The Application of Planar Anistropy to Millimeter-Wave Ferrite Phase Shifters

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Abstract —The torque produced by the planar anistropy that occurs in some hexagonal ferrites is included in the equation of motion of the magnetization. The elements of the permeability tensor, derived here for a lossless approximation, are affected by the planar anistropy in much the same way as by an increase in saturation magnetization. This modified permeability has been incorporated into a model for a planar ferrite loaded rectangular waveguide, and the calculated values for differential nonreciprocal phase shift are found to increase substantially over those for a conventional isotropic ferrite.

I. Introduction

THE TERM "FERRITE" commonly applies to materials with spinel, garnet, and hexagonal crystal structures [1]. Over the past 30 years, these compounds have been studied extensively and their properties optimized for microwave device use [2]. Grain-oriented hexagonal ferrites with uniaxial anistropy have been developed for use in isolators and circulators at millimeter-wave frequencies [3], but no practical applications have yet been realized for hexagonal ferrites with planar anistropy [4], and planar materials are not now commercially available. It appears that the natural attributes of planar ferrites may be used to advantage in switchable, remanent phase shifters. This paper discusses the theory of their application to phasers at millimeter-wave frequencies.

In conventional isotropic ferrite phase shifters, a figure of merit is the ratio of saturation magnetization to operating frequency. Because the upper limit of saturation magnetization is $\sim 4.4\times10^5$ A/m (5500 gauss), the performance of millimeter-wave phasers declines with increasing frequency. While the anisotropy of uniaxial ferrites can replace the applied magnetic fields required for resonance isolators, their use seems impractical for variable phase shifters where the remanent magnetization must be switched and varied to control the phase shift.

Planar ferrites with their plane of easy magnetization oriented perpendicular to the direction of propagation of EM waves seem to have natural advantages in such applications. These materials have not received the same degree of attention accorded isotropic or uniaxial ferrites, and to date the properties of these materials have not been so highly developed.

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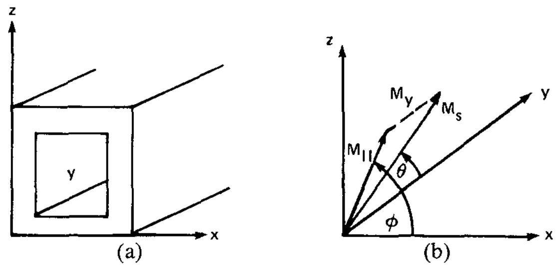

Fig. 1. (a) Planar ferrite toroid. (b) Coordinate system used in describing planar anisotropy.

II. Planar Anisotropy and Tensor Permeability

Magnetic materials all exhibit an angular dependence of their internal energy that is expressed in terms of an anisotropy energy. This energy has a symmetry reflecting that of the unit cell of the particular crystal structure. In hexagonal ferrites, the anisotropy energy may be minimized when the magnetization is parallel to the C-axis (uniaxial anisotropy) or when the magnetization lies in the plane perpendicular to the C-axis (planar anisotropy). This angular dependence of the anistropy energy is expressed phenomenologically by

$$W_{AN} = K_1 \sin^2\theta + K_2 \sin^4\theta + \cdots \tag{1}$$

where K_1 and K_2 are first- and second-order anisotropy constants ($K_1 + 2K_2$ being negative for planar materials) and θ is the angle between the C-axis and the magnetization. Fig. 1 indicates the appropriate coordinate system.

Any directional dependent energy produces a torque given by

$$T_A = \frac{-\partial W_{AN}}{\partial\theta} = -\left(2K_1 \sin\theta\cos\theta + 4K_2 \sin^3\theta\cos\theta\right). \tag{2}$$

For planar ferrites, the anisotropy energy causes the magnetization to remain near the plane described by $\theta = \pi/2$.

