Surface polaritons on uniaxial antiferromagnets

R. E. Camley
Department of Physics, University of Colorado, Colorado Springs, Colorado 80907

D. L. Mills
Department of Physics, University of California, Irvine, California 92717

(Received 1 February 1982)

We investigate the properties of surface polaritons on uniaxial antiferromagnets. In zero applied field we find that surface polaritons may exist in a band of frequencies forbidden to bulk polaritons and that propagation is reciprocal, i.e., if the mode has frequency ω and wave vector $\vec{k}_{\parallel}$, then $\omega(+\vec{k}_{\parallel})=\omega(-\vec{k}_{\parallel})$. When a field is applied parallel to the surface, along the anisotropy axis, surface and bulk polaritons may exist in the same frequency range, and the dispersion relation of the surface polaritons is no longer reciprocal. We show that surface polaritons may be easily detected through a standard attenuated total reflection measurement. The case of an antiferromagnetic metal is also discussed, since the negative dielectric constant of the metal influences the dispersion relation in a qualitative manner.

I. INTRODUCTION

For many years, it has been known that a long-wavelength surface spin wave may propagate on the surface of a ferromagnetic with magnetization parallel to the surface. This wave, referred to frequently as the Damon-Eshbach surface spin wave,[1] may be described within the framework of a theory which includes the Zeeman energy and magnetic dipole interactions in the spin Hamiltonian; the influence of exchange on its properties has also been explored.[2] The striking and unusual properties of this wave have been studied extensively in recent years by the light scattering method,[3–7] and an excellent account of the data has been provided by theoretical analyses.[2,6–8]

One of us has shown recently[9] that on the surface of uniaxial antiferromagnets, with sublattice magnetization parallel to the surface, long-wavelength magnetostatic spin waves exist with properties similar to those of the Damon-Eshbach wave. For example, the waves propagate only within a limited range of angles, and with a magnetic field applied parallel to the surface their dispersion relation is nonreciprocal, i.e., if $\vec{k}_{\parallel}$ is the wave vector of the mode and $\omega(\vec{k}_{\parallel})$ its frequency, then $\omega(-\vec{k}_{\parallel})\neq\omega(+\vec{k}_{\parallel})$. These modes will lie in the infrared for typical antiferromagnets, rather than in the microwave frequency range appropriate to the ferromagnetic surface magnon described in the preceding paragraph. If we consider a surface magnon with wave vector $\vec{k}_{\parallel}$ parallel to the surface, then in the magnetostatic limit the spin motion associated with the wave extends into the bulk a distance the order of $k_{\parallel}^{-1}$; as a consequence the properties of these deeply penetrating waves may be described by macroscopic theory. They thus constitute a new class of surface magnon on the antiferromagnet, quite distinct in nature from the modes explored some years ago, which require a microscopic theory for their description and which typically have spin motion confined to the outermost few atomic layers.[10–12]

As remarked above, the properties of these antiferromagnetic surface magnons were explored in Ref. 9 within the framework of a magnetostatic theory which, in essence, discards the retardation terms in Maxwell's equations. Such a theory is valid when a $ck_{\parallel}\gg\omega$, where c is the velocity of light. The purpose of the present paper is to explore the influence of retardation on these waves since, at infrared frequencies, experiments may well probe the regime where $ck_{\parallel}$ is not large compared to ω. Indeed, we show that the modes may be excited by means of the attenuated total reflection (ATR) method used widely in the study of surface polaritons on metals and dielectric crystals, and this method probes regions of the dispersion curve where retardation influences it importantly.[13]

We refer to the modes explored in this paper as surface polaritons, since they fall into a class of surface electromagnetic waves often referred to by this nomenclature. The surface polaritons commonly studied on dielectric media have transverse magnetic (TM) character, while the antiferromagnetic surface polaritons of interest here are transverse electric (TE) waves, with magnetic field in the sagittal plane and electric field normal to it. (The modes have pure TE character only for propagation normal to the sublattice magnetizations, actually.) Although the properties of bulk polaritons in antiferromagnets were discussed some time ago[14] and were subsequently studied experimentally in FeF_2,[15] the nature of surface polaritons on the surfaces of these materials has not been explored before. We do note that a general description of surface-polariton propagation on magnetic media has been given earlier,[16] though explicit results were provided only for propagation on ferromagnetic surfaces, where large sample sizes would be required before an approximation to the semi-infinite geometry could be realized in an experiment. Earlier studies have also explored nonreciprocal propagation on gyrotropic media for the semi-infinite geometry, but with primary emphasis on wave guides loaded with ferrites, with geometry suitable for propagation at microwave frequencies.[17] Cylindrical ferrite wave guides have also been analyzed in the microwave literature.[18,19] At the infrared frequencies appropriate for surface-polariton propagation on typical antiferromagnets, the semi-infinite geometry can be readily realized with samples of modest size.

