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Kammersgaard, Nikolaj Peter Brunvoll; Kvist, Søren H.; Thaysen, Jesper; Jakobsen, Kaj Bjarne

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Geometrical Theory of Diffraction Formulation for On-Body Propagation

Nikolaj P. B. Kammersgaard, Søren H. Kvist, Jesper Thaysen, and Kaj B. Jakobsen

Abstract—A Geometrical Theory of Diffraction model for on-body propagation is developed in the article. The exact solution to the canonical problem of a plane wave incident on an infinitely long cylinder, with arbitrary constitutive parameters, is found. The same is done for a magnetic and an electric infinitesimal dipole source of any orientation, located on the surface of the cylinder. The exact solutions are transformed with the Watson transformation to yield asymptotic expressions that are valid in the deep shadow region. These asymptotic expressions are validated by comparison to the numerically evaluated exact solution. It is found that the expressions are valid as long as the object is opaque, with a geometry down to the size of \( \frac{\sqrt{\kappa^2 + \tau^2}}{\kappa} > \lambda_0/2 \), and the rays not too torsional \( \tau/\kappa < 2 \), where \( \kappa \) and \( \tau \) are the curvature and torsion of the local geometry, respectively. The asymptotic expressions are found to approximate the exact solution significantly better than the asymptotic expression of an equivalent perfect electric conductor (PEC) geometry. The same is the case for the impedance boundary condition asymptotic approximation for low dielectric constant materials. Finally, the asymptotic expressions are generalized so they can be applied to any convex geometry of the human body or an opaque lossy dielectric of electrically large size.

Index Terms—Geometrical Theory of Diffraction, asymptotic approximation, on-body communication, WBAN, lossy dielectric cylinder.

I. INTRODUCTION

O N-BODY antennas and propagation have received enormous attention in recent years. It has been a major research area, but has recently also resulted in many commercial uses. Within the past two years, for example, more than 10 sets of so-called ‘truly’ wireless headphones have been launched. The ‘truly’ wireless term refers to the fact that these headsets have been launched. The same is the case for the impedance boundary condition asymptotic approximation for low dielectric constant materials. Finally, the asymptotic expressions are generalized so they can be applied to any convex geometry of the human body or an opaque lossy dielectric of electrically large size.

A GTD/UTD formulation for the Impedance Boundary Condition (IBC) was done in [11], [12] and for dielectric coated cylinders in [13], [14]. These approximations would be more appropriate to use for on-body propagation than the PEC assumption. The IBC approximation is based on an assumption only valid for a material with high conductivity, however. This assumption is not fulfilled for human tissue at 2.45 GHz. The dielectric coated cylinders has a PEC core which is not a correct assumption for the human body either. Therefore, it is relevant to find a formulation valid for human tissue. However, the approaches used in [11]–[14] are useful.

The purpose of this work is to provide a GTD formulation for the on-body propagation. The model should be as precise as possible but yet intuitive. This is why GTD is chosen in this work instead of UTD. The intuitiveness of the exponential decaying fields in GTD outweighs the precision of the UTD in the transition region. Still the model should include effects from the electromagnetic properties of the human tissue, the frequency and polarization of the electromagnetic field, as well as the curvature and torsion of the geometry. Out of consideration for the intuitiveness and from the good correlation obtained in [1], [3], [7], a homogenous model is used. This choice is further motivated by the relatively small impact of a layered model seen in [15].

The model addresses two types of propagation. The first is the off-body connection between an antenna on the body and an antenna off the body. By assuming the off-body antenna is placed in the far-field, this can be modeled by calculating the fields on an infinite long, lossy dielectric (human tissue) cylinder caused by an incident plane wave. The second is the on-body connection between two antennas placed on the body. This is modeled by calculating the fields on an infinite long, lossy dielectric (human tissue) cylinder caused by an infinitesimal dipole source on the surface of the cylinder. The cylinder is chosen since it gives rise to geodesics (shortest
paths) with all combinations of curvature and torsion. This makes it possible to approximate any kind of geometry by the assumption that the propagation only depends on the local geometry. Therefore, general GTD expressions for the off-body and on-body propagation can be found from these two canonical problems.

The paper is organized as follows. In Section II the eigenfunction solution to the two canonical problems are found. In Section III the asymptotic or GTD expressions are found. In Section IV, numerical calculations of the eigenfunction solution to the plane wave incidence case is compared to the asymptotic solution to validate the expression. Furthermore, the limitations of the GTD solution is investigated. In Section V the GTD expressions for the two cases are generalized to any geometry. Finally, Section VI contains the conclusion.

The approach followed in the paper is found in [16] which covers the topic in detail and is a good source for introduction and details. An $e^{j\alpha z}$ time dependence for the electromagnetic fields are assumed and suppressed throughout the paper. The constitutive parameters of human tissue are taken from [17], [18].