Thus

$$\sin\theta \simeq 1.0 \tag{3}$$

$$\cos\theta \simeq \frac{M_y}{M_s} \tag{4}$$

and

$$T_A = -(2K_1 + 4K_2)\frac{M_y}{M_s} \tag{5}$$

where M_y is the y-component of M_s.
From Fig. 1(b), we see that the component of the magnetization in the $x-z$ plane, the hexagonal plane, is given by

$$\vec{M}_{11} = M_s \sin\theta[\cos\phi \hat{a}_x + \sin\phi \hat{a}_z]. \tag{6}$$

The effects of the anisotropy torque are modeled as arising from an effective anisotropy field H_A, also assigned to be in the "easy" plane. Then the torque can be written as

$$\vec{T} = M_s \times \mu_0 \vec{H}_A. \tag{7}$$

Written in the rectangular coordinates of Fig. 1(b), this is

$$\vec{T} = |M_s|[\sin\theta\cos\varnothing\hat{a}_x + \cos\theta\hat{a}_y + \sin\theta\sin\phi\hat{a}_z] \times \mu_0|H_A|[\cos\phi\hat{a}_x + 0\hat{a}_y + \sin\phi\hat{a}_z]. \tag{8}$$

Again, because the magnetization lies very near the easy plane

$$\sin\theta \simeq 1$$

$$\cos\theta \simeq \frac{M_y}{M_s}$$

$$\cos\phi = \frac{M_x}{M_{11}} = \frac{M_x}{M_s \sin\theta} \simeq \frac{M_x}{|M_s|}$$

and

$$\sin\phi = \frac{M_z}{M_{11}} \simeq \frac{M_z}{|M_s|}.$$

Thus

$$\vec{T} = \frac{M_y H_A}{M_s}\mu_0[M_z\hat{a}_x - M_x\hat{a}_z]. \tag{9}$$

The torque of (9) is in the same direction as that of (5), so that

$$\mu_0 M_y H_A = -(2K_1 + 4K_2)\frac{M_y}{M_s}$$

or

$$H_A = \frac{-(2K_1 + 4K_2)}{\mu_0 M_s}. \tag{10}$$

The usual equation of motion of the magnetization is written as

$$\frac{\partial \overline{M}}{\partial t} = \overline{T} = -|\gamma|\overline{M} \times \overline{B} \tag{11}$$

where γ is the gyromagnetic ratio.

The effects of anisotropy can be added through the anisotropy torque as

$$\frac{\partial \overline{M}}{\partial t} = -|\gamma|[\vec{M} \times \mu_0\vec{H} + \vec{T}_A]. \tag{12}$$

In component form for the case of $\overline{M}_s$ along the Z-axis

$$\overline{M} = m_x\hat{a}_x + m_y\hat{a}_y + M_s\hat{a}_z \tag{13}$$

$$\overline{H} = h_x\hat{a}_x + h_y\hat{a}_y + H_0\hat{a}_z \tag{14}$$

$$H_0 = \text{internal dc field.} \tag{15}$$

Inserting (9), (13), and (14) into (12) and assuming exponential time dependence for the RF fields, we obtain

$$j\omega m_x = -|\gamma|\mu_0[m_y H_0 - M_s h_y + H_A m_y] \tag{16}$$

$$j\omega m_y = -|\gamma|\mu_0[M_s h_x - m_x H_0] \tag{17}$$

$$\frac{\partial}{\partial t}m_z = \frac{\partial}{\partial t}M_s \doteq 0 \doteq -|\gamma|\mu_0\left[m_x h_y - m_y h_x - \frac{H_A}{M_s}m_y m_x\right]. \tag{18}$$