Before we begin, it will be useful to recall briefly the properties of surface polaritons on a semi-infinite, isotropic dielectric. Such a material with a single infrared-active optical phonon may be described by the model dielectric constant

$$\epsilon(\omega)=\epsilon_\infty+\frac{\Omega_p^2}{\omega_{\rm TO}^2-\omega^2}\ , \tag{1.1}$$

where $\omega_{\rm TO}$ is the transverse-optical (TO) phonon frequency, Ω_p^2 a measure of its oscillator strength, and ϵ_∞ the contributions to the dielectric constant from high-frequency excitations, such as electronic interband transitions. Surface polaritons of TM character may propagate on the surface in the frequency regime $\omega_{\rm TO}\le\omega\le\omega_{\rm LO}$, with $\omega_{\rm LO}$ the frequency of the longitudinal-optical (LO) phonon. The dielectric constant vanishes when $\omega=\omega_{\rm LO}$, so the surface waves propagate in a frequency regime, where $\epsilon(\omega)$ is negative, which is a forbidden frequency band for propagation of bulk polaritons.

With the $\hat{z}$ axis aligned along the sublattice magnetization directions, the antiferromagnet is described by a frequency-dependent magnetic permeability tensor of the form[14]

$$\overleftrightarrow{\mu}(\omega)=\begin{bmatrix}\mu_1(\omega) & i\mu_2(\omega) & 0\\ -i\mu_2(\omega) & \mu_1(\omega) & 0\\ 0 & 0 & 1\end{bmatrix}\ , \tag{1.2}$$

where $\mu_2(\omega)$ vanishes in the absence of an external Zeeman field H_0. Thus, if one considers zero external field and propagation of a TE wave perpendicular to the sublattice magnetization axis, so the magnetic field of the surface wave lies in the xy plane, there is a direct analogy between the TE surface polariton on the antiferromagnet and the TM wave on the dielectric. The surface polariton may propagate only in the frequency regime where $\mu_1(\omega)$ is negative and bulk-polariton propagation is forbidden (we assume for the moment that the dielectric constant of the material is positive). One finds $\mu_1(\omega)$ is negative in the frequency regime between

$$\omega_0=\gamma[H_A(2H_E+H_A)]^{1/2}\ ,$$

and

$$\omega_1=(\omega_0^2+8\pi\gamma^2M_sH_A)^{1/2}\ ,$$

with H_A, H_E, and M_s the anisotropy field, the exchange field, and the magnetization of one of the two sublattices.

When a magnetic field is applied parallel to the surface along the anisotropy axis, we find interesting and varied results. Propagation is no longer reciprocal. There are still gaps in frequency where bulk waves do not propagate, and one finds surface solutions in these gaps. However, the surface polaritons are no longer limited to these frequency ranges. There are now frequency regions where both bulk and surface polaritons exist, but with different wave vectors. Also in the presence of an applied field, one finds solutions which have no magnetostatic analog.

In the case of a metallic antiferromagnet $[\epsilon(\omega)<0]$ we also find bulk- and surface-polariton solutions. In this case we find that the phase velocity of both the bulk and surface polaritons is oppositely directed to the group velocity. This agrees with a general discussion of the propagation of bulk electromagnetic waves in materials where both ϵ and μ are negative.[20]

The remainder of the paper is outlined as follows. In Sec. II we present the theory for surface polaritons on an antiferromagnet, and calculate dispersion curves for MnF_2. In addition, results are presented appropriate to a metallic antiferromagnet. In Sec. III we present the ATR calculation, and examples, using realistic parameters, are calculated.

II. THEORY

The geometry considered here is a semi-infinite antiferromagnet lying in the half-space $y>0$. The sublattice magnetizations and anisotropy fields are parallel to the surface along the $+\hat{z}$ and $-\hat{z}$ axes. For simplicity we restrict propagation to be along the $+\hat{x}$ and $-\hat{x}$ axes, perpendicular to the field and parallel to the surface. The geometry is illustrated in Fig. 1.

