane wave given in (2). By estimating the field inside the cylinder as the field inside a
similar infinite cylinder we can write the cylindrical components of the inner field as [9]

$$\bar{E}_n(r') = \sum_{n=-\infty}^{\infty} E_{\rho n} \rho' + E_{\theta n} \theta' + E_{\phi n} \phi'$$  \hspace{1cm} (22)

where

$$E_{\rho n} = \frac{k}{\lambda_n} \sum_{n=-\infty}^{\infty} \left\{ j \cos \theta \varepsilon_{\rho n} J_n(\lambda_n \rho') \right\} F_n$$

$$E_{\theta n} = \frac{k}{\lambda_n} \sum_{n=-\infty}^{\infty} \left\{ -n \cos \theta \varepsilon_{\theta n} J_n(\lambda_n \rho') \right\} F_n$$

$$E_{\phi n} = \frac{k}{\lambda_n} \sum_{n=-\infty}^{\infty} \left\{ j \varepsilon_{\phi n} J_n(\lambda_n \rho') \right\} F_n$$

and

$$F_n = E_0 \tau^{-n} \exp \left\{ jn(\phi' - \phi_0) + jkz' \cos \theta_0 \right\}$$

$$\varepsilon_{\rho n} = \frac{j \sin \theta}{R_n J_n(u)} \left\{ \frac{H_n^{(2)}(v)}{v J_n(u)} - \frac{J_n'(u)}{u J_n(u)} \right\}$$

$$\eta_{\theta n} = -\frac{\sin \theta}{R_n J_n(u)} \left\{ \frac{H_n^{(2)}(v)}{v J_n(u)} - \frac{J_n'(u)}{u J_n(u)} \right\}$$

$$\varepsilon_{\phi n} = \frac{j \sin \theta}{R_n J_n(u)} \left\{ \frac{H_n^{(2)}(v)}{v J_n(u)} - \frac{J_n'(u)}{u J_n(u)} \right\}$$

$$\eta_{\phi n} = \frac{j \sin \theta}{R_n J_n(u)} \left\{ \frac{H_n^{(2)}(v)}{v J_n(u)} - \frac{J_n'(u)}{u J_n(u)} \right\}$$

$$R_n = \frac{\pi v^2 H_n^{(2)}(v)}{2} \left\{ J_n'(u) - \frac{J_n'(u)}{u J_n(u)} \right\}$$

$$u = \lambda_n \alpha = \frac{ka}{\sqrt{\varepsilon_{\rho n} - \cos^2 \theta_0}}$$

$$v_j = ka \sin \theta_0.$$  \hspace{1cm} (24)

Then by substituting (22) and (23) into (1) and performing the integration over $d\rho'$ we can write the scattering amplitude tensor element as

$$F_{xy}(\hat{s}, \hat{i}) = 2k^2 h \mu(\hat{s}, \hat{i}) (\varepsilon_{\rho n}, -1) \sum_{m=-\infty}^{\infty} e^{j\varepsilon_{\rho n} \cos \theta_0} \left\{ \frac{k}{2\lambda_n} \left[ (\eta_{\theta n} - j \varepsilon_{\phi n} \cos \theta_0) (\hat{p}_x \cdot \hat{x} + \hat{p}_\phi \cdot \hat{y}) z_{n+1} e^{i\phi_0} \right. \right.$$}

$$- (\eta_{\rho n} + j \varepsilon_{\phi n} \cos \theta_0) (\hat{p}_x \cdot \hat{x} - \hat{p}_\phi \cdot \hat{y}) z_{n-1} e^{-i\phi_0} + e_{\rho n} \hat{z} \right\}$$  \hspace{1cm} (25)

with

$$A_n = \frac{k}{2\lambda_n} (z_{n+1} - z_{n-1})$$

$$B_n = \frac{k}{2\lambda_n} (z_{n+1} + z_{n-1}).$$  \hspace{1cm} (28)
It is clear from (27) that the cross-polarized terms in the scattering amplitude tensor vanish in the backward or forward direction. Also from (26) as the cylinder length becomes very large, the scattered field will be propagating along a cone around the cylinder axis with scattered angle equal to \( \theta_i \).

