A theoretical investigation of electromagnetic waves obliquely incident upon a plasma slab

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A THEORETICAL INVESTIGATION OF
ELECTROMAGNETIC WAVES OBLIQUELY
INCIDENT UPON A PLASMA SLAB.

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INCIDENT UPON A PLASMA SLAB

A Dissertation
Presented to
The Faculty of the Department of Physics
The College of William and Mary in Virginia

In Partial Fulfillment
of the Requirements for the Degree of
Doctor of Philosophy

By
Calvin T. Swift
August 1969
This dissertation is submitted in partial fulfillment of
the requirements for the degree of
Doctor of Philosophy

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Approved, August 1969

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Dedicated to the memory of my brother,

Frank Wyatt (Babe) Swift

January 16, 1947 - June 19, 1969
I am grateful to Dr. Frederic R. Crownfield, Jr., for his guidance and suggestions during the course of this research, I especially appreciate the constructive comments he expressed each time we met.

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ABSTRACT

The problem of an electromagnetic wave, obliquely incident upon a plasma slab is considered as a boundary-value problem, using a self-consistent solution of the coupled linearized Vlasov and Maxwell equations. Power reflection, transmission, and absorption coefficients are derived under the assumption that all particles undergo specular reflection at the surfaces of the plasma slab. Although our analysis is valid for arbitrary slab thickness, computational results are presented for slabs which are thin compared to a wavelength. The results show that a series of resonances occur which are attributed to the finite temperature of the plasma. The results further show that the resonances are Landau damped as the thermal velocity of the plasma electrons increases. It is shown that similar resonances can be predicted from the coupled linearized hydrodynamic-Maxwell equations; however, as is well known such a model does not predict Landau damping. The effects of a finite collision frequency are then included via a simple B.G.K. collision term. The numerical computations vividly indicate that the resonances undergo severe damping for extremely small values of the collision frequency to signal frequency ratio.

Finally, the plasma capacitor problem is considered, and the results indicate that the longitudinal resonances have very similar characteristics to the plane wave resonances.
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CHAPTER I. INTRODUCTION

The interaction of electromagnetic waves with plasmas has been of continuing interest to those engaged in the study of the ionosphere, of the radar return from meteor trails, and of reentry plasma sheaths. In the last category, most of the research emphasis has been directed toward the solution of boundary-value problems, varying in complexity from relatively simple plane wave\(^1\) interactions to rather complicated ones involving antennas under plasmas\(^2\). With few exceptions\(^3\), a simple "cold" plasma model has been used. A cold plasma is defined here as one in which the/thermal velocity is zero, so that the plasma behaves as an incompressible fluid which exerts no pressure.

One of the major shortcomings of this model is that no mechanism is provided in which to excite/longitudinal plasma waves. The importance of these longitudinal oscillations lies in the fact that they have been observed experimentally in the laboratory, as far back as 1931\(^4\) and later in connection with the radar scattering from cylindrical plasma columns\(^5\). It was observed that the radar return consisted of a series of resonances, the characteristics of which are shown in figure 1. The interesting feature of these resonances is that cold plasma theory predicts only the main resonance which, for the cylindrical column, occurs at \(\omega_p = \sqrt{2}\omega\) (for a plane slab, this resonance occurs when \(\omega = \omega_p\)). Much more comprehensive and convincing experiments were initiated in 1957 by Dattner\(^6\). His experiment consisted of placing a small cylindrical discharge tube into the wall of a waveguide and monitoring the reflection coefficient as a function of
Figure 1.- Radar return from a cylindrical plasma column.
increasing discharge tube current. Because of the thoroughness of the experiment, there is no doubt that the secondary resonances exist, and they have since been termed Tonks-Dattner resonances. Although excellent experimental results were available, a satisfactory analytical explanation for these resonances did not appear until the classic work of Parker, Nickel, and Gould(7) was published. Upon applying a fluid model of the plasma, they were able to conclude that a spectrum of resonances were generated at the frequencies

$$\omega = \omega_0^2 + \left(\frac{3kT_e}{m}\right)k^2$$

where $T_e$ is the electron temperature, $K$ is the Boltzmann constant, $m$ is the electron mass, and $k$ is a wave number which depends on the radius of the plasma column. These frequencies correspond to longitudinal plasma oscillations, which couple more strongly because the phase velocity of such waves is of the order of the thermal velocity of the electrons in the plasma.

One of the major shortcomings of an approach based on the linearized fluid equations lies in the fact that, for finite thermal velocity, the fluid equations are valid only for $\omega \sim \omega_p(7)$. Analytical results based on such a model should therefore become less accurate as the ratio of $\omega_p/\omega$ decreases. Another major shortcoming is the failure of the fluid equations to predict Landau damping(8). The proper description requires a detailed solution of the Vlasov equation, where we identify the Vlasov to be solved self-consistently with Maxwell's equations as the collisionless Boltzmann equation. The importance of the Landau damping lies in the fact that the collisionless damping should be
the ratio of the thermal velocity to
pronounced as/the phase velocity of the longitudinal wave increases. It
will be shown that this ratio increases as the order of the resonance in-
creases, which could account for the damping of the secondary resonances
shown in figure 1. In order to determine how the width of the resonances
at half-maximum behave in detail as a function of all the parameters it is
necessary to solve a boundary-value problem. Our model consists of a
plane wave obliquely incident upon a plasma slab with the electric
vector polarized in the plane of incidence, so that longitudinal plasma
oscillations are excited. This particular problem also has applications
to the study of antennas under reentry plasmas, because the radiation
characteristics of such antennas can be described by a spectrum of plane
waves(9). The kinetic treatment of this problem has previously been
considered by Hinton(10) and by Bowman and Weston(11) in this country
and by Kondratenko and Miroshnichenko(12) in the Soviet Union. Hinton(10)
solved the problem by expressing the currents as integrals over particle
orbits. This procedure is equivalent to solving the Vlasov equation.
The approach, however, requires several ponderous perturbation
expansions and leads to an integral equation solution of the problem.
Bowman and Weston(11), on the other hand, used the singular eigenfunction
techniques of Case(13) Shure(14), Felderhof(15), and Van Kampen/to obtain
solutions to the Vlasov-Maxwell equations. The major disadvantage of
this approach is that analytical and numerical results appear to be
rather difficult to obtain. Kondratenko and Miroshnichenko(11) published
an excellent and concise piece of work. Proceeding as Landau(8) did for
the half-space problem they used an integrating factor to solve the
Vlasov equation. This resulted in a solution in the form of an integral
equation which was reduced by means of a Fourier series. Our treatment
derect from theirs, largely in the initial formulation procedure.