In the material, Maxwell's equations without sources or currents hold. Thus

$$\nabla\cdot\vec{D}=0\ , \tag{2.1}$$

$$\nabla\times\vec{E}=-\frac{1}{c}\frac{\partial\vec{B}}{\partial t}\ , \tag{2.2}$$

$$\nabla\cdot\vec{B}=0\ , \tag{2.3}$$

$$\nabla\times\vec{H}=+\frac{1}{c}\frac{\partial D}{\partial t}\ . \tag{2.4}$$

The constituitive relations for an antiferromagnet are

$$\vec{D}=\overleftrightarrow{\epsilon}(\omega)\vec{E}\ , \tag{2.5}$$

where

$$\overleftrightarrow{\epsilon}(\omega)=\begin{bmatrix}\epsilon_{\perp}(\omega) & 0 & 0\\ 0 & \epsilon_{\perp}(\omega) & 0\\ 0 & 0 & \epsilon_{\parallel}(\omega)\end{bmatrix}, \tag{2.6}$$

and

$$\vec{B}=\overleftrightarrow{\mu}(\omega)\vec{H}\ , \tag{2.7}$$

where

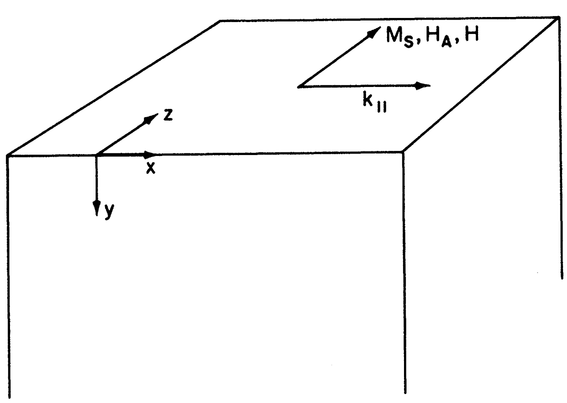

FIG. 1. Geometry considered in this paper. The sublattice magnetizations and anisotropy fields are parallel to the surface and directed along the $+z$ and $-z$ axes. Propagation of the surface polariton is restricted to be along the $+x$ or $-x$ direction.

$$\overleftrightarrow{\mu}(\omega)=\begin{bmatrix}\mu_1(\omega) & i\mu_2(\omega) & 0\\ -i\mu_2(\omega) & \mu_1(\omega) & 0\\ 0 & 0 & 1\end{bmatrix}. \tag{2.8}$$

$\mu_1(\omega)$ and $\mu_2(\omega)$ contain the dynamics of the spin system and may be derived by considering Bloch's equations for each magnetic sublattice. Such a calculation for a uniaxial antiferromagnet gives[14]

$$\mu_1(\omega)=1+4\pi\gamma^2H_AM_S\left[\frac{1}{\omega_0^2-(\omega+\gamma H_0)^2}+\frac{1}{\omega_0^2-(\omega-\gamma H_0)^2}\right], \tag{2.9}$$

$$\mu_2(\omega)=4\pi\gamma^2H_AM_S\left[\frac{1}{\omega_0^2-(\omega+\gamma H_0)^2}-\frac{1}{\omega_0^2-(\omega-\gamma H_0)^2}\right]. \tag{2.10}$$

In the material we look for a transverse electric (TE) surface wave. Thus

$$\vec{E}(\vec{x},t)=\hat{z}E(x,y)e^{-i\omega t}\ . \tag{2.11}$$

The corresponding B field is [from Eq. (2.2)]

$$\vec{B}(\vec{x},t)=\left[-\hat{x}\frac{ic}{\omega}\frac{\partial E(x,y)}{\partial y}+\hat{y}\frac{ic}{\omega}\frac{\partial E(x,y)}{\partial x}\right]e^{-i\omega t}\ . \tag{2.12}$$

In the above equations we have assumed that $\vec{E}$ is a function of x and y, but not z. This is reasonable since the propagation is in the x direction, and we expect the amplitude to decrease in the y direction in order to have a surface wave, but nothing depends on z. Using Eqs. (2.5)–(2.8) and Eq. (2.11) and Eq. (2.12) we obtain the following equation:

$$\left[\frac{\partial^2}{\partial x^2}+\frac{\partial^2}{\partial y^2}+\left[\frac{\mu_1^2-\mu_2^2}{\mu_1}\right]\frac{\omega^2}{c^2}\epsilon_{\parallel}\right]E(x,y)=0\ . \tag{2.13}$$