II. EIGENFUNCTION SOLUTION

A. Plane Wave Incidence

The solution to the canonical problem of an infinite long, lossy dielectric cylinder illuminated by a oblique incident plane wave is shown in the following. The geometry is seen in Fig. 1. A standard polar coordinate system is used, based on the coordinate system shown in the figure. The original solution can be found in [19], [20]. Note that [19] has a sign error in the coefficient named $c_n$ as pointed out in [9]. The problem has also previously been solved for a IBC cylinder [11]. The notation of [11] is used.

The incident TM$_z$ and TE$_z$ polarized plane wave is given by:

$$E_{z \text{TM}}(\rho, \phi, z) = E_0 (\hat{\epsilon} \cos \alpha - \hat{\rho} \sin \alpha) e^{jk_0 (\rho \sin \alpha + \rho \cos \phi \cos \alpha)} \quad (1a)$$

$$E_{z \text{TE}}(\rho, \phi, z) = E_0 \hat{\rho} e^{jk_0 (\rho \sin \alpha + \rho \cos \phi \cos \alpha)} \quad (1b)$$

where $k_0 = \omega / \sqrt{\epsilon_0 \mu_0}$ is the free space wave number with the free space permittivity and permeability given by $\epsilon_0$ and $\mu_0$, respectively. $E_0$ is the amplitude of the incident wave. And $\alpha$ is the angle of incidence as shown in Fig. 1.

Vector potentials for TM$_z$ and TE$_z$ polarized incident waves are introduced with the notation of [21]:

- For a TM$_z$ incident wave: $A_{z \text{TM}} = \hat{z} A_{z \text{TM}}$, $F_{z \text{TM}} = \hat{z} F_{z \text{TM}}$
  
  $$E_1 = -\nabla \times \vec{F}_{z \text{TM}} + \frac{1}{j\omega \mu_0} \nabla \times \nabla \times \vec{A}_{z \text{TM}} \quad (2a)$$
  $$H_1 = \nabla \times \vec{A}_{z \text{TM}} + \frac{1}{j\omega \mu_0} \nabla \times \nabla \times \vec{F}_{z \text{TM}} \quad (2b)$$

- For a TE$_z$ incident wave: $F_{z \text{TE}} = \hat{z} F_{z \text{TE}}$, $A_{z \text{TE}} = \hat{z} A_{z \text{TE}}$
  
  $$E_2 = -\nabla \times \vec{F}_{z \text{TE}} + \frac{1}{j\omega \mu_0} \nabla \times \nabla \times \vec{A}_{z \text{TE}} \quad (3a)$$
  $$H_2 = \nabla \times \vec{A}_{z \text{TE}} + \frac{1}{j\omega \mu_0} \nabla \times \nabla \times \vec{F}_{z \text{TE}} \quad (3b)$$

The total field will be a superposition of these fields:

$$\vec{E} = \vec{E}_1 + \vec{E}_2 \quad (4a)$$
$$\vec{H} = \vec{H}_1 + \vec{H}_2 \quad (4b)$$

It can be shown that the TM$_z$ and TE$_z$ potentials satisfy the same wave equation as the fields [21]:

$$(\nabla^2 + k_0^2) \begin{bmatrix} A_z \\ F_z \end{bmatrix} = 0 \quad (5)$$

The approach in [19] is followed. The $z$ variation is isolated in the term $e^{jk_z z}$ where $k_z = k_0 \sin \alpha$. The incident plane wave is written with the expansion of the $e^{j\rho \cos \phi \sin \alpha}$ term by the use of first-order Bessel functions. This gives the following equations for the fields outside the cylinder:

$$\begin{bmatrix} A_z \\ F_z \end{bmatrix} = e^{jk_z z} \sum_{n=-\infty}^{\infty} j^n e^{-j n \phi} \begin{bmatrix} A_{z} \\ F_{z} \end{bmatrix} \quad (6)$$

where $\phi$ and $z$ are the coordinates for the observation point and $A_z$ and $F_z$ are given by:

$$\begin{bmatrix} A_{z \text{TM}} \\ F_{z \text{TM}} \end{bmatrix} = \begin{bmatrix} C_{0}^n \\ C_{0}^n \end{bmatrix} \left( J_n (k_{10} \rho) - A_n^m A_m^n H_n^{(2)} (k_{10} \rho) \right) \quad (7a)$$
$$\begin{bmatrix} A_{z \text{TE}} \\ F_{z \text{TE}} \end{bmatrix} = \begin{bmatrix} -C_{0}^n \\ C_{0}^n \end{bmatrix} A_n^m A_m^n H_n^{(2)} (k_{10} \rho) \quad (7b)$$

where $\rho$ is the coordinate of the observation point, $k_{10} = k_0 \cos \alpha$ and $\eta_0 = \sqrt{k_0^2}$ is the intrinsic impedance of free
space, \( J_n \) are first-order Bessel functions, \( H_n^{(2)} \) are second-order Hankel functions and:

\[
C_m^{(2)} = E_0 \frac{j}{\eta_0 k_{10}} \quad (8a)
\]

\[
C_0^{(2)} = E_0 \frac{1}{j k_{10}} \quad (8b)
\]