IV. Numerical Results and Discussion

A quantity of interest in remote sensing is the backscattering cross section. It has the value [16]

\[
\Gamma_{pq} = 4\pi \left| F_{pq}(-\hat{i}, \hat{i}) \right|^2
\]  

(29)

for linear polarization, and the value [16]

\[
\sigma = \pi \left| F_{vv}(-\hat{i}, \hat{i}) \mp F_{hh}(-\hat{i}, \hat{i}) \right|^2
\]  

(30)

for circular polarization. In (30) the + and − signs stand for left-hand (LHC) and right-hand (RHC) circular polarization, respectively.

This section provides an illustration of the backscattering cross sections for the disk and needle-shaped objects and cylinders. In addition, the effects due to different estimates of the inner fields on the backscattering cross sections are also shown. Finally, comparisons are made with measurements from real leaves and branches.

A. Backscattering Cross Sections of Disk, Needle, and Cylinders

The geometry of the scattering problem and the associated symbols are defined in Fig. 1. Fig. 2 shows the backscattering cross sections under the Rayleigh and the generalized Rayleigh–Gans approximation (GRG) as a function of the incidence angle. It is seen that the results between the two approximations are quite different except at normal incidence, which corresponds to a low frequency condition. The HH polarization gives higher returns at larger angles of incidence since this is a thin disk. In Fig. 3 the backscattering cross sections are plotted versus the wavenumber radius product \((ka)\) at an incidence angle of 50 degrees. As expected there is agreement between the backscattering cross sections from the two approximations at low frequencies. At higher frequencies the GRG leads to oscillatory behavior due to phase variations and differs substantially from the Rayleigh approximation. In Fig. 4 comparisons are made between calculations and measurements reported by Allen and McCormick [16]. The overall agreement is good except at locations near the null where the experimental results are much higher. This appears to be due mainly to the finite equivalent bandwidth of the measurement system and in part to the approximation used for the inner field. This difference is not expected to be of importance when we deal with scattering from randomly oriented disks where returns from disks are summed incoherently.

Next, we consider scattering from a thin needle. This being a cylindrical structure we have three possible approximations: the Rayleigh, the GRG, and the inner field of the same needle of infinite length (infinite cylinder approximation). Both VV and HH polarizations are computed for the three approximations and shown in Fig. 5. For the needle zero degree means incidence along the needle axis. From Fig. 5 the Rayleigh approximation does not agree with other approximations except when the wave is incident perpendicular to the axis of the needle (90 degree incidence). This is analogous to normal incidence on
Fig. 3. The backscattering cross sections of a circular disk \( \theta = 50^\circ, \, a = 5 \, \text{cm}, \, h = 0.1 \, \text{mm}, \, \epsilon_r = 28.04 - j13.34 \) under Rayleigh and GRG approximations versus the wave number.

![Graph of backscattering cross sections](image)

**Table 1**

<table>
<thead>
<tr>
<th>SAMPLE NO.</th>
<th>(1)</th>
<th>(2)</th>
<th>(3)</th>
</tr>
</thead>
<tbody>
<tr>
<td>( k_a )</td>
<td>0.762</td>
<td>1.521</td>
<td>3.642</td>
</tr>
<tr>
<td>( h/a )</td>
<td>0.1060</td>
<td>0.1</td>
<td>0.1009</td>
</tr>
<tr>
<td>( \text{Re} , \epsilon_r )</td>
<td>3.12</td>
<td>3.11</td>
<td>3.13</td>
</tr>
</tbody>
</table>

Fig. 4. Comparisons between calculated (14), (18), (30) and the measured backscattering cross sections of a circularly polarized plane wave exciting circular disks [16] \( f = 2.86175 \, \text{GHz}, \, \text{Im} (\epsilon_r) = 0.036 \).

the disk. The GRG gives similar results to those using the infinite cylinder approximation. Its results are slightly lower around zero and higher around ninety degree incidence than the corresponding infinite cylinder results. In