We now assume for $E(x,y)$ a surface-wave solution with propagation along $\hat{x}$ and decay in amplitude (decay parameter $\alpha_>$) in the $\hat{y}$ direction, the direction into the material:

$$E(x,y)=E^{>}e^{ik_{\parallel}x}e^{-\alpha_> y} \tag{2.14}$$

With the use of Eqs. (2.13) and (2.14) we obtain an equation for the decay parameter $\alpha_>$

$$\alpha_{>}^{2}=k_{\parallel}^{2}-\left[\frac{\mu_1^2-\mu_2^2}{\mu_1}\right]\frac{\omega^2\epsilon_{\parallel}}{c^2} . \tag{2.15}$$

The dispersion relation for bulk polaritons propagating perpendicular to the saturation magnetization may be obtained from Eq. (2.15) by replacing $i\alpha$ by k_y (i.e., assuming a plane-wave solution appropriate to a material infinite in extent). One then has

$$\frac{c^2(k_x^2+k_y^2)}{\omega^2}=\left[\frac{\mu_1^2-\mu_2^2}{\mu_1}\right]\epsilon_{\parallel} . \tag{2.16}$$

Bulk polaritons in the antiferromagnet FeF_2 have recently been studied by Sanders *et al.*,[15] as remarked in Sec. I.

Outside in the vacuum, one has the usual wave equation for $\vec{E}(\vec{x},t)$:

$$\left[\nabla^2+\frac{\omega^2}{c^2}\right]\vec{E}(\vec{x},t)=0 . \tag{2.17}$$

We again look for a surface wave with electric field component only in the z direction:

$$\vec{E}(\hat{x},t)=\hat{z}E^{<}e^{ik_{\parallel}x}e^{+\alpha_{<}y}e^{-i\omega t} . \tag{2.18}$$

From the wave equation, we then find

$$\alpha_{<}^2=k_{\parallel}^2-\frac{\omega^2}{c^2} . \tag{2.19}$$

We now match the solutions inside the material to those outside in the vacuum. Conservation of the tangential components of $\vec{E}$ at the surface plane $y=0$ gives

$$E^{<}=E^{>} . \tag{2.20}$$

To match the tangential components of $\vec{H}$ across the boundary, we need expressions for H_x both inside and outside. From Eqs. (2.2), (2.7), and (2.14) it is easy to show that for $y>0$ one has

$$H_x(x,t)=-\frac{c}{i\omega}\left[\frac{\mu_1}{\mu_1^2-\mu_2^2}\right]\times\left[\alpha_{>}+\frac{\mu_2}{\mu_1}k_{\parallel}\right]E_z(x,t) . \tag{2.21}$$

In the vacuum H_x is given by

$$H_x(x,t)=\frac{c}{i\omega}\alpha_{<}E_z(x,t) . \tag{2.22}$$

Matching H_x across the boundary, we obtain an implicit dispersion relation for the surface polariton:

$$\alpha_{<}+\frac{\mu_1\alpha_{>}+\mu_2k_{\parallel}}{\mu_1^2-\mu_2^2}=0 . \tag{2.23}$$

The terms $\alpha_{<}$ and $\alpha_{>}$ are even in $k_{\parallel}$; μ_1 and μ_2 are independent of it, so it is the term $-\mu_2k_{\parallel}$ in the above equation that causes $\omega(k_{\parallel})$ to differ from $\omega(-k_{\parallel})$. In the absence of an applied field, $\mu_2=0$, and we obtain the simpler relation

$$\frac{c^2k_{\parallel}^2}{\omega^2}=\frac{\mu_1(\epsilon-\mu_1)}{1-\mu_1^2} , \tag{2.24}$$

which is even in $k_{\parallel}$.

We present results below for the dispersion relation for MnF_2. We have used the following parameters: $H_{ex}=550$ kG, $H_A=3.8$ kG, $M_s=0.6$ kG, and $\epsilon=5.5$. The g factor is 2.05, so $\gamma=-1.803\times10^7$ rad/G s. The antiferromagnetic resonance (AFMR) frequency ω_0 is then 65.05 kG or 186 GHz. This frequency is rather low; use of FeF_2 would give an AFMR frequency of 1.58 THz which is well into the far infrared.

Figure 2 shows the dispersion relation for zero applied field. In this figure the bulk polaritons appear as a broad band because we fix only $k_{\parallel}$ for the bulk waves. If we restrict propagation for the bulk waves to be perpendicular to M_s, a range of k_y values are also allowed. This range of values leads to a band of frequencies for the bulk polaritons. We see in Fig. 2 two regions where bulk polaritons may exist, and in the gap between these regions we find surface polaritons.