Similar expressions for the fields inside the cylinder can be found in [19], [20]. The tangential fields inside and outside the cylinder are equated at the surface of the cylinder to satisfy the boundary conditions. This gives the following coefficients:

\[
A_m^{(n)} = \frac{J_n(k_{10} a)}{H_n^{(2)}(k_{10} a)} \left( \frac{F_n N_n}{G_n M_n} - q_e^2 \right) \quad (9a)
\]

\[
A_n^{(n)} = \frac{2q_e}{\pi(k_{10} a)} \left( H_n^{(2)}(k_{10} a) \right)^2 \left( F_n G_n - q_e^2 \right) \quad (9b)
\]

with

\[
\begin{bmatrix}
F_n \\
G_n \\
M_n \\
N_n
\end{bmatrix} = \begin{bmatrix}
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a)
\end{bmatrix} - \begin{bmatrix}
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a)
\end{bmatrix} \quad (10a)
\]

\[
\begin{bmatrix}
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a)
\end{bmatrix} = \begin{bmatrix}
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a)
\end{bmatrix} - \begin{bmatrix}
J_{n}^{(2)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(1)}(k_{10} a) \\
J_{n}^{(2)}(k_{10} a)
\end{bmatrix} \quad (10b)
\]

where \( k_{10} = \sqrt{k_{-1}^2 - (k_0 \sin \alpha)^2} \). The wave number of the lossy dielectric \( k_1 \) is given by \( k_1 = \omega \sqrt{\varepsilon_1 \mu_1} \) with the permittivity and permeability of the lossy dielectric given by \( \varepsilon_1 \) and \( \mu_1 \), respectively, and finally,

\[
q_m = \frac{\varepsilon_1 k_{10}}{\eta_0 k_{10}} \quad (11a)
\]

\[
q_e = \frac{\mu_1 k_{10}}{\mu_0 k_{10}} \quad (11b)
\]

\[
q_e = \frac{n k_{10}}{k_{10} \eta_0} \left[ 1 - \left( \frac{k_{10}}{k_1} \right)^2 \right] \quad (11c)
\]

**B. Infinitesimal Dipole Sources**

The canonical problem of the fields from an infinitesimal dipole source on the surface of an infinitely long dielectric cylinder is solved in the following. The focus is on the case with the observation point placed on the surface of the cylinder as well. This makes it possible to calculate the coupling or path loss between two infinitesimal dipoles on the surface. The geometry is defined in Fig. 2. Again a standard polar coordinate system is used. Without loss of generalization, the source point is assumed to be at \( (x = 0, y = 0, z = 0) \). The location of the observation point is then simply given by its coordinates. A similar approach for the impedance boundary conditions has been made in [12], [22], [23] and for a perfect electric conductor in [24]–[26]. The case with the source away from the cylinder is solved in [27]. The same potentials as defined in Eq. 2 and Eq. 3 are used. Furthermore, the following Fourier transforms of the potentials are defined:

\[
A_z = \frac{1}{2\pi} \sum_{n=-\infty}^{\infty} \frac{e^{jn\phi}}{e^{j\eta_0}} \int_{-\infty}^{\infty} A_x e^{jk_z z} dk_z \quad (12a)
\]

\[
F_z = \frac{1}{2\pi} \sum_{n=-\infty}^{\infty} \frac{e^{jn\phi}}{e^{j\eta_0}} \int_{-\infty}^{\infty} F_x e^{jk_z z} dk_z \quad (12b)
\]

where \( \phi \) and \( z \) are the coordinates of the observation point.

By solving the boundary condition on the cylinder for the tangential fields, the potentials from tangential magnetic and electric infinitesimal dipoles can be found. The potentials from the \( z \)-oriented sources become:

\[
\tilde{A}_z = \frac{H_n^{(1)}(k_{10} \rho)}{2\pi k_{10} (F_n G_n - q_e^2) H_n^{(2)}(k_{10} a)} \left( -J_n J_n \frac{\rho}{\eta_0} \right) \quad (13a)
\]

\[
\tilde{F}_z = \frac{H_n^{(2)}(k_{10} \rho)}{2\pi k_{10} (F_n G_n - q_e^2) H_n^{(2)}(k_{10} a)} \left( j\eta_0 q_e J_n - G_n M_n \right) \quad (13b)
\]

where \( \rho \) is the radial coordinate of the observation point, which will be equal to \( a \) in our case. \( J_n \) and \( M_n \) are the \( z \)-oriented electric and magnetic dipole moments, respectively. And the
potentials from the $\phi$-oriented sources are:

\[
\hat{A}_\phi = \frac{H_n^{(2)}(k_0\rho)}{2\pi ak_0(F_0 G_n - q_c^2 H_n^{(2)}(k_0\rho))} \cdot \left[ jk_1 J_n'(k_0\rho) G_n - j\eta k_1 \frac{H_n^{(2)}(k_0\rho)}{k_0 H_n^{(2)}(k_0\rho)} q_e \right] M_\phi \tag{14a}
\]

\[
\hat{F}_\phi = \frac{H_n^{(2)}(k_0\rho)}{2\pi ak_0(F_0 G_n - q_c^2 H_n^{(2)}(k_0\rho))} \cdot \left[ jk_1 J_n'(k_0\rho) G_n - j\eta k_1 \frac{H_n^{(2)}(k_0\rho)}{k_0 H_n^{(2)}(k_0\rho)} q_e \right] M_\phi \tag{14b}
\]

where $J_\phi$ and $M_\phi$ are the $\phi$-oriented electric and magnetic dipole moments, respectively. In the potentials the symbols are defined as for the plane wave incidence, except for $k_z$ that is a variable here. This also effects $k_0$ and $k_1$, which will be given as $k_0 = \sqrt{k_0^2 - k_m^2}$ and $k_1 = \sqrt{k_0^2 - k_a^2}$. The definition of $q_m$, $q_e$ and $q_c$ are the same, but the updated expressions for $k_0$, $k_1$ and $k_m$ should be used.

The potentials from the magnetic and electric infinitesimal dipoles oriented normal to the surface of the cylinder can be found by the use of the reciprocity theorem and the fields from the $z$-oriented sources as shown here:

\[
E_p^z J_p = E_z^z J_z = \frac{1}{j\omega \epsilon_0} k_0^2 A_p^z J_z \tag{15a}
\]

\[
E_p^M J_p = -H_z^M M_z = -\frac{1}{j\omega \mu_0} k_0^2 A_p^M J_z \tag{15b}
\]

\[
H_z^M J_p = -E_z^M M_z = -\frac{1}{j\omega \mu_0} k_0^2 A_p^M J_z \tag{15c}
\]

\[
M_p^M J_p = H_z^M M_z = \frac{1}{j\omega \mu_0} k_0^2 A_p^M E_z J_z \tag{15d}
\]

The potentials from the $\rho$-oriented sources become:

\[
\hat{A}_\rho = \frac{H_n^{(2)}(k_0\rho)}{2\pi ak_0(F_0 G_n - q_c^2 H_n^{(2)}(k_0\rho))} \cdot \left[ jk_1 \frac{H_n^{(2)}(k_0\rho)}{k_0 H_n^{(2)}(k_0\rho)} q_e \right] J_\rho \tag{16a}
\]

\[
\hat{F}_\rho = \frac{H_n^{(2)}(k_0\rho)}{2\pi ak_0(F_0 G_n - q_c^2 H_n^{(2)}(k_0\rho))} \cdot \left[ jk_1 \frac{H_n^{(2)}(k_0\rho)}{k_0 H_n^{(2)}(k_0\rho)} q_e \right] M_\rho \tag{16b}
\]

where $J_\rho$ and $M_\rho$ are the $\rho$-oriented electric and magnetic dipole moments, respectively.
where $\tau_p$ is the $p$th singularity of $A_n^m$, $A_n^r$, and $A_n^t$ given by the $p$th root of

$$D_W(\tau) = \left( W_2(\tau) + jm\tau q \sqrt{1 - \left( \frac{k_0 a + m \tau}{k_1 a} \right)^2} W_2(\tau) \right) \times \left( W_2(\tau) + jm\tau q \sqrt{1 - \left( \frac{k_0 a + m \tau}{k_1 a} \right)^2} W_2(\tau) \right) - \left[ m\tau q \tau W_2(\tau) \right]^2$$

and $D'_W(\tau)$ is the derivative of $D_W(\tau)$ with respect to $\tau$. This can, at $\tau_p$, be expressed as

$$D'_W(\tau_p) = \left( \frac{P_2(\tau_p)}{W_2(\tau_p)} \right)^2 - \frac{j m^2 q \sqrt{1 - \left( \frac{k_0 a + m \tau}{k_1 a} \right)^2}}{W_2(\tau_p)^2}$$

$$W_2(\tau_p) + jm\tau q \sqrt{1 - \left( \frac{k_0 a + m \tau}{k_1 a} \right)^2} W_2(\tau_p)^2$$

$$+ \left( \frac{P_2(\tau_p)}{W_2(\tau_p)} \right)^2 - \frac{2 m^2 q \tau W_2(\tau_p)}{W_2(\tau_p) + jm\tau q \sqrt{1 - \left( \frac{k_0 a + m \tau}{k_1 a} \right)^2} W_2(\tau_p)^2}$$

$$\frac{m^2 q \tau W_2(\tau_p)}{W_2(\tau_p) + jm\tau q \sqrt{1 - \left( \frac{k_0 a + m \tau}{k_1 a} \right)^2} W_2(\tau_p)^2}$$