In none of the above papers were numerical results presented.
In fact, the only computations which have appeared, were done by
Melnyk(17), who considered a plasma whose equilibrium statistics are
governed by degenerate Fermi-Dirac statistics. We shall consider
Maxwell-Boltzmann statistics and approach the problem by initially
assuming specular reflection of electrons as the plasma boundaries. This,
we will show, automatically allows us to immediately choose a Fourier
series representation of the problem, and we do not obtain an integral
equation. It is in this way that our formulation differs from that in
reference 11. The usual electromagnetic boundary conditions are used
in connection with the boundary condition of specular reflection. We
then solve for and calculate the reflection, transmission, and absorption
coefficients as functions of the plasma electron density and thermal
velocity for a slab which is thin compared to a free-space wavelength
and for zero collision frequency. A series of resonances, i.e., peaks
in the reflection coefficient, occur which exhibit features of the
Tonks-Dattner resonances, and which become Landau damped as the thermal
velocity of the plasma increases. The reflection coefficient described
by a continuous fluid model of a plasma is also computed; and similar
resonances are noted except that they are not Landau damped.

A kinetic analysis of the plasma capacitor(18) is included to
strengthen our physical deduction concerning the predominance of longi-
tudinal oscillations in our plane wave solution. The results show that
the plasma capacitor, which contains only a longitudinal electric field, resonates at precisely the same slab thickness, plasma frequency, thermal velocity, and propagating frequency as for the plane wave interacting with the slab. These resonances are more conventionally defined in the sense that a peak in resistance, and a zero in reactance is noted at the resonant frequency.

Finally, a finite collision frequency is considered using a simple B.G.K.\(^{(19)}\) collision term and for purposes of nomenclature, we shall continue to refer to the kinetic equation as the Vlasov equation. The results show that the higher-order resonances are completely washed out at such a small value of the ratio of collision to propagating frequency that laboratory reproduction of such resonances would be difficult to achieve at normal radio and microwave frequencies. It is concluded that although the present model exhibits some characteristics of the Tonks-Dattner resonances on a qualitative basis, the detailed structure of the resonances is influenced by another mechanism, probably the inhomogeneity of the plasma.
CHAPTER II. INTERACTION OF A PLANE WAVE WITH A UNIFORM PLASMA SLAB

Figure 2 shows the geometry of the problem. A plane wave is incident upon a plasma slab with the electric vector polarized in the plane of incidence. The incident electromagnetic wave is assumed to have a harmonic dependence of the form \( E = E_0 \exp i(k_0 x \cos \theta + k_0 z \sin \theta - \omega t) \). The faces of the plasma slab are \( x = 0 \) and \( x = L \); the plane of incidence is the \( x-z \) plane, and the angle of incidence is \( \theta \). Here \( k_0 \) is the free-space propagation constant, \( k_0 = \omega/c \) and \( \omega \) is the angular frequency of the incident wave. The case where the electric vector of the incident wave is perpendicular to the plane of incidence is discussed in an analogous manner in a later section.

Kinetic effects, however, depend upon the ratio of thermal velocity to the phase velocity of the plasma waves involved. This ratio is appreciable only for longitudinal plasma waves. These, however, do not couple to incident electromagnetic waves polarized perpendicular to the plane of incidence. The reflection coefficient for the case where the electric vector is parallel to the plane of incidence will be discussed, first for the linearized cold plasma model, then for the linearized fluid model, and finally for the linearized Vlasov equation (with a B.G.K. collision term). The equations describing the plasma in each case are solved self-consistently with Maxwell's equations. In each case, the tangential field components at the left of the slab \( (x < 0) \) are given by:

\[
H_y = H_0 \left[ e^{i k_0 x \cos \theta} + R e^{-i k_0 x \cos \theta} \right] e^{i (k_0 z \sin \theta - \omega t)} \tag{1}
\]
Figure 2.- Geometry of a plane wave obliquely incident upon a plasma slab.
Where \( R \) is the complex reflection coefficient for the magnetic field and \( H_0 \) is the magnetic field amplitude of the incident wave. MKSA units will be used throughout.

To the right of the slab \((x > L)\), we have:

\[
E_z = -H_0 \left[ e^{i k_x \cos \theta} - R e^{-i k_x \cos \theta} \right] \sqrt{\frac{\mu_0}{\varepsilon_0}} \cos \theta e^{i(k_0 z \sin \theta - \omega t)}
\]

(2)

Here \( T \) is the complex transmission coefficient for the magnetic field. The boundary conditions across the surfaces \( x = 0 \) and \( x = L \) require that the \( z \)-dependence of the fields within the slab be the same as those outside; therefore, the fields inside the slab are of the form

\[
E, H = E(x), H(x)e^{i(k_0 z \sin \theta - \omega t)}.
\]

As such, the exponential dependence \( e^{i(k_0 z \sin \theta - \omega t)} \) need not explicitly appear in any of the subsequent expressions.

A. Interaction of a Plane Wave With a Uniform Cold Plasma Slab

If the plasma is cold, the random velocity of the free electrons is assumed to be zero, and the dielectric constant of the plasma can be determined without resorting to kinetic theory. Instead, the equations of motion of a free electron interacting with an electromagnetic wave are solved in order to deduce the polarization per particle, from which
one obtains the following expression for the relative dielectric constant of the plasma:

\[
\frac{\varepsilon}{\varepsilon_0} = 1 - \frac{(\frac{\omega_p}{\omega})^2}{1 + (\frac{\nu}{\omega})^2} + i \frac{(\frac{\nu}{\omega}) (\frac{\omega_p}{\omega})^2}{1 + (\frac{\nu}{\omega})^2}
\]

(5)

where \( \omega_p \) is the plasma frequency, \( \omega_p^2 = \frac{n_0 e^2}{m_0} \) (\( n_0 \) = electron density) and \( \nu \) is the collision frequency for momentum transfer, both of which are assumed to be constant.

The solution to this problem appears in Stratton\(^{(20)}\) but not in a form which will be useful when we compare these results with those that we will obtain for the fluid and kinetic models.