Neglecting the higher order terms in time varying quantities, m_x and m_y are found to be

$$m_x = \frac{\mu_0|\gamma|\omega_m(H_A + H_0)}{\Delta}h_x + \frac{j\omega\omega_m}{\Delta}h_y \tag{19}$$

$$m_y = -\frac{j\omega\omega_m}{\Delta}h_x + \frac{\omega_0\omega_m}{\Delta}h_y \tag{20}$$

where

$$\omega_0 = \mu_0|\gamma|H_0 \tag{21a}$$

$$\omega_m = \mu_0|\gamma|M_s \tag{21b}$$

$$\Delta = \mu_0|\gamma|\omega_0(H_0 + H_A) - \omega^2. \tag{21c}$$

The susceptibily matrix is then

$$\overline{M} = [x]h = \begin{bmatrix} \chi_{xx} & \chi_{xy} & 0 \\ \chi_{yx} & \chi_{yy} & 0 \\ 0 & 0 & 0 \end{bmatrix}\begin{bmatrix} h_x \\ h_y \\ h_z \end{bmatrix} \tag{22}$$

where

$$\chi_{xx} = \frac{\mu_0|\gamma|\omega_m(H_A + H_0)}{\Delta} \tag{23a}$$

$$\chi_{xy} = -\chi_{yx} = j\frac{\omega\omega_m}{\Delta} \tag{23b}$$

and

$$\chi_{yy} = \frac{\omega_0\omega_m}{\Delta}. \tag{23c}$$

Note that, for planar anisotropy, the diagonal elements of the susceptibility tensor are not the same. We can define a relative permeability as

$$[\mu] = 1 + [\chi] \tag{24}$$

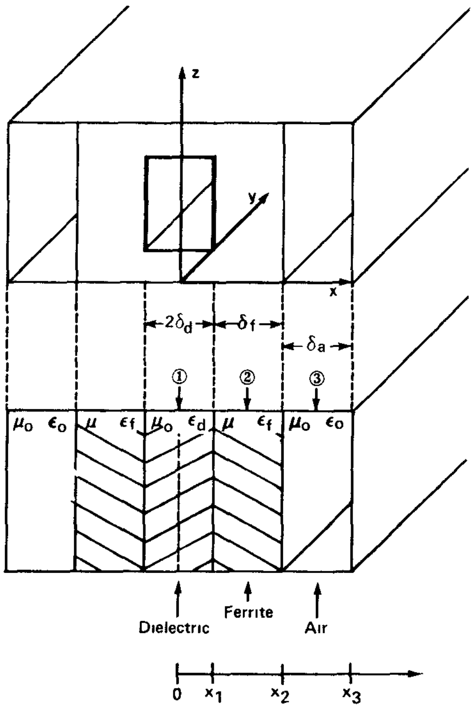

Fig. 2. Rectangular waveguide loaded with ferrite and dielectric.

and thus obtain

$$[\mu]=\begin{bmatrix}\mu_{xx} & j\kappa & 0\\ -j\kappa & \mu_{yy} & 0\\ 0 & 0 & 1\end{bmatrix} \tag{25}$$

where

$$\mu_{xx}=1+\chi_{xx}=1+\frac{\mu_0|\gamma|\omega_m(H_A+H_0)}{\Delta} \tag{26a}$$

$$\mu_{yy}=1+\chi_{yy}=1+\frac{\omega_0\omega_m}{\Delta} \tag{26b}$$

$$\kappa=\frac{1}{j}\chi_{xy}=\frac{\omega\omega_m}{\Delta}. \tag{26c}$$

Unlike the case of isotropic or uniaxial ferrites, $\mu_{xx}\neq\mu_{yy}$ for planar materials.