The surface-polariton curves start at the inter-

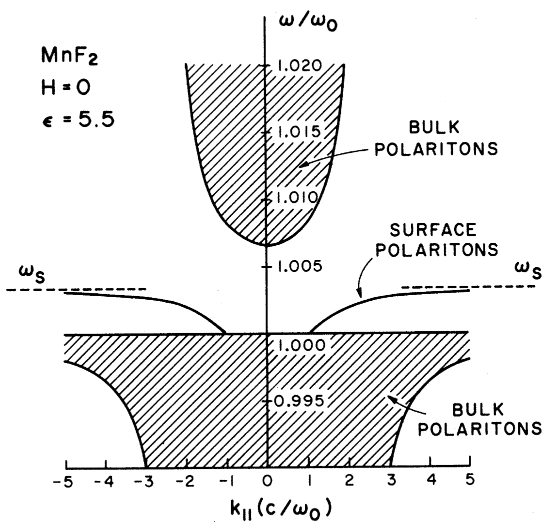

FIG. 2. Dispersion relation for surface polaritons on MnF_2 in zero applied field. ω_s is the frequency of the surface polariton in the magnetostatic limit. The dispersion relation here is reciprocal.

section of the top of the bulk-polariton region and the light line. As $k_{\parallel}$ increases, the curves bend over and approach the large-k magnetostatic limit given by[9]

$$\omega_s = (\omega_0^2 + 4\pi\gamma^2 H_A M_s)^{1/2} . \tag{2.25}$$

We note once again that in this case, the dispersion curves for both bulk and surface polaritons are even in $k_{\parallel}$.

With an applied field, Eq. (2.23) must be solved numerically. In Fig. 3 we show the results for the dispersion relation for MnF_2 in an applied field of 300 G. We see here three frequency regions where bulk polaritons may propagate. There are also three surface-polariton curves. Two of these curves approach the magnetostatic limits in the large $k_{\parallel}$ region, but one mode has no magnetostatic analog. We also see here that the bulk-polariton curves are even in $k_{\parallel}$, while the surface-polariton curves clearly are not. Finally, in contrast to the zero-field case, the surface polaritons are not limited in frequency to the gaps between the bulk-polariton regions.

We have also investigated the case $\epsilon < 0$ which would be appropriate to a metallic antiferromagnet. In Fig. 4 we plot the result for bulk and surface polaritons for this case and in zero applied field. We see in this case that both bulk and surface waves have $d\omega/dk_{\parallel}$ negative, indicating that the group velocity is oppositely directed to the phase velocity. This is, as mentioned earlier, consistent with a general discussion of the propagation of bulk electromagnetic waves in materials where both μ and ϵ are negative.[20]

In Fig. 5 we present the dispersion relation for the bulk and surface polaritons where ϵ is negative and there is an applied field of 250 G. In this case the surface modes again exhibit nonreciprocal behavior. Also we see that in some regions the *surface* polariton may have $d\omega/dk_{\parallel} > 0$ indicating a positive group velocity.

III. THE ATR METHOD

Surface polaritons may be easily probed by an ATR experiment. The geometry of the experiment is illustrated in Fig. 6. The incident beam in the prism propagates at an angle θ with respect to the normal. If θ is greater than the critical angle, when the incident beam reaches the prism—air gap interface the beam will normally be nearly 100% specularly reflected. However, in the air gap there will be exponentially increasing and decreasing fields. If the wave vector parallel to the surface of these fields and the frequency of the fields are close to that of the surface polariton, the exponentially decaying and increasing fields can couple to the surface polariton. In this case, some of the energy of the incident beam is transferred to the surface polariton, and the reflection coefficient drops strongly.

If we fix the frequency of the incident wave to be near the frequency for the surface polaritons, we may adjust the parallel wave vector by changing the angle θ. Thus if one plots the reflection coefficient versus angle, and a large narrow dip is found,

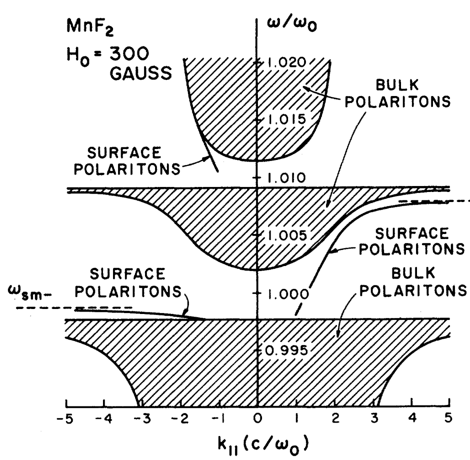

FIG. 3. Dispersion relation for surface polaritons on MnF_2 with an applied field of 300 G. ω_{sm+} and ω_{sm-} are the frequencies of the surface polaritons in the magnetostatic limit. The dispersion relation is no longer reciprocal. Also note near $\omega/\omega_0 = 1.01$ there is a surface mode for which no magnetostatic analog exists.