In order to develop the desired solution, we start with the Maxwell curl equations:

\[
\nabla \times \vec{H} = \frac{\partial \vec{D}}{\partial t} = -i \omega \varepsilon \vec{E}
\]

(6)

\[
\nabla \times \vec{E} = -\frac{\partial \vec{B}}{\partial t} = i \omega \mu_0 \vec{H}
\]

(7)

Since the plasma is non-paramagnetic, the permeability \( \mu_0 \) is assumed to be that of free space, and the dielectric constant \( \varepsilon \) given by (5). Using (6) and (7), \( \vec{H} = H_y(x)\hat{u}_y \) (\( \hat{u}_y \) = unit vector in the y-direction) can formally be derived. The result is:

\[
H_y(x) = \frac{H_y(L) \sin k_p x - H_y(0) \sin k_p (x-L)}{\sin k_p L}
\]

(8)
where $H_y(0)$ and $H_y(L)$ are the values of $H_y$ at $x = 0$ and $x = L$, respectively, and $k_p = k_0 \sqrt{\frac{\epsilon}{\epsilon_0}} \sin^2 \theta$ for $\omega_p/\omega < \cos \theta$ and $k_p = ik_0 \sqrt{\sin^2 \theta - \frac{\epsilon}{\epsilon_0}}$ for $\omega_p/\omega > \cos \theta$. Reflection, transmission, and absorption coefficients can now be determined from the boundary conditions, i.e., continuity of tangential $E$ and tangential $H$ at $x = 0$ and $x = L$. Using (1)-(4), (8), and (6), the boundary conditions lead to the following relationships:

\[
(1 + R) H_0 = H_y(0) \tag{9}
\]

\[
H_0 T e^{i k_p L \cos \theta} = H_y(L) \tag{10}
\]

\[
-\sqrt{\frac{\epsilon_0}{\epsilon}} (1 - R) H_0 \cos \theta = -i \omega \mu_0 \left[ H_y(0) G_1 L - H_y(L) G_2 L \right] \tag{11}
\]

\[
-\sqrt{\frac{\epsilon_0}{\epsilon}} H_0 T e^{i k_p L \cos \theta} \cos \theta = -i \omega \mu_0 \left[ H_y(0) G_2 L - H_y(L) G_1 L \right] \tag{12}
\]

It is important to note here that the functions $G_1$ and $G_2$ will be defined separately for (a) the cold plasma model (b) the fluid model and (c) the Vlasov model. In this way, we can use a single algebraic relationship to solve (9)-(12) once $G_1$ and $G_2$ are given. For the cold plasma, we have:

\[
G_1 = \left[ \left( \frac{k_p}{k_0} \right)^2 \frac{1}{\epsilon/\epsilon_0} \right] \cot \frac{k_p L}{k_p L} \tag{13}
\]

\[
G_2 = \frac{G_1}{\cos k_p L} = \left[ \right] \csc \frac{k_p L}{k_p L}
\]
If we also define $Z$ as the surface impedance of the plasma at $x = 0$, we have

$$Z = \frac{i k_0 L}{\cos \theta} G_1 - \frac{G_z^2}{G_1 - i \frac{\cos \theta}{k_0 L}}$$  \hspace{1cm} (15)$$

then it follows that

$$R = \frac{1 - Z}{1 + Z}$$  \hspace{1cm} (16)$$

and

$$T = \frac{(1+R)G_2 e^{-i k_0 L \cos \theta}}{G_1 - i \frac{\cos \theta}{k_0 L}}$$  \hspace{1cm} (17)$$

The absorption coefficient may be defined as:

$$A = 1 - |R|^2 - |T|^2$$  \hspace{1cm} (18)$$

In the absence of collisions ($\nu = 0$) the absorption coefficient of the cold plasma slab is zero.

In the limit as $L \to \infty$, only forward traveling waves exist in the plasma slab, and the expression for the surface impedance reduces to

$$Z_{L \to \infty} = \left( \frac{k_p}{k_0} \right) \frac{1}{\epsilon_0} \frac{1}{\cos \theta}$$ \hspace{1cm} (19)$$
Near \((\omega_p/\omega)^2 = 1\) and \((\nu/\omega) \sim 0\), the dielectric constant approaches zero, and \(G_1\) and \(G_2\) become large so that the surface impedance can approximately be written as:

\[
Z_{\epsilon \to 0} = -\left(\frac{k_p}{k_0}\right)^2 \frac{1}{\epsilon \epsilon_0} \tan \frac{k_p L}{k_0} \cdot \frac{i k_p L}{\cos \theta} \tag{20}
\]

which, for a thin slab \((k_p L \to 0)\) reduces to

\[
Z_{\epsilon \to 0} \bigg|_{L \to 0} = -\frac{i k_0 L}{\cos \theta} \left(\frac{k_p}{k_0}\right)^2 \frac{1}{\epsilon \epsilon_0} \tag{21}
\]

We see that the reflection coefficient of even a thin slab, for which \(\nu = 0\) and \(\epsilon \to 0\), should be unity at \(\omega = \omega_p\) because the impedance becomes infinite. We also note that the limit given by (21) depends upon the order in which the limits are taken with respect to \(\epsilon\) and \(L\). Naturally, for \(L = 0\), \(Z\) goes to one, not zero as implied by (21).

The thin slab will be investigated in more detail later. When the slab is not thin, we see from (13), (14), and (15) that the impedance approaches zero when \(k_p L = (n + \frac{1}{2})\pi\) for \(n = 0, 1, 2, \ldots\). These are Fabry-Perot resonances which are familiar in optics.

Numerical results will be given in section F.

B. Fluid Description of the Plane Wave Problem

The linearized Vlasov equation, with a B.G.K. collision term of the form \(-\nu(f - f_o) = -\nu f_{11}\) (where \(f_o\) is the unperturbed distribution function and \(f_{11}\) is the perturbation) and with \(\partial/\partial t = -i\omega\), may be written as
If the zeroth and first moments/equation (22) are calculated with respect to the particle velocity, the following expressions are obtained (see ref. 1, for example) for conservation of mass and momentum:

\begin{equation}
(-i \omega + v) n_1 + n_0 \nabla \cdot \vec{u} = 0
\end{equation}

\begin{equation}
(-i \omega + v) n_0 \vec{u} = -\frac{n_0 e \vec{E}}{m} - \nabla \cdot \vec{P}
\end{equation}

where \( \vec{u} \) is the fluid velocity, \( \vec{P} \) is the pressure tensor, and \( n_0 \) and \( n_1 \) are the unperturbed electron density, and its perturbation, respectively.