Planar Ferrite in Rectangular Waveguide

For the coordinate system shown in Fig. 2 and under the assumption that the waveguide is excited by a TE_{10} mode, the H_x, H_y, and E_z components will be nonzero, while $H_z=E_x=E_y=0$. Maxwell's curl equations can then be written in component form as

$$-\Gamma E_z=-j\omega\mu_0\left[\mu_{xx}h_x+j\kappa h_y\right] \tag{27a}$$

$$\frac{-\partial E_z}{\partial x}=-j\omega\mu_0\left[-j\kappa h_x+\mu_{yy}h_y\right] \tag{27b}$$

$$\frac{-\partial h_y}{\partial z}=0 \tag{28a}$$

$$\frac{\partial h_x}{\partial z}=0 \tag{28b}$$

$$\frac{\partial h_y}{\partial x}+\Gamma h_x=j\omega\epsilon E_z. \tag{28c}$$

Here it is assumed that all fields propagate as $e^{-\Gamma y}$ and $\frac{\partial}{\partial y}=-\Gamma$. Through the usual manipulation of these curl equations, they can be cast as a wave equation in E_z

$$\frac{\partial^2 E_z}{\partial x^2}+\left[\frac{\Gamma^2\mu_{yy}}{\mu_{xx}}+\omega^2\mu_0\epsilon\left(\frac{\mu_{xx}\mu_{yy}-\kappa^2}{\mu_{xx}}\right)\right]E_z=0. \tag{29}$$

The results of this derivation can be simplified by defining some of the following terms:

$$\rho=\frac{\mu_{xx}}{\mu_{xx}\mu_{yy}-\kappa^2} \tag{30}$$

$$\theta=\frac{\mu_{xx}}{j\kappa} \tag{31}$$

$$\psi=\frac{\mu_{yy}}{\mu_{xx}}. \tag{32}$$

Now, we can rewrite (29) as

$$\frac{\partial^2 E_z}{\partial x^2}+\left[\Gamma^2\psi+\frac{\omega^2\mu_0\epsilon}{\rho}\right]E_z=0. \tag{33}$$

The magnetic fields can be expressed in terms of the electric field as

$$h_y=\frac{\rho}{j\omega\mu_0}\left[\frac{\partial E_z}{\partial x}+\frac{\Gamma}{\theta}E_z\right] \tag{34}$$

$$h_x=\frac{\rho}{j\omega\mu_0}\left[\Gamma\psi E_z-\frac{1}{\theta}\frac{\partial E_z}{\partial x}\right] \tag{35}$$

where we define

$$k_m^2=\Gamma^2\psi+\frac{\omega^2\mu_0\epsilon}{\rho}. \tag{36}$$

Note that ψ is a measure of asymmetry introduced by planar anisotropy. If $H_A=0$, $\psi=1$.

Equations (33)–(35) may be applied to a dielectric region by noting that for nonmagnetic material

$$\begin{aligned}\omega_m&=0\\ \mu_{xx}&=1\\ \mu_{yy}&=1\\ \kappa&=0\\ \theta&=-j\infty\\ \rho&=1\\ \psi&=1.\end{aligned} \tag{37}$$

For a dielectric region

$$\frac{\partial^2 E_z}{\partial x^2}+\left[\Gamma^2+\omega^2\mu_0\epsilon\right]E_z=0 \tag{38}$$

$$h_y=\frac{1}{j\omega\mu_0}\frac{\partial E_z}{\partial x} \tag{39}$$

$$h_x=\frac{1}{j\omega\mu_0}E_z \tag{40}$$

and

$$k_d^2 = \Gamma^2 + \omega^2 \mu_0 \epsilon. \tag{41}$$

Fig. 2 illustrates a rectangular waveguide loaded with dielectric and ferrite materials. The form of E_z in the three regions is

$$E_{z1} = A(\cos k_d x)e^{-\Gamma y} \tag{42a}$$

$$E_{z2} = \left[Be^{jk_m x} + Ce^{-jk_m x}\right]e^{-\Gamma y} \tag{42b}$$

$$E_{z3} = \left[D \sin k_a (x - x_3)\right]e^{-\Gamma y} \tag{42c}$$

where

$$k_a^2 = \Gamma^2 + \omega^2 \mu_0 \epsilon_0 \epsilon_a \tag{43a}$$

$$k_m^2 = \Gamma^2 \psi + \frac{\omega^2 \mu_0 \epsilon_0 \epsilon_f}{\rho} \tag{43b}$$