where the Airy differential equation $W_2'(\tau) - \tau W_2(\tau) = 0$ has been used to simplify the expression. In Eq. 20b the right branch of the square root has to be chosen to ensure that the equal sign is correct. By substitution of $n$ with $v$ and then with $k_0 a + m \tau$ the expressions for $q_c(\tau)$ and $q'_c(\tau)$ can be found as:

$$q_c(\tau) = \left( 1 + \frac{\tau}{2 m^2 \tau} \right) \frac{k_0}{k_2} \left[ 1 - \left( \frac{k_0}{k_1} \right)^2 \right]^2$$

$$q'_c(\tau) = \frac{1}{2 m^2 \tau} \frac{k_2}{k_0} \left[ 1 - \left( \frac{k_0}{k_1} \right)^2 \right]^2$$

The attenuation factor is given by

$$\alpha_p = j m\tau q \frac{\cos \alpha}{a} \tau_p$$

The $\tau_p$’s are found to be approximately located along the line, which originates in origo and makes an angle of $-60^\circ$ with the real axis in the complex plane. Therefore, the $\alpha_p$’s will have an imaginary part, which will effectively increase the wave number. The wavelength along the cylinder is therefore shorter than the free space wavelength.

### B. Infinitesimal Dipole Sources

The same procedure is followed as for the case of plane wave incidence. The expressions for the potentials are simplified by the use of Eq. 17. The Watson transformation is used as well as the Fock substitution from Eq. 18. The integral over $k_2$ is evaluated by the use of the steepest decent method where only the first order approximation in the asymptotic expansion is used. The details of this process are well outlined in [31]. The result for the potentials caused by a $\rho$-oriented electric source are given by:

$$A_x = j \rho \left[ \sqrt{\frac{k_0 \pi}{2j}} \cos \alpha \frac{\pi}{4} \frac{1}{2j k_0 a} H_v^{(2)}(k_0 a) e^{-j k_0 a t - \alpha t} + \left( \frac{j v}{k_0 a} D_{\rho} H_v^{(2)}(k_0 a) + \sin(\alpha) D_{\rho} H_v^{(2)}(k_0 a) \right) \right]$$

$$F_x = j \rho \left[ \sqrt{\frac{k_0 \pi}{2j}} \cos \alpha \frac{\pi}{4} \frac{1}{2j k_0 a} H_v^{(2)}(k_0 a) e^{-j k_0 a t - \alpha t} + \left( \frac{j v}{k_0 a} D_{\rho} H_v^{(2)}(k_0 a) - \sin(\alpha) D_{\rho} H_v^{(2)}(k_0 a) \right) \right]$$

The expressions for the electric and magnetic sources of other orientations have been left out in order to reduce the paper length.

From the potentials the $\rho$-oriented electric field from a $\rho$-oriented electric source is calculated and results in:

$$E_{\rho} = j \rho \left[ \sqrt{\frac{k_0 \pi}{2j}} \cos \alpha \frac{\pi}{4} \frac{1}{2j k_0 a} \frac{1}{k_0 a} H_v^{(2)}(k_0 a) e^{-j k_0 a t - \alpha t} \right]$$

$$\left[ H_v^{(2)}(k_0 a) + \frac{j v}{k_0 a} D_{\rho} H_v^{(2)}(k_0 a) + \sin^2(\alpha) D_{\rho} H_v^{(2)}(k_0 a) \right]$$

### IV. NUMERICAL RESULTS

The eigenfunction solution for the plane wave incidence given in Section II-A has been evaluated numerically. This is done by truncating the summation of Eq. 6. Truncation at $n = k_0 a + 20$ yields a sufficient accuracy.

The GTD solution found in Section III-A has also been evaluated. The most challenging part of the evaluation is to find the zeros of Eq. 21. This can be done by different numerical methods as listed in [11]. Here the zeros have been found by a simple algorithm that searches an area around origo in the complex plane to find minimums of error of Eq. 21. When the requested amount of minimums have been found,
the error for each minimum is reduced below a threshold by zooming in on each of these. The minimums are then considered to be solutions to the equation. Here the number of zeroes or modes used is four. Two modes are more than enough in the deep shadow. The error threshold is set to $10^{-5}$, but less will do.

In Fig. 3 the proposed approximation has been compared to the exact solution, an IBC approximation and a PEC approximation. The results shown are from a cylinder with radius $a = 160 \text{ mm}$, incidence angle of $\alpha = 30^\circ$ and constitutive parameters of fat (see Table I). This could be a helix shaped path around the human torso. The solution frequency is 5.8 GHz. The results for the E-fields on the cylinder from a TE incident wave is seen in Fig. 3a. The results for the H-fields on the cylinder from a TM incident wave is seen in Fig. 3b. The H-field has been chosen for the TM case (b). The solution is for a helix shaped path around the human torso. The solution frequency is 2.45 GHz. The results are for a cylinder with radius $a = 160 \text{ mm}$, incidence angle of $\alpha = 30^\circ$ and constitutive parameters as specified by the standard [18] (see Table I). This corresponds to a typical path on the human head. The solution frequency is 2.45 GHz. All polarizations are seen to be modeled well with less than 1 dB of error in the deep shadow. The errors are a bit larger than the results for the 80 mm cylinder at 2.45 GHz since the electrical length of the 80 mm cylinder at 2.45 GHz is close to the limit of an electrically large structure.