The pressure tensor corresponds to terms in the next higher moment, which can be eliminated by assuming a scalar pressure and using the equation of state

\begin{equation}
P = n K T_e
\end{equation}

\( K \) is the Boltzmann constant and \( T_e \) is the electron temperature) in connection with an adiabatic equation

\begin{equation}
P/n^\gamma = \text{constant}
\end{equation}

If we linearize the pressure term in (24), and eliminate \( n_1 \) using \( P_1 = n_1 K T_e \) the following equations are obtained:

\begin{equation}
m n_0 (\nu - i \omega) \vec{u} = n_0 e \vec{E} - \nabla P_1
\end{equation}

\begin{equation}
a^2 m n_0 \nabla \cdot \vec{u} = i \omega P_1
\end{equation}
Where \( n_0 \) is the unperturbed electron density \( a^2 = \gamma v_T^2 = \gamma kT_e/m \) (\( v_T \) is the thermal velocity) and \( m \) is the electron mass. Equations (27) and (28) are the same ones Wait\(^3\) used to derive the reflection coefficients for a uniform half-space. Our procedure is similar to his, except we have an additional boundary at \( x = L \). The electric field and the fluid velocity are related via the Maxwell equations

\[
\nabla \times \vec{E} = i \omega \mu_0 \vec{H} \tag{29}
\]

\[
\nabla \times \vec{H} = -i \omega \varepsilon_0 \vec{E} + n_0 e \vec{u} \tag{30}
\]

when the last term in (30) is the macroscopic convection current.

Equations (27)-(30) can be used to develop wave equations for \( P_\perp \) and \( \vec{H} \). Since the wave equation is a second-order differential equation, a total of four unknown coefficients must be determined within the slab (two coefficients for \( P_\perp \) and two for \( \vec{H} \)). However, equations (27) and (29) can be used to show that the boundary condition of specular reflection, i.e., \( u = 0 \) at \( x = 0 \), and \( x = L \), implies that

\[
\left. \frac{\partial H_y}{\partial z} \right|_{x=0}^{x=L} = \left. \frac{i \omega \varepsilon_0}{n_0 e} \frac{\partial P_\perp}{\partial x} \right|_{x=0}^{x=L}
\]

the two unknown coefficients for \( P_\perp \) can therefore be expressed in terms of those for \( \vec{H} \), giving the following solutions:

\[
H_y(x) = \frac{H_y(L) \sin k_p x - H_y(0) \sin k_p (x-L)}{\sin k_p L} \tag{31}
\]
\[ P_i(x) = \frac{\omega_p^2}{\omega^2} \frac{m}{\hbar} \sin \theta \left[ \frac{H_y(0) \cos k_m (x-L) - H_y(L) \cos k_m L}{\sin k_m L} \right] \quad (32) \]

where

\[ k_m = \frac{\omega^2}{a^2} \left( 1 - \frac{\omega_p^2}{\omega^2} + i \frac{\nu}{\omega} \right) - k_0 \sin^2 \theta \quad (33) \]

Note that the expression for \( H_y(x) \) is identical to what was obtained for the cold plasma. It also follows that:

\[
E_z = \frac{i \omega \mu_0}{\hbar} \left\{ \begin{array}{c}
H_y(0) \left[ - \frac{k_p \cos k_p (x-L)}{\varepsilon_0 \sin k_p L} - \frac{\omega_p^2}{\omega^2} \frac{k_0^2 \sin^2 \theta \cos k_m (x-L)}{1 + i \frac{\nu}{\omega} \frac{\varepsilon_0}{\varepsilon_0} k_m \sin k_m L} \right] \\
+ H_y(L) \left[ \frac{k_p \cos k_p x}{\varepsilon_0 \sin k_p L} \frac{\omega_p^2}{\omega^2} \frac{k_0^2 \sin^2 \theta \cos k_m L}{1 + i \frac{\nu}{\omega} \frac{\varepsilon_0}{\varepsilon_0} k_m \sin k_m L} \right] \end{array} \right\} \quad (34) \]

Therefore, if we set

\[
G_1 = \left( \frac{k_p}{k_0} \right)^2 \frac{\cos k_p L}{\varepsilon_0 k_p L \sin k_p L} + \frac{\omega_p^2}{\omega^2} \frac{\sin^2 \theta \cos k_m L}{1 + i \frac{\nu}{\omega} \frac{\varepsilon_0}{\varepsilon_0} k_m \sin k_m L} \quad (35) \]
We obtain an expression for the impedance at $x = 0$ identical with the cold plasma result (15), but with our functions $G_1$ and $G_2$ given now by (34) and (35). When $k_u = \infty$, i.e., when $a^2 = 0$, $G_1$ and $G_2$ reduce to those for the cold plasma. By inspecting (35) and (36), we note that for $\nu = 0$, Fabry-Perot type resonances occur when $k_u L = l\pi$

($l = 0, 1, 2, \ldots$) in addition to the cold plasma resonances $\epsilon/\epsilon_0 = 0$ and $k_p L = (n + 1/2)\pi$ ($n = 0, 1, 2, \ldots$). It is interesting to note at this point that the phase velocity can be very low for longitudinal waves, and we can therefore expect resonances for slabs which are thin compared to a free-space wavelength. If the thermal velocity $v_T = a/\sqrt{\gamma}$ is small compared to the speed of light these resonances occur when

$$k_w L = n\pi \approx \frac{k_p L}{\sqrt{\gamma} \frac{v_T}{c}} \sqrt{1 - \frac{\omega_p^2}{\omega^2}}$$

or at a phase velocity $v_{ph} = \frac{k_w L}{n\pi} c$. The phase velocity of longitudinal waves therefore becomes smaller as the slab dimension decreases and as the order of the resonances increases. The quantity $\gamma$ is normally assumed to be three for an electron gas.
Results of calculations for the reflection and absorption coefficients are presented in Section F, and comparisons are made with those obtained from the other models of the plasma.

C. Direct Solution of the Plane Wave Problem Using the Linearized Vlasov Equation

In the last section, we wrote the linearized Vlasov equation, with a B.G.K. collision term of the form \(- v(f - f_0) = - vf_1\) and with \(\partial / \partial t = - i\omega\), as

\[-i\omega f_1 + \vec{v} \cdot \frac{\partial f_1}{\partial \vec{x}} - \frac{\mathbf{E}}{m} \cdot \frac{\partial f_0}{\partial \vec{v}} = -vf_1\]  \hspace{1cm} (22)

where \(f_0\) is the unperturbed distribution function, \(f_1\) is the perturbation of the distribution function, \(\vec{v}\) is the particle velocity, and \(\vec{x}\) is the particle position. Equation (22) is to be solved self-consistently with the Maxwell Curl equations:

\[\nabla \times \vec{H} = -i\omega \varepsilon_0 \vec{E} - e \int f_1 \vec{v} d\vec{v}\]  \hspace{1cm} (38)

\[\nabla \times \vec{E} = i\omega \mu_0 \vec{H}\]  \hspace{1cm} (39)

where the last term in (38) is the density of the convection current.