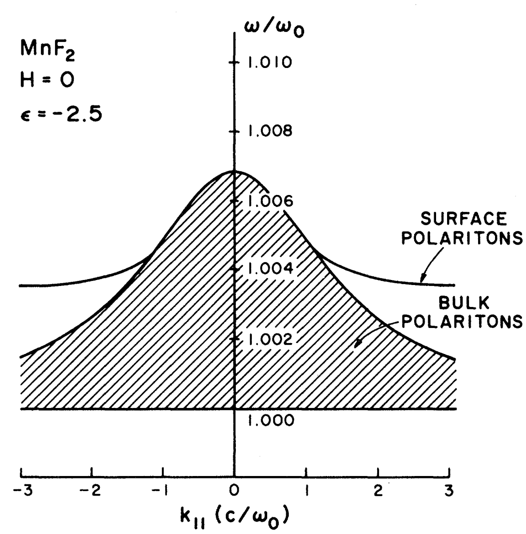

FIG. 4. Dispersion relation for surface polaritons on metallic antiferromagnet in zero applied field. Note that the group velocity is negative for both bulk and surface polaritons.

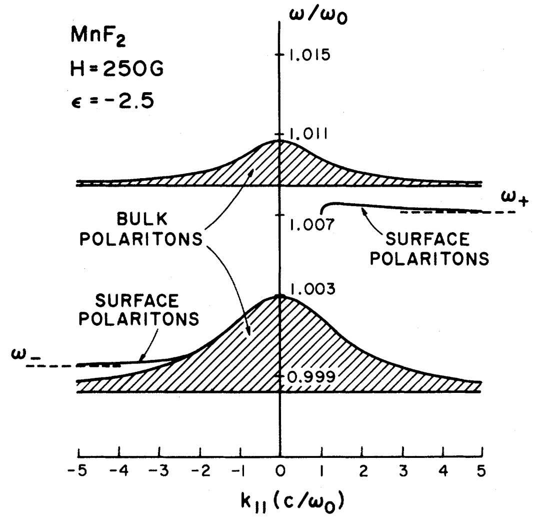

FIG. 5. Dispersion relation for surface polaritons on metallic antiferromagnet in a 250-G field. The surface polariton now exists in the gap between the bulk polariton bands. ω_+ and ω_- are the frequencies in the magnetostatic limit.

this indicates the excitation of a surface polariton. Bulk polaritons may also be detected, although these usually appear as a broad dip.

The calculations for the ATR measurement are similar to those in the preceding section, so we only sketch the derivation below. We assume electric fields of the following form:

region (1),

$$\vec{E}(\vec{x},t)=\hat{z}e^{ik_{\parallel}x-i\omega t}\times(e^{ik_{\perp}(y+d)}+Re^{-ik_{\perp}(y+d)}) , \tag{3.1}$$

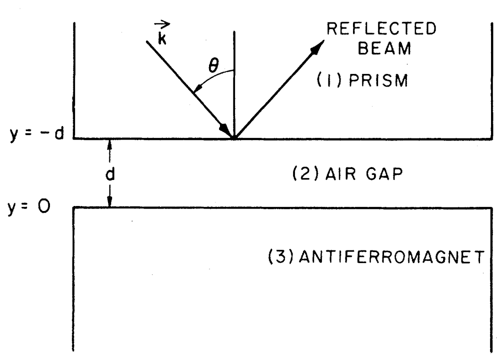

FIG. 6. Geometry for the ATR experiment.

region (2),

$$\vec{E}(\vec{x},t)=\hat{z}e^{ik_{\parallel}x-i\omega t}\times(E_+e^{+\alpha_< y}+E_-e^{-\alpha_< y}) , \tag{3.2}$$

region (3),

$$\vec{E}(\vec{x},t)=\hat{z}e^{ik_{\parallel}x-i\omega t}(E_Ae^{-\alpha_> y}) . \tag{3.3}$$