$$k_d^2 = \Gamma^2 + \omega^2 \mu_0 \epsilon_0 \epsilon_d. \tag{43c}$$

The transverse operator method utilizes the fact that the fields in the various regions of the waveguide can be related by transfer matrices [5]–[7]. The transfer matrices relate the E_z and h_y fields at one boundary of a medium to those at the other boundary along transverse directions in the waveguide. Thus, we can relate the fields at points x_1 and x_2 by

$$\begin{bmatrix} E_z \\ h_y \end{bmatrix}_{x_1} = \begin{bmatrix} A & B \\ C & D \end{bmatrix} \begin{bmatrix} E_z \\ h_y \end{bmatrix}_{x_2}. \tag{44}$$

Because tangential E and H fields are continuous at boundaries, the transfer matrices can be multiplied to obtain a transfer matrix for the entire waveguide that relates the tangential fields of many regions. Fig. 2 illustrates the geometry for the problem with which we are concerned. Assuming the guide to be symmetric about the center, we need consider only half of the guide as shown. Then

$$\begin{bmatrix} E_z \\ h_y \end{bmatrix}_{\text{wall}} = \begin{bmatrix} A & B \\ C & D \end{bmatrix}_{\text{air}} \begin{bmatrix} A & B \\ C & D \end{bmatrix}_{\text{ferrite}} \begin{bmatrix} A & B \\ C & D \end{bmatrix}_{\text{diel.}} \begin{bmatrix} E_z \\ h_y \end{bmatrix}_{\text{center}} \tag{45}$$

$$\begin{bmatrix} E_z \\ h_y \end{bmatrix}_{\text{wall}} = \begin{bmatrix} A & B \\ C & D \end{bmatrix}_{\text{total}} \begin{bmatrix} E_z \\ h_y \end{bmatrix}_{\text{center}}. \tag{46}$$

Applying the boundary condition, $\hat{n} \times \vec{E} = 0$, at the waveguide wall, we conclude that $E_z(\text{wall}) = 0$. Also, for TE_{10} propagation in rectangular waveguide, the longitudinal component of magnetic field vanishes at the center of the guide, $H_y(\text{center}) = 0$. Thus, we obtain

$$\begin{bmatrix} 0 \\ h_y \end{bmatrix}_{\text{wall}} = \begin{bmatrix} A & B \\ C & D \end{bmatrix}_{\text{total}} \begin{bmatrix} E_z \\ 0 \end{bmatrix}_{\text{center}} \tag{47}$$

or

$$0 = AE_{z_c}$$

$$h_{y_w} = CE_{z_c}.$$

Because $E_{z_{\text{center}}}$ does not vanish, then A must equal zero.

The A, B, C, and D parameters for the transfer matrix of any medium can be obtained by utilizing the equations for E_z and h_y of that particular medium. For a ferrite medium, they are found to be

$$A = \cos k_m \delta_f + \frac{\Gamma}{k_m \theta} \sin k_m \delta_f \tag{48}$$

$$B = \frac{-j\omega\mu_0}{\rho k_m} \sin k_m \delta_f \tag{49}$$

$$C = \frac{j\rho}{\omega\mu_0 k_m}\left[\frac{-\Gamma^2}{\theta^2} - k_m^2\right] \sin k_m \delta_f \tag{50}$$

$$D = \cos k_m \delta_f - \frac{\Gamma}{k_m \theta} \sin k_m \delta_f. \tag{51}$$