In order to proceed, the distribution functions for the velocity half-space \(v_x > 0\) and the velocity half-space \(v_x < 0\) are considered separately. If the former is denoted by \(f_1^+(\ldots v_x \ldots)\) and the latter denoted by \(f_1^-(\ldots -v_x \ldots)\) then \(f_1^+\) and \(f_1^-\) satisfy the following equations:

\[(-i\omega + v_x) f_1^+ + v_x \frac{\partial f_1^+}{\partial x} + ik_0 v_x \sin \theta f_1^+ - \frac{e}{m} \left( E_z \frac{\partial f_0}{\partial v_z} + E_x \frac{\partial f_0}{\partial v_x} \right) = 0\]  \hspace{1cm} (40)
If we now introduce $F^+ = f_1^+ + f_1^-$ and $F^- = f_1^+ - f_1^-$, the following second-order differential equation in $x$ is obtained for $F^-:
abla F^-$

$$
\frac{\partial^2 F^-}{\partial x^2} + \left(\omega + i\nu - k_0 \nu \sin \theta\right)^2 F^- - \frac{2i e (\omega + i\nu - k_0 \nu \sin \theta)}{m \nu^2} \frac{d}{d\nu_x} \frac{dE_x}{d\nu_x} = 0 
$$

If all particles are specularly reflected at $x = 0$ and $x = L$, then $F^-$ must vanish at $x = 0$ and $x = L$. This condition can be satisfied identically by a Fourier sine series for $F^-$ as a function of $x$:

$$
F^- = \sum_{l=1}^{\infty} F_0 \left( \nu^2 \right) \sin \frac{l \pi x}{L} 
$$

with

$$
F_0^- = \frac{2}{L} \int_{0}^{L} F^- \sin \frac{l \pi x}{L} \, dx 
$$

Examining (42) and (43), we see that they then imply a Fourier sine series expansion for $E_x$ and a cosine series for $E_z$. If $E_{1x}$, $E_{1z}$, and $H_{1y}$ are the corresponding Fourier coefficients for $E_x$, $E_z$ and $H_y$, we have
\[
F_L = \frac{2ie \omega v_x F_0 \left[ \sqrt{\frac{\mu_0}{\varepsilon_0}} \frac{v_x}{c} H_y - (1 + i \frac{v}{\omega}) E_L \right]}{m v_T^2 \left[ (\omega - k_0 v_x \sin \theta + i\nu)^2 - \left( \frac{q_0 v_x}{L} \right)^2 \right]}
\]

where

\[
f_0 = n_0 F_0 = n_0 e^{-\frac{v^2}{2} v_T^2} \left( \frac{2\pi v_T^2}{4\pi v_T^2} \right)^{3/2}
\]

is the equilibrium distribution function, assumed to be Maxwellian.

Since the x-component of the current is defined as:

\[
J_x = -e \int\int d\nu_x^2 d\nu_x^3 \int_0^\infty v_x F^+ d\nu_x,
\]

the Fourier components of \( J_x \) are

\[
J_{lx} = \frac{i\nu^2}{v_T^2} \frac{\omega_0 e}{L} \frac{\nu}{\nu_T} \left\{ J_{1l} E_{lx} - \sqrt{\frac{\mu_0}{\varepsilon_0}} \frac{J_{2l}}{c} H_y \right\}
\]

where:

\[
J_{1l} = \left( 1 + i \frac{\nu}{\omega} \right) \frac{2L}{L} \int\int d\nu_x^2 d\nu_x^3 \int_0^\infty \frac{\nu_x^2 F_0 d\nu_x}{(\omega - k_0 v_x \sin \theta + i\nu)^2 - \left( \frac{q_0 v_x}{L} \right)^2}
\]

\[
J_{2l} = \frac{2L}{L} \int\int d\nu_x^2 d\nu_x^3 \int_0^\infty \frac{\nu_x^2 v_x F_0 d\nu_x}{(\omega - k_0 v_x \sin \theta + i\nu)^2 - \left( \frac{q_0 v_x}{L} \right)^2}
\]
We have assumed that the unperturbed distribution function is Maxwellian, and we define the two-dimensional form as follows:

\[ \int_{-\infty}^{\infty} \mathrm{d}v_x \ F_0 \ (v_x, v_y, v_z) = F_0 \ (v_x, v_z) \]

The evenness of the integrand allows us to extend the range of integration over the entire real \( v_x \) axis; and we also use a partial fraction expansion to obtain

\[ J_{1,2} = (1 + i \frac{v}{\omega}) \int \int_{-\infty}^{\infty} \frac{\mathrm{d}v_x \ \mathrm{d}v_z \ F_0 \ v_x}{\omega - k_0 v_z \sin \theta - \frac{\omega v_x}{L} + i \nu} \tag{49} \]

\[ J_{2,2} = \int \int_{-\infty}^{\infty} \frac{\mathrm{d}v_x \ \mathrm{d}v_z \ v_x v_z F_0}{\omega - k_0 v_z \sin \theta - \frac{\omega v_x}{L} + i \nu} \tag{50} \]

We can now make use of Fourier transforms in velocity space\(^{(22)}\) to reduce (49) and (50) to single integrals. For example, (49) can be written as a convolution in the form

\[ J_{1,2} = (1 + i \frac{v}{\omega}) \int \int_{-\infty}^{\infty} \frac{\mathrm{d}v_x \ \mathrm{d}v_z \ F_0}{\omega - k_0 v_z \sin \theta - \frac{\omega v_x}{L} + i \nu} \]

\[ \equiv \frac{(1 + i \frac{v}{\omega})}{(2\pi)^2} \int \int_{-\infty}^{\infty} H_1 (\Lambda_x, \Lambda_z) \ H_2^* (\Lambda_x, \Lambda_z) \ d\Lambda_x \ d\Lambda_z \tag{51} \]

where \( \Lambda_x \) and \( \Lambda_z \) are the transform variables of the velocity components \( v_x \) and \( v_z \), respectively, \( H_1 (\Lambda_x, \Lambda_z) \) and \( H_2^* (\Lambda_x, \Lambda_z) \)
are given by:

$$H_1(\Lambda_x, \Lambda_z) = \frac{L}{2\pi} \int_{-\infty}^{\infty} \frac{e^{i(\Lambda_x v_x + \Lambda_z v_z)}}{v_x - \left[ \frac{L\omega}{2\pi} - \frac{v_z}{L} \right] + i \frac{L v}{2\pi}} \, dv_x \, dv_z$$  \hspace{1cm} (52)$$