In the above equations we have assumed the incident beam in region (1) has amplitude unity and the reflected beam has amplitude R. The wave vectors $k_{\parallel}$ and $k_{\perp}$ are related to the incident angle θ by

$$k_{\parallel}=k\sin\theta , \tag{3.4}$$

$$k_{\perp}=k\cos\theta , \tag{3.5}$$

where

$$k=\sqrt{\epsilon_p}\,\omega/c . \tag{3.6}$$

In region (2) $\alpha_<$ is given by Eq. (2.19) and in region (3) $\alpha_>$ is given by Eq. (2.15). From the forms of the electric fields in each region, one may find the expressions for $\vec{H}(\vec{x},t)$ in each region through the use of Maxwell's equations. One then uses the boundary conditions of the continuity of the tangential components of $\vec{E}$ and $\vec{H}$ to match the solutions at the boundaries. In this way one may obtain the following expression for the reflection coefficient R:

$$R=\frac{A-1}{A+1} , \tag{3.7}$$

where

$$A=\left[\frac{Fe^{-2\alpha_< d}+1}{Fe^{-2\alpha_< d}-1}\right]\frac{ik_{\perp}}{\alpha_<} , \tag{3.8}$$

$$F=\frac{\alpha_<-\tilde{\alpha}}{\alpha_<+\tilde{\alpha}} , \tag{3.9}$$

and

$$\tilde{\alpha}=\frac{\mu_1\alpha_>+\mu_2k_{\parallel}}{\mu_1^2-\mu_2^2} . \tag{3.10}$$

In Fig. 7 we plot the magnitude of the reflection coefficient R versus the angle of propagation θ of the incident beam in the prism with respect to the normal. We fix the frequency of the exciting beam to be either $\omega/\omega_0=1.004$ (187.3 GHz) or $\omega/\omega_0=1.005$ (187.5 GHz). The applied field is 300 G, so we are probing the surface polaritons whose dispersion relations are illustrated in Fig. 3. We have also included a phenomenological

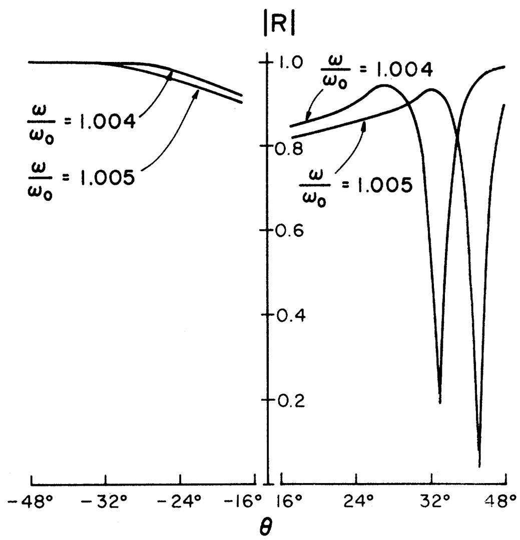

FIG. 7. ATR spectrum: the magnitude of the reflection coefficient R vs angle θ. The sharp dips correspond to excitation of surface polaritons, the broad dips correspond to excitation of bulk polaritons.

linewidth of 20 G for MnF_2. The air-gap thickness is $d = 0.0025$ cm, which is approximately $1/k$.

The angle θ is allowed to be positive or negative. For positive θ, the wave propagates left to right across the surface ($k_{||}$ positive), and for negative θ the wave propagates from right to left ($k_{||}$ negative). Changing the sign of θ is equivalent to reversing the magnetic field. For a given frequency on the $-\theta$ side, we see a broad dip for small θ, and then a level region. The dip corresponds to the loss of energy to bulk polaritons. On the $+\theta$ side, we again see a broad dip for small θ, but for large θ there is a sharp narrow dip corresponding to the excitation of a surface polariton. The nonreciprocity between $+\theta$ and $-\theta$ is just a result of the nonreciprocity of the surface polariton in the presence of an aplied field. In Fig. 3 for $\omega/\omega_0 = 1.005$ there is a surface polariton for $+k_{||}$ but not for $-k_{||}$.

We also see in Fig. 7 that the dispersion curve of the surface polariton can be determined from a series of ATR experiments at different frequencies. When we change the frequency from $\omega/\omega_0 = 1.004$ to $\omega/\omega_0 = 1.005$ (a change of 0.2 GHz) the position of the dip in the reflectivity changes indicating that a surface polariton with different $k_{||}$ is being excited.