In a dielectric region, $\rho = 1$ and $\theta = -j\infty$, and the elements become

$$A = \cos k_d \delta_d \tag{52}$$

$$B = \frac{-j\omega\mu_0}{k_d} \sin k_d \delta_d \tag{53}$$

$$C = \frac{-jk_d}{\omega\mu_0} \sin k_d \delta_d \tag{54}$$

$$D = \cos k_d \delta_d. \tag{55}$$

III. Results

Phase shifters are normally made of a ferrite toroid with a core of dielectric. Fig. 2 illustrates a slab of dielectric sandwiched by slabs of ferrite. This model can be applied to the toroid configuration because the RF fields at the top and bottom of the toroid are essentially parallel to the direction of magnetization. For remanence state phasers, the internal dc field vanishes; thus, $H_0 = 0$ and (26a, b, and c) become

$$\mu_{xx} = 1 - \frac{\mu_0 |\gamma| \omega_m H_A}{\omega^2} \tag{56a}$$

$$\mu_{yy} = 1 \tag{56b}$$

$$\kappa = \frac{-\omega_m}{\omega}. \tag{56c}$$

These approximate representations of the ferrite material can be used in (30)–(32) to determine ρ, θ, and ψ. In turn then, k_a, k_m, and k_d are obtained from (43), and the elements of the transfer matrices from (48)–(55). With given values for widths of air, ferrite, and dielectric regions, the propagation constant Γ is determined to be that value necessary to cause the A element of the overall transfer matrix to vanish.

TABLE I
COMPUTED DIFFERENTIAL PHASE SHIFT FOR VARIOUS FREQUENCIES AND MATERIAL PROPERTIES

Freq. (GHz)	$4\pi M$ (g)	H_A (Oe)	$\Delta\phi$ (deg/cm·GHz)
35	3500	12 000	11.14
35	3500	0	4.67
35	5500	0	7.6
50	3500	12 000	6.22
50	3500	0	3.23
50	5500	0	5.17

$\delta_a = 0.13\lambda_0$, $\delta_d = 0.04\ \lambda_0$, $\delta_f = 0.08\lambda_0$, $\epsilon_f = \epsilon_d = 16$.

As a specific example of the geometry of Fig. 2, consider as typical values

$$f = 35$$
$$a = 0.5\lambda$$
$$b = 0.294\lambda$$
$$\lambda = \frac{c}{f} = \frac{3\times 10^{10}}{35\times 10^{9}} = 0.857 \text{ cm}$$
$$\epsilon_a = 1.0$$
$$\epsilon_d = 16$$
$$\epsilon_f = 16$$

Results of sample calculations are tabulated in Table I. The computed differential phase shift is the difference between the computed propagation constants $\Gamma_{(+)}$ and $\Gamma_{(-)}$ for the two states of magnetization, corresponding to positive and negative signs for κ. Absolute differential phase shift can be obtained from the normalized values presented by multiplying by length (in cm) and operating frequency (in GHz). A planar anisotropy field of 12 000 Oe in a material with $4\pi M = 3500$ g almost triples the computed differential phase shift found for an isotropic ferrite. These results indicate that the planar anisotropy increases differential phase shift as an increase in magnetization does, or as an applied magnetic field would.

It should be noted that this material is still far removed from ferromagnetic resonance and its attendant losses. Resonance would occur at a frequency of about 13.5 GHz, more than 20 GHz below the operating frequency.

Shown in Fig. 3 is a plot of computed differential phase shift as a function of normalized anisotropy field ($MA = \gamma H_A/\omega$) for a fixed geometry and for a frequency of 9 GHz. This was run to verify the computer program with known results for an isotropic ferrite at 9 GHz [7] with $\delta_f = 0.08\lambda_0$, $\delta_d = 0.04\lambda_0$, $\delta_a = 0.13\lambda_0$, and $\epsilon_d = \epsilon_f = 16$. ($\lambda_0$ represents c/f at 9 GHz). The point at $MA = 0$ agrees exactly with results for isotropic ferrites, 6.84°/cm·GHz, or 156°/in at 9 GHz. The apparent effect of the planar anisotropy field is quite large, but only relatively low values of MA have any relevance here. For MA near 4.0, the material undergoes ferromagnetic resonance, and the permeability of a $+cp$ wave becomes negative near $MA = 2.5$. For a 35-GHz frequency, however, a planar anisotropy field of 11 000 Oe corresponds to an MA value of 0.96, which this calculation indicates would more than double