$$H^*_2(\Lambda_x, \Lambda_z) = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} e^{-i(\Lambda_x v_x + \Lambda_z v_z)\left(-\frac{v_x^2 + v_z^2}{2}v^2\right)} \frac{e^{-(\Lambda_x^2 + \Lambda_z^2)v^2}}{2\pi v^2}$$

$$= e^{-(\Lambda_x^2 + \Lambda_z^2)v^2}$$  \hspace{1cm} (53)$$

Equation (52) can be evaluated using contour integration in the complex $$v_x$$ plane. The integration over $$v_x$$ gives:

$$H_1 = \frac{-2\pi i L}{\omega} \int_{-\infty}^{\infty} dv_x e^{i\Lambda_x v_x} e^{i\Lambda_x(\omega+i\nu)v} \left(1+\frac{L v}{2\pi}\right) e^{-i\Lambda_x \frac{k v}{L} \sin \theta}$$

$$= 0$$  \hspace{1cm} (\Lambda_x > 0)$$

And, the integration over $$v_z$$ gives:

$$H_1 = \frac{-2\pi i L}{\omega} \int_{-\infty}^{\infty} dv_z e^{i\Lambda_z v_z} \left(1+\frac{L v}{2\pi}\right) e^{-i\Lambda_z \frac{k v}{L} \sin \theta}$$

$$= 0$$  \hspace{1cm} (\Lambda_x < 0)$$

Equation (51) therefore reduces to

$$J_{1/2} = -(1+\frac{v^2}{\omega}) \frac{v^2}{L^2} \int d\Lambda_x e^{i\frac{\Lambda_x \omega}{L^2} \left(1+\frac{v^2}{\omega}\right)} e^{-\left(1+\frac{\lambda^2 L^2 \sin^2 \theta}{4\pi^2 v^2}\right) \frac{\Lambda_x^2 v^2}{2}}$$  \hspace{1cm} (56)$$
The dispersion integrals can also be reduced by transforming the x and z velocity components to those perpendicular to and parallel to the wave vector. However, the form of the integral given by (56) allows one to determine by inspection if the integrands of the dispersion integrals are oscillatory and therefore difficult to evaluate. Approximations for small $v_T$ can be made by expanding the Gaussian term in (56).

A similar reduction for $J_{2l}$ yields:

$$J_{2l} = \frac{i l}{\ell \pi} \frac{v_T^4}{r} \sin \theta \int_0^\infty d\Lambda_x e^{i \frac{\Lambda x}{2} \left(1 + i \frac{\omega}{\omega_p}\right)} e^{-\left(1 + \frac{k_0^2 l^2}{\ell^2 \pi^2 \sin^2 \theta}\right) \frac{\Lambda_x^2}{2} \Lambda_x^2}$$

(57)

The Fourier coefficients of the z-component of the current can be determined in an analogous manner, with the result:

$$j_{dz} = \frac{\omega_p^2 \omega_0}{v_T^2} \frac{l}{\ell \pi} \left[ J_{3l} E_{dx} - \sqrt{\mu_0 \epsilon_0} J_{2l} H_{dy} \right]$$

(58)

where

$$J_{3l} = -\left(1 + i \frac{\omega}{\omega_p}\right) \frac{k_0 l}{\ell \pi} \sin \theta \frac{v_T^2}{r} \int_0^\infty d\Lambda_x e^{i \frac{\Lambda x}{2} \left(1 + i \frac{\omega}{\omega_p}\right)} e^{-\left(1 + \frac{k_0^2 l^2}{\ell^2 \pi^2 \sin^2 \theta}\right) \frac{\Lambda_x^2}{2} \Lambda_x^2}$$

(59)

and
Further comments on the integrals, including the expression of the integrals in terms of the plasma dispersion function (23), are made in Appendix A.

A Fourier analysis of the Maxwell Curl equations (38) and (39) gives:

\[
\frac{\ell \pi}{L} E_{Rx} = -i k_0 \sin \theta \, E_{Rx} - i \omega \mu_0 \, H_{Ly} 
\]

(61)

\[
E_{Rx} = \frac{k_0 \sin \theta}{\omega \epsilon_0} \, H_{Ly} + \frac{1}{i \omega \epsilon_0} \, j_{fz}
\]

(62)

\[
H_{Ly} \left[ k_0^2 - k_0^2 \sin^2 \theta - \left( \frac{\ell \pi}{L} \right)^2 \right] + \frac{\ell \pi}{L} \left[ \frac{H_y(0) - (-1)^l H_y(L)}{L/2} \right] = -i k_0 \sin \theta \, j_{fz} - \frac{\ell \pi}{L} \, j_{Rx}
\]

(63)

Substituting the current expressions (44) and (56) into the above set of equations, explicit expressions for the field components may be obtained. The results are:
\[
H_{y} = - \left( \frac{L}{L'} \right) \left[ \frac{H_y(0) - (-1)^l H_y(L)}{L/2} \right] \\
\frac{k_0^2 - k_0^2 \sin^2 \theta - \left( \frac{L}{L'} \right)^2 - k_0^2 \frac{\omega_p^2}{\omega^2} \frac{\omega}{v_T^2} \left( J_{4e} - k_0 L \sin \theta J_{2e} \right)}{L''}
\]

\[
E_{_{_{\omega,0}}} = \frac{L}{\omega \lambda_0 L} \left[ \frac{H_y(0) - (-1)^l H_y(L)}{L/2} \right] \left[ k_0 \sin \theta - \frac{\omega_p^2}{\omega^2} \frac{k_0 L}{L'} \frac{J_{2e}}{L} \right] \\
\left[ 1 - \frac{\omega_p^2}{\omega^2} \frac{L}{L'} J_{2e} \right] \left[ k_0^2 - k_0^2 \sin^2 \theta - \left( \frac{L}{L'} \right)^2 - k_0^2 \frac{\omega_p^2}{\omega^2} \frac{\omega}{v_T^2} \left( J_{4e} - k_0 L \sin \theta J_{2e} \right) \right]
\]

\[
E_{_2} = \frac{i \omega \mu_0}{\omega \lambda_0 L} \left[ \frac{H_y(0) - (-1)^l H_y(L)}{L/2} \right] \left[ 1 - \sin^2 \theta - \frac{\omega_p^2}{\omega^2} \frac{\omega}{v_T^2} \left( J_{5e} \right) \right] \\
\left[ 1 - \frac{\omega_p^2}{\omega^2} \frac{L}{L'} J_{5e} \right] \left[ k_0^2 - k_0^2 \sin^2 \theta - \left( \frac{L}{L'} \right)^2 - k_0^2 \frac{\omega_p^2}{\omega^2} \frac{\omega}{v_T^2} \left( J_{4e} - k_0 L \sin \theta J_{2e} \right) \right]
\]