In principle, surface polaritons on metallic antiferromagnets could also be detected through an ATR measurement. However, the linewidth for metallic antiferromagnets is likely to be larger than in insulating antiferromagnets, and this may render detection of the dip difficult. We have done calculations for a material with parameters the same as MnF_2 except with $\epsilon = -2.5$. We find that if the linewidth is greater than 60 G one cannot find the dip due to the surface polariton.

Above we have presented results showing that a surface polariton may be detected by an ATR experiment where the applied field and the frequency of the probing beam are held constant, but the angle θ is varied. One may also detect surface polaritons in a geometry where θ and the frequency remain fixed, but the field is changed. In this way one changes ω and $k_{||}$ of the surface polariton by changing the applied field strength. Again, when ω and $k_{||}$ of the exciting beam is close to that of the surface polariton, one obtains a large dip in reflectivity.

IV. SUMMARY

In conclusion we have derived an implicit dispersion relation for surface polaritons on antiferromagnets. This dispersion relation has been solved numerically, and we find have found the following features:

(1) In zero field, surface polaritons exist in frequency ranges forbidden to bulk polaritons. The propagation in zero field is reciprocal.

(2) When a field is applied parallel to the surface along the anisotropy direction, bulk and surface polaritons may exist in the same frequency range, but with different wave vectors. Propagation is no longer reciprocal. When a field is applied, one finds surface-polariton solutions which have no magnetostatic analog.

(3) In metallic antiferromagnets ($\epsilon < 0$) surface polaritons also exist. In most regions both the surface and bulk polaritons have negative group velocity, although under some conditions surface polaritons may have a positive group velocity.

(4) The properties of surface polaritons on antiferromagnets may be easily measured through an ATR experiment.

ACKNOWLEDGMENT

This research has been supported by the Air Force Office of Scientific Research, under Contract No. F49620-78-C-0019. One of the authors (R. E. C.) wishes to acknowledge the University of California, Irvine, where this work was performed.

[1]R. W. Damon and J. R. Eshbach, J. Phys. Chem. Solids 19, 308 (1961).
[2]R. E. Camley and D. L. Mills, Phys. Rev. B 18, 4821 (1979).
[3]P. Grunberg and F. Metawe, Phys. Rev. Lett. 39, 1561 (1977).
[4]J. R. Sandercock and W. Wettling, J. Appl. Phys. 50, 7784 (1979).
[5]A. P. Malozemoff, M. Grimsditch, J. Aboaf, and A. Brunsch, J. Appl. Phys. 50, 5885 (1979).
[6]R. E. Camley and M. Grimsditch, Phys. Rev. B 22, 5420 (1980).
[7]P. Grunberg, M. G. Cottam, W. Vach, C. Mays, and R. E. Camley, J. Appl. Phys. 53, 2078 (1982).
[8]M. G. Cottam, J. Phys. C 12, 1709 (1979).
[9]R. E. Camley, Phys. Rev. Lett. 45, 283 (1980).
[10]D. L. Mills and W. Saslow, Phys. Rev. 171, 488 (1968).
[11]R. E. de Wames and T. Wolfram, Phys. Rev. 185, 752 (1969).
[12]T. Wolfram and R. de Wames, Phys. Rev. 185, 762 (1969).
[13]Extensive discussions of this method may be found in Surface Polaritons; Electromagnetic Waves at Surfaces and Interfaces (North-Holland, Amsterdam, in press).
[14]D. L. Mills and E. Burstein, Rep. Progr. Phys. 37, 817 (1974).
[15]R. W. Sanders, R. M. Belanger, M. Motokawa, and V. Jaccarino, Phys. Rev. B 23, 1190 (1981).
[16]A. Hartstein, E. Burstein, A. A. Maradudin, R. Brewer, and R. F. Wallis, J. Phys. C 6, 1226 (1973).
[17]H. Suhl and L. R. Walker, Bell Syst. Tech. J. 33, 579 (1954); 33, 939 (1954); 33, 1133 (1954).
[18]F. W. Schott, T. F. Tao, and R. A. Freibrun, J. Appl. Phys. 38, 3015 (1967).
[19]T. F. Tao, J. W. Tully, and W. F. Schott, Appl. Phys. Lett. 14, 106 (1969).
[20]V. G. Veselago, Usp. Fiz. Nauk 92, 517 (1964) [Sov. Phys.—Usp. 10, 509 (1968)].