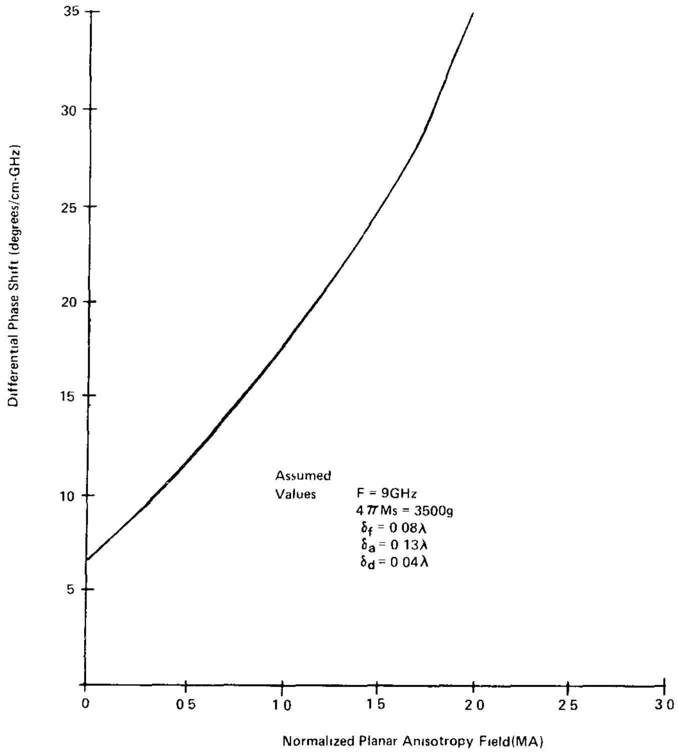

Fig. 3. Computed differential phase shift versus planar anisotropy field.

the differential phase shift. This result is consistent with those of Table I for 35 Ghz.

A planar anisotropy field will also alter the optimum polarization, and therefore the optimum geometry, for differential phase shift. Figs. 4 and 5 show curves representing the computed differential phase shift versus the normalized ferrite width for isotropic and planar ferrites, respectively. The width of the dielectric rib is the parameter in each set of curves. For these curves

$$f_0 = 35 \text{ GHz}$$
$$4\pi M = 3500 \text{ gauss}$$
$$\epsilon_d = \epsilon_f = 16.$$

Waveguide dimensions were set at $a = 0.428$ cm and $b = 0.252$ cm.

The isotropic ferrite (Fig. 4) shows a maximum in differential phase shift at a ferrite width that produces approximately circular polarization of RF magnetic fields in the ferrite region. For the planar material, Fig. 5 indicates further increase in differential phase shift with increasing ferrite width, where the RF magnetic fields will be more elliptically polarized. Of course, a truly optimum geometry could only be determined after the effects of dielectric, magnetic, and ohmic loss are included.

IV. CONCLUSION

The theory developed here and the computer-aided calculations of differential phase shift show that a planar anisotropy field could be used to increase substantially the phase shift in a ferrite-loaded rectangular waveguide. Results also indicate that the planar anisotropy alters the