We recognize from (35)-(35) that

\[
E_{_2} = \sum_{l=0}^{\infty} E_{_{_{\omega,0}}} = -i \omega \mu_0 \left[ H_y(0) G_1 L - H_y(L) G_2 L \right]
\]

where \( G_1 \) and \( G_2 \) are the functions previously introduced to define the surface impedance for the cold plasma and fluid models. From (64), (66), and (67) it follows that \( G_1 \) and \( G_2 \) for the Vlasov model are:
\[ G_1 = \sum_{\ell = 0}^{\infty} \frac{2}{(1 + \delta_0^2)} \left( \frac{1}{k_0 L} \right)^2 \left[ 1 - \sin^2 \theta - \frac{\omega^2}{\omega^2} \frac{\omega}{v_T^2} J_{2 \ell} \right] \]

\[ G_2 = \sum_{\ell = 0}^{\infty} \frac{2}{(1 + \delta_0^2)} \left( \frac{1}{k_0 L} \right)^2 \left[ 1 - \sin^2 \theta - \frac{\omega^2}{\omega^2} \frac{\omega}{v_T^2} J_{2 \ell} \right] \cdot (-1)^\ell \]

where we have defined \( J_{2 \ell} = \left( J_{2 \ell} - \frac{J_{2 \ell}}{c} \sin \theta \right) \frac{\omega L}{2\pi} \). The reflection and transmission coefficients can be derived substituting \( G_1 \) and \( G_2 \) into equations (15), (16) and (17).

We should emphasize again that \( G_1 \) and \( G_2 \) as given above specify the surface impedance of the plasma, and therefore uniquely specify the reflection, transmission, and absorption coefficients for the plane wave problem. We show in Appendix B, that our results reduce to those of the half space in the limit as \( L \to \infty \), as they should.

As the ratio of \( \frac{v_T}{v_{\text{phase}}} \) (\( v_{\text{phase}} \) being the phase velocity of the wave) becomes small, the imaginary parts of the dispersion integrals \( J_{2n} \) are negligible, and the real parts of the dispersion integrals can be expanded in increasing powers of \( \frac{v_T}{v_{\text{phase}}} \), as shown in Appendix A. The results, when applied to (68) and (69) are:
Resonances occur in $G_1$ and $G_2$ when

$$L_{\pi} = \frac{1}{k_0 L} \sqrt{1 - \frac{\omega_p^2}{\omega^2}}$$

(72)

which is approximately the same as the hydrodynamic results for

$\omega_p/\omega \sim 1$. With $\kappa$ defined as $2\pi/L$, we note that the resonances

defined by (72) occur when

$$\frac{v_{\text{phase}}}{v_T} = \frac{\omega}{\kappa v_T} = \frac{k_0 L}{L_{\pi} v_T/c}$$

(73)
Thus the kinetic effect, i.e. Landau damping, should become more pronounced as the ratio of \( \frac{v_{\text{phase}}}{v_T} \) becomes smaller. Therefore, for fixed \( \frac{v_T}{c} \) and fixed \( k_c L \), the Landau damping should become more severe as the order of the resonance increases. Numerical results of the reflection and absorption coefficients are discussed in Section F.

D. The Kinetic Results for the Electric Field Perpendicular to the Plane of Incidence

We have indicated that most of the interesting effects associated with a plane wave obliquely incident upon a plasma layer occur with longitudinal plasma waves excited in the plasma. As such, the case of the electric vector perpendicular to the plane of incidence \( \vec{E} = E_y \hat{u}_y \) in figure 1) is only of secondary interest. However, we include the results here, for completeness. If we proceed as in Section C, we find that the Fourier coefficients of \( F^- \) are:

\[
F_e^- = -\frac{2i\omega^2}{c} \frac{k_y^2}{\omega} \frac{F_0 v_y v_x H_{ex}}{[(\omega - k_x v_x \sin \theta + i\nu)^2 - (\frac{2\pi v_x}{L})^2]}
\]

(74)

It further follows from the Maxwell Curl equations that:

\[
\nabla^2 H_x + k_0^2 H_x = -\frac{\partial j_x}{\partial x}
\]

(75)
\[
\frac{\partial^2 \zeta}{\partial x^2} + \left(k_0^2 - k_0^2 \sin^2 \theta\right) \zeta = -ie \int d\nu^z \frac{v_y}{v_x} \frac{\partial F_i^+}{\partial x}
\]

\[= -ie \int d\nu^z \frac{v_y}{v_x} (\omega - k_0 v_\perp \sin \theta) F_i^-
\]

(76)

A Fourier expansion of (76) gives

\[
\zeta = \sum_{k=1}^{\infty} k_0 \frac{L\pi}{L/2} \left[ \frac{\zeta(0) - (-1)^k \zeta(L/2)}{1 - \sin^2 \theta - \left(\frac{\omega}{k_0 L}\right)^2 - \omega \frac{\omega^2}{\omega^2} J_x} \right]
\]

(77)

where

\[
J_x = \int \int \frac{dv_x dv_y F_0}{\omega - k_0 v_\perp \sin \theta - \frac{\omega}{k_0 L} + i \nu}
\]

\[= -i \frac{L\pi}{L} \int d\lambda_x e^{-\frac{\lambda_x^2}{2} \left(1 + \frac{k_0^2 L^2}{\omega^2 \sin^2 \theta} \right)} e^{i \lambda_x \frac{\omega L}{\omega^2} (1 + \frac{i \nu}{\omega})}
\]

(78)

If we set

\[E_y = E_0 \left[ e^{ik_0 x \cos \theta} + R^T e^{-ik_0 x \cos \theta} \right] e^{i (k_0 z \sin \theta - \omega t)}
\]

for \(x < 0\) and

\[E_y = E_0 T^T e^{i (k_0 x \cos \theta + k_0 z \sin \theta - \omega t)}
\]

for \(x > L\), where \(R^T\) and \(T^T\) are the reflection and transmission coefficients for the electric field perpendicular to the plane of
incidence. If we complete the boundary-value problem, we find that

\[
\frac{1 + R^\perp}{1 - R^\perp} = Z = -i k_o L \cos \theta \left[ G_{2}^{\perp} - \frac{(G_{2}^{\perp})^2}{G_{1}^{\perp} + \frac{i}{k_o L \cos \theta}} \right]
\]