![Graph showing comparison between Theory and Measurement](image)

Fig. 6 scattering from the same cylinder as the one in Fig. 5 is plotted versus \( k_a \). It is seen that beyond \( k_a \) equal to about 0.5 the infinite cylinder approximation leads to a decreasing cross section versus frequency along with VV approaching HH. This means that there are more cancellations due to phase variations than what is included in the GRG approximation. In addition, VV and HH remain different at higher frequencies under the GRG approximation, which is not a correct behavior. To verify the scattering calculations using the infinite cylinder approximation we show comparisons for four different cylinders in Fig. 7. Here again, the agreement is good except near null locations where the experimental measurements are higher than the calculated. The possible causes are as stated in the preceding paragraph.

B. Comparison with Leaf and Branch Measurements

In this subsection we shall examine how well the disk and cylinder can be used to model leaves and branches by comparing their cross sections with measurements on real leaves and branches. For the leaf model we shall use a relative dielectric constant derived for a leafy vegetation in terms of the leaf gravimetric moisture content (Mg) [18]. For the branch model we assume a relative dielectric value of 9.6 \(- j4.03\).

The disk and finite cylinder models are compared with measurements from an aspen leaf and a birch stick when they are illuminated by a circularly polarized plane wave. The measurements were acquired using a specially designed polarimetric scattering amplitude tensor measurement facility [19]. It uses circularly polarized (RHC or LHC) transmit signals. The amplitudes of the two orthogonal RHC and LHC components of the backscatter return
are measured. From these measurements, the scattering amplitude tensor for circular polarization can be determined. A receiver with very high sensitivity and wide dynamic range permits accurate measurements at a single frequency. The apparatus is operated out-of-doors at an unobstructed site, with the antenna directed vertically upward toward the target. This virtually eliminates unwanted reflections from structures. The targets are supported using fine nylon threads that exhibit negligible effect.

Fig. 8 shows the comparison between the calculated and measured values for the backscattering cross sections of an aspen leaf illuminated by a circularly polarized incident plane wave. The leaf is located in the x-y plane with
Fig. 7. Comparisons between the calculated (27), (30) and the measured [16] backscattering cross sections of a circularly polarized plane wave exciting dielectric rods \( f = 2.86175 \text{ GHz}, \text{Im}(\varepsilon_r) = -0.036 \).

Fig. 8. Comparisons between the calculated (14), (18), (30) and the measured backscattering cross sections for an aspen leaf illuminated by a circularly polarized plane wave \( a = 2.275 \text{ cm}, h = 0.1 \text{ mm}, f = 9.6 \text{ GHz}, Mg = 0.5 \).

The incident wave is in the \( x-z \) plane \( (\phi_i = 0) \). We let the disk radius be equal to the leaf base. There is a general agreement between the model and the measurements in the angular trend except where the nulls occur. Since a real leaf is not flat, perfect agreement is not expected.

The comparison between the measured cross section of a birch stick and that calculated from a cylinder using ex-
Fig. 9. Comparisons between the calculated \( \sigma = 0.95 \text{ cm}, h = 6.25 \text{ cm}, f = 9.6 \text{ GHz}, \epsilon = 9.6 - j4.03 \) and the measured backscattering cross sections for a birch stick illuminated by a circularly polarized plane wave. There is a very good agreement between the theory and the measurement at large angles of incidence. For small angles of incidence the higher value of the measurement in Fig. 9 is due possibly to diffraction from the stick ends. In actual canopy modeling, this edge diffraction effect does not exist since branches do not have sharp cut-offs.

V. CONCLUSION

The GRG approximation seems to be a reasonable choice for calculating the scattered field from a circular disk and its use to construct scattering models for deciduous vegetation appears acceptable.

A finite length circular cylinder can be used to model scattering from tree branches by using the infinite cylinder approximation.

ACKNOWLEDGMENT

The authors want to thank Dr. L. E. Allen for his assistance in acquiring measurements.

REFERENCES

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