In Fig. 4 the results of the proposed approximation is compared to the exact solution for each of the three polarizations for each of the incident waves. The results are for a cylinder with radius $a = 80 \text{ mm}$, incidence angle of $\alpha = 10^\circ$, and constitutive parameters as specified by the standard [18] (see Table I). This corresponds to a typical path on the human head. The solution frequency is 2.45 GHz. All polarizations are seen to be modeled well with less than 1 dB of error in the deep shadow. The errors are a bit larger than the results for the 160 mm cylinder at 5.8 GHz since the electrical length of the 80 mm cylinder at 2.45 GHz is close to the limit of an electrically large structure.

To investigate the area of validity for the proposed approximation, the relative root-mean-square error for a sweep of the radius of the cylinder $a$ and the angle of incidence $\alpha$ has been...
calculated. The relative error $RE$ was calculated as:

$$RE = \frac{|\vec{E}_{\text{exact}} - \vec{E}_{\text{approx}}|}{|\vec{E}_{\text{exact}}|}$$ \hspace{1cm} (27)

The relative root-mean-square error was found for the $\phi$-angles from $135^\circ$ to $225^\circ$. This is the area where all points are at least $45^\circ$ into the shadow. The cylinder’s constitutive parameters were $\epsilon_r = 39.2$ and $\sigma = 1.8 \text{ S/m}$ as given by the standard [18] for 2.45 GHz. The solution frequency was 2.45 GHz. The decibel value of the RMS-error is shown in Fig. 5. In Fig. 5a and Fig. 5b the error is plotted for the TE and TM cases versus the electrical radius of the cylinder $k_0a$ and the angle of incidence $\alpha$. In Fig. 5c and Fig. 5d the error is plotted for the TE and TM cases versus the electrical curvature of the cylinder $\kappa/k_0$ and the electrical torsion $\tau/k_0$.

It is seen that the approximation is precise for large structures at low angles of incidence. A good rule of thumb for the validity could be that $k_0a > \pi$ or equivalently $a > \lambda_0/2$ and $\alpha < 60^\circ$. If this is converted to general geometric parameters it corresponds to $\frac{\kappa^2 + \tau^2}{\kappa^2} > \lambda_0/2$ and $\tau/\kappa < 2$. The inaccuracy in the paraxial region might be caused by modes, that are not captured by this approximation. The fields in the paraxial region for an IBC cylinder is discussed in [32]. Following the approach outlined in [32] it might be possible to derive accurate results in the paraxial region for a lossy dielectric cylinder as well.

To further investigate the validity of the proposed approximation, the relative root-mean-square error for a sweep of the imaginary part of $-k_{1t}a$ and the radius of
the cylinder $a$ has been calculated. For the approximation in Eq. 17 to be valid the imaginary part of $k_{11} a$ has to be large which is equivalent to the cylinder being opaque. The cylinder’s permittivity was kept constant at $\varepsilon_r = 39.2$, the solution frequency was 2.45 GHz, and the angle of incidence was kept at $\alpha = 0^\circ$. If these parameters are changed the result is only a little different, but the conclusion remains the same. The decibel value of the RMS-error is shown in Fig. 6.

The proposed approximation is seen to be valid as long as the imaginary part of $-k_{11} a$ is large. For smaller cylinders and for TE incident waves the approximations remains valid for even a small imaginary part of around 2. A possible explanation of the improved area of validity for the case of a small cylinders could be the fact that the path through the cylinder is a combination of the reflection on the interfaces and the attenuation through the cylinder. For small cylinders the path through the cylinder is weaker compared to the path around than it is for larger cylinders caused by the initial high loss due to the reflections. The lower loss in the TE creeping waves is believed to be the cause of the improved area of validity for the case of a TE incident wave compared to the case of a TM incident wave. It is seen that for small $k_{11} a$ the approximation breaks down as quantified from Fig. 5.