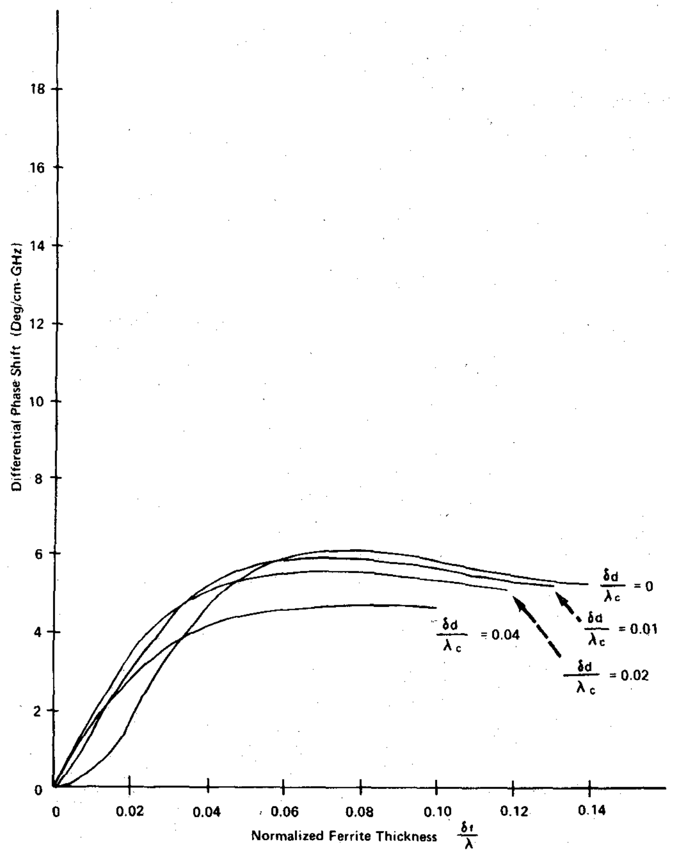

Fig. 4. Differential phase shift versus ferrite thickness with dielectric thickness as a parameter. These curves are for an isotropic ferrite with $4\pi M_s = 3500$ g.

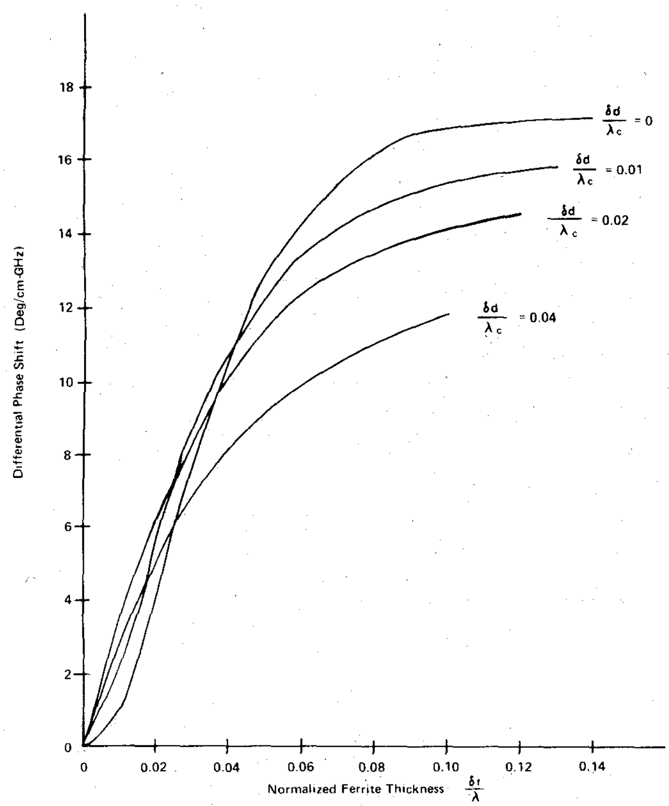

Fig. 5. Differential phase shift versus ferrite thickness with dielectric thickness as a parameter. These curves are for a planar material with $H_A = 12\,000$ Oe and $4\pi M_s = 3500$ g.

polarization condition for optimum performance, and therefore the optimum ferrite and dielectric geometry.

These results offer some promise to meet the needs of millimeter-wave phase shifters where maximum realizable magnetization currently limits performance. However, at this time, planar hexagonal materials are not available on the open market.

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