V. GTD FORMULATION FOR GENERAL GEOMETRY

The expressions for the diffracted electric and magnetic fields on and away from the cylinder can be described by the use of a few constants. The coupling between the TE and TM mode in air and the modes on the cylinder are given by the square root of $D^m_p$ and $D^c_p$. This can be realized by using the first order Debye expansion of the Hankel functions in Eq. 19. Thereby the diffracted field $E^d$ in the shadow part of the cylinder, caused by an incident plane wave $E^i$, can be described as:

$$E^d \sim E^i \cdot T e^{-j k_0 s} e^{-j k_0 t - \alpha_p t}$$  

(28)

where

$$T \sim \sum_{p=1}^\infty \left( b^p D^m_{p,i} + b^i D^c_{p,i} - n^p \tilde{b} D^c_{p,i} + n^i \tilde{b} D^c_{p,i} \right)$$  

(29)

where $(\tilde{t}, \tilde{n}, \tilde{b})$ are the tangential, normal, and binormal unit vectors, respectively, at the source or ingoing grazing point and similarly $(t, n, b)$ are the unit vectors at the observation or outgoing grazing point. The shift in sign of $D^c_p$ is explained by the shift of direction of $\tilde{b}$ dependent on the way of travel. The equation can be rewritten to:

$$T \sim \sum_{p=1}^\infty \left( b^p D^m_{p,i} + n^i \tilde{b} D^c_{p,i} \right) \left( b D^m_{p,o} + n \tilde{b} D^c_{p,o} \right)$$  

(30)

where the ‘$i$’ and ‘$o$’ subscript refers to incoming or outgoing waves with respect to the cylinder. The magnitudes of the constants are given by:

$$|D^m_{p,i}| = |D^m_{p,o}| = \sqrt{|D^m_p|}$$  

(31a)

$$|D^c_{p,i}| = |D^c_{p,o}| = \sqrt{|D^c_p|}$$  

(31b)

where phases are chosen in such a way that:

$$D^m_{p,i} D^m_{p,o} = D^m_p$$  

(32a)

$$D^m_{p,i} D^c_{p,o} = D^c_p$$  

(32b)

$$D^c_{p,i} D^m_{p,o} = -D^c_p$$  

(32c)

$$D^c_{p,i} D^c_{p,o} = D^c_p$$  

(32d)

The fields on the cylinder from an incident plane wave can be described as well. By evaluating Eq. 19 for points on the surface the so-called attachment coefficients can be determined. In the deep shadow on the cylinder the diffracted fields $E^d$ from an incident plane wave $E^i$ will be given by:

$$E^d \sim E^i \cdot S e^{-j k_0 t - \alpha_p t}$$  

(33)

where

$$S \sim \sum_{p=1}^\infty \left( b^p D^m_{p,i} + n^i \tilde{b} D^c_{p,i} \right) \left( i A_{p,i} + b A_{p,b} + n A_{p,n} \right)$$  

(34)

with the attachment coefficient given by (The argument, $k_{10} a$, of the Hankel functions has been suppressed):

$$A_{p,n} = \frac{k_{10} \pi}{2j} \cos \alpha \left( \frac{\nu}{k_{10} a} D^e_{p,o} H_v(2) - j \sin \alpha D^m_{p,o} H_v^{(2)*} \right)$$  

(35a)

$$A_{p,b} = \frac{k_{10} \pi}{2j} \cos \alpha \left[ \frac{\nu}{k_{10} a} D^c_{p,o} H_v^{(2)*} \right]$$  

(35b)

$$A_{p,t} = \frac{k_{10} \pi}{2j} \cos \alpha \left[ -j \cos \alpha D^e_{p,o} H_v^{(2)*} \right]$$  

(35c)

$$A_{p,e} = \frac{k_{10} \pi}{2j} \cos \alpha \left[ \frac{\nu}{k_{10} a} D^m_{p,o} H_v^{(2)*} \right]$$  

(35d)

Equivalently, the diffracted fields $E^d$ in the shadow part of the far-field from a point source $J_1$ on the surface can be found by reciprocity to be:

$$E^d \sim \frac{j k_{10} a}{4\pi} J_1 \cdot T e^{-j k_0 t - \alpha_p t} e^{-j k_0 s}$$  

(36)

where

$$T \sim \sum_{p=1}^\infty \left( b^p L^m_{p,t} + \tilde{n}^p L^e_{p,t} \right) \left( b D^m_{p,o} + n \tilde{b} D^c_{p,o} \right)$$  

(37)
The only thing that remains is to replace source case. This can be seen by comparing to Eq. 26.

The coefficients first determined from the plane wave incidence case matches with the ones from the infinitesimal dipole source case. This can be seen by comparing to Eq. 26.

The behavior of the creeping wave on the body has been described by the use of the unit vectors, which are defined by the geodesic paths. The only thing that remains is to replace \( \alpha \), \( \cos \alpha \) and \( \sin \alpha \) with expressions not specific to the cylinder.

This can be done by the use of:

\[
\begin{align*}
\alpha &= \frac{\kappa}{\kappa^2 + \tau^2} \\
\cos \alpha &= \frac{\kappa}{\sqrt{\kappa^2 + \tau^2}} \\
\sin \alpha &= \frac{-\tau}{\sqrt{\kappa^2 + \tau^2}}
\end{align*}
\]  

where \( \kappa \) and \( \tau \) is the surface ray curvature and torsion, respectively. Furthermore, when the term \( \frac{1}{\sqrt{t}} \) or \( \frac{1}{\sqrt{1+t}} \) appears one should replace it with:

\[
\frac{1}{\sqrt{t}} \int_0^t \frac{d\chi_0}{d\chi} = \frac{d\psi_0}{d\chi}
\]

where \( d\chi \) and \( d\psi_0 \) are the distance between adjacent rays at the observation and source point, respectively. \( d\psi_0 = d\chi_0/t \) is the angle between adjacent rays at the source. Finally, the attenuation will be variable over the surface since the curvature and torsion along a geodesic may change for a general geometry contrary to a cylinder. Therefore, the term \( \alpha_p(t) \) should be replaced by \( \int_0^t \alpha_p(t')dt' \). From these equations the fields in the proximity of the human body or any opaque lossy dielectric object and in the deep shadow can be determined for any convex geometry.

VI. CONCLUSION

A general geometrical theory of diffraction formulation of the on-body propagation was found. The canonical problem of an infinitely long cylinder was solved. It was done for plane wave incidence as well as for a magnetic or an electric infinitesimal dipole source on the surface of the cylinder. The exact solution was transformed to an asymptotic form valid in the deep shadow region of an opaque electrically large human body or lossy dielectric cylinder with radius \( a \). The model was shown to be valid as long as \( a > \lambda_0/2 \) and \( \alpha < 60^\circ \) or \( \frac{\kappa}{\kappa^2 + \tau^2} > \lambda_0/2 \) and \( \tau/\kappa < 2 \). Furthermore, the improvement by the use of this model compared to a PEC
approximation of the human body was significant. Especially, for the TM case and torsional cases. The improvement over an IBC model was smaller, but significant for low dielectric constant tissue such as fat. The asymptotic formulation for the cylinder was generalized to any convex geometry. It was shown that the two different excitations of the cylinder resulted in the same generalized constants. Future work could be to solve other canonical geometries to further validate the model. The relevant cases would especially be an ellipsoid and an elliptical cylinder. A layered model that might describe the human body more accurately could be relevant future work as well.

REFERENCES


Nikolaj Peter Brunvoll Kammersgaard received the B.Sc. and M.Sc. degrees in electrical engineering from the Technical University of Denmark, Lyngby, in 2013 and 2014, respectively. Since 2014, he has been affiliated with GN Hearing A/S, a Danish hearing aid manufacturer, where he is currently pursuing an Industrial Ph.D. degree in cooperation with the Technical University of Denmark. His research interests include on-body antennas and propagation, as well as electrically small antennas and physical limitations of antennas. He was awarded the Student Paper Award for best student paper at the International Workshop on Antenna Technology 2015. In 2015, he was awarded with “Kandidatprisen” from the Danish Association of Engineers (IDA). The award is given to three recently graduated promising electrical engineers for their Master’s thesis.
Søren Helstrup Kvist received the B.Sc., M.Sc. and Ph.D. degrees in electrical engineering from the Technical University of Denmark in 2007, 2009 and 2014, respectively. Since 2009 he has been affiliated with GN Hearing A/S, a Danish hearing aid manufacturer, where he is currently employed as manager of the Radio Systems group. His research interests include on-body antennas and propagation, as well as electrically small antennas and physical limitations of antennas. Dr. Kvist was awarded the ”Elektroprisen 2015” from the Danish Association of Engineers (IDA). The award is given for excellent technical and scientific work within electrical engineering.

Jesper Thaysen received the B.Sc., M.Sc. and Ph.D. degrees in electrical engineering from the Technical University of Denmark in 1998, 2000 and 2005, respectively, and an MBA from Middlesex University in 2015. Since 2008 he has been employed at GN Hearing A/S, a Danish hearing aid manufacturer, where he currently acts as the Head of the Global Manufacturing Organization. His research interests include small antennas and on-body antennas and propagation. Dr. Thaysen has overseen more than 40 B.Sc. and M.Sc. students, as well as 3 Ph.D. students, as the company representative in university-industry cooperative projects.

Kaj Bjarne Jakobsen received the B.Sc.EE and the M.Sc.EE degree from the Technical University of Denmark, Kgs. Lyngby, in 1985 and 1986, respectively, the Ph.D. degree in Electrical Engineering from University of Dayton, Dayton, OH, in 1989, and the HD in Organization and Management, Copenhagen Business School, Copenhagen in 2000. From 1986-1989 he was a Fulbright Scholar at the Department of Electrical Engineering, University of Dayton, OH. Since 1990 he has been with the Department of Electrical Engineering, Technical University of Denmark, Kgs. Lyngby, where he is Associate Professor. His research interests are in body-centric wireless network, wireless body area network, and body sensor network. He received in 1989 the NCR Stakeholder Award, Ohio, USA, and was appointed Teacher-of-the-Year at the Technical University of Denmark in 1994.