(79)

where \( Z \) is the surface impedance, and the functions \( G_{1}^{\perp} \) and \( G_{2}^{\perp} \) are given by:

\[
G_{1}^{\perp} = \left( \frac{1}{k_o L} \right)^2 \sum_{\ell=0}^{\infty} \frac{2}{1 + \delta_{\ell}^2} \frac{1}{1 - \sin^2 \theta - \left( \frac{\ell \pi}{k_o L} \right)^2 - \omega \frac{\omega_o^2}{\omega^2} J_{\ell}}
\]

(80)

\[
G_{2}^{\perp} = \left( \frac{1}{k_o L} \right)^2 \sum_{\ell=0}^{\infty} \frac{2}{1 + \delta_{\ell}^2} \frac{(-1)^\ell}{1 - \sin^2 \theta - \left( \frac{\ell \pi}{k_o L} \right)^2 - \omega \frac{\omega_o^2}{\omega^2} J_{\ell}}
\]

(81)

Equation (79) is slightly different in form from what we previously obtained for the case of perpendicular incidence; this is because \( R^\perp \), in this case, is the ratio of the reflected to the incident electric field.

E. Relationship to the Plasma Capacitor Problem

We have already pointed out that while the ratio of \( v_{T}/v_{\text{phase}} \) is ordinarily negligible for transverse waves, \( p \)-polarized waves might be expected to show interesting kinetic effects, as we have shown. As further support of this supposition, it is useful to compare the plane
wave calculations with corresponding computations in which only longitudinal plasma oscillations are excited. The plasma capacitor provides this for us.

The problem under consideration consists of a plane-parallel plasma-filled capacitor, whose plates are located at \( x = 0 \) and \( x = L \). An electric field, oscillating at an angular frequency \( \omega \), is applied normal to the plates. In this section we show that the capacitor exhibits resonance behavior at the same values of \( k_0 L, \frac{(\omega_p/\omega)^2}{v_p^2} \), and \( (v_T/v_{\text{phase}}) \). In order to affect this, we will solve for the impedance of the capacitor, as did Hall (18) and Shure (29).

If we set \( E_z \) and \( v \) equal to zero in (40) and (41), and proceed in a manner analogous to that which led to the expression for the \( x \)-component of the current density (equation (46)), we find that the Fourier coefficients of the current density in the capacitor are:

\[
\hat{j}_x = \frac{i\omega_p^2 \omega \varepsilon_0}{v_T^2} \frac{L}{\ell \pi} \hat{J}_{12} E_x \tag{82}
\]

which is nothing more than (46) with \( H_{1y} = 0 \). The continuity equation, relating charge density and current density, gives:

\[
\frac{\nabla}{\Lambda} = j_x - i\omega \varepsilon_0 E_x = \text{constant} \tag{83}
\]

where \( I/\Lambda \) is the current per unit area on the plate of the capacitor. A Fourier expansion of (83) gives:
\[ \frac{I}{A} \cdot \frac{2}{L} \left[ 1 - \frac{(-1)^{\ell}}{\ell \pi} \right] = i \omega \varepsilon_0 \left[ 1 - \frac{\omega_p^2}{\nu_T^2} \frac{L}{\ell \pi} J_{1, \ell} \right] E_0 \]  

(84)

Since the voltage between the plates is given by \( V = -\int_0^L E_x dx \), we find that the impedance of the capacitor is given by

\[ Z = -\frac{1}{i \omega C} \left( \frac{2}{\pi} \right)^2 \sum_{\ell=1}^{\infty} \frac{\left[ 1 - (-1)^{\ell} \right]}{\ell^2 \left[ 1 - \frac{\omega_p^2}{\nu_T^2} \frac{L}{\ell \pi} J_{1, \ell} \right]} \]  

(85)

where \( C \) is the capacity in the absence of a plasma. If we identify \( \Lambda^+ = 1 - \frac{\omega_p^2}{\nu_T^2} \frac{L}{\ell \pi} J_{1, \ell} \), the above expression becomes identical to that obtained by Shure\(^{24}\). For \( \nu_T/\nu_{\text{phase}} \ll 1 \), (85) reduces to

\[ Z \approx \frac{1}{i \omega C} \left( \frac{2}{\pi} \right)^2 \sum_{\ell=1}^{\infty} \frac{\left[ 1 - (-1)^{\ell} \right]}{\ell^2 \left[ 1 - \frac{\omega_p^2}{\omega^2} - \frac{3 \omega_p^2}{\omega^2} \left( \frac{\nu_T L \pi}{\omega L} \right)^2 \right]} \]  

(86)

An inspection of the denominator of the sum in (86) shows that resonances occur when

\[ \ell \pi = \frac{k_0 L}{\sqrt{3} \nu_T / \omega} \sqrt{1 - \frac{\omega_p^2 / \omega^2}{\omega_p / \omega}} \]  

(87)

Equation (87) is identical to equation (72), which defined the resonance condition for a plane wave incident upon the slab. We should therefore expect similarities between the plane wave and capacitor results. These similarities will be exhibited in the form of numerical calculations in the next section.
F. Numerical Results

In order to clearly delineate the differences between using cold, fluid, and kinetic models for a plane electromagnetic wave obliquely incident upon a plasma layer, computations of the reflection coefficient for each model were made as a function of \((\omega_p/\omega)^2\) with the following parameters fixed:

\[
\begin{align*}
\theta &= 15^\circ \\
k_0L &= 0.1 \\
v/\omega &= 0 \\
v_T/c &= 10^{-3} \\
(\nu_T/\nu_{\text{phase}} &= 0.157)
\end{align*}
\]

The results are shown in figure 3. We note that for the cold plasma (fig. 3(a)), only one resonance occurs, and that is located at \(\omega = \omega_p\). In the fluid limit (fig. 3(b)), a series of resonances occur, corresponding to the zeros of \(\sin\left(\frac{\omega}{\omega_p} \sqrt{1 - \frac{\omega_p^2}{\omega^2} L}\right)\). We also note that the main resonance is displaced from \((\omega_p/\omega)^2 = 1\), and that each higher order resonance is narrower than its predecessor. The results obtained from the Vlasov equation are similar to those obtained via the hydrodynamic equations, except that the resonances do not occur at the same values of \((\omega_p/\omega)^2\). The departure becomes more pronounced with increasing order, but this is to be expected since the fluid approximation becomes less valid (for example, compare eqs. (37) and (72)). We also note that a resonance is associated with each odd term of the Fourier expansion.

In order to demonstrate the effects of Landau damping, the \(l = 5\) resonance was investigated for additional values of the parameters.
Figure 3. - Reflection coefficient of a plane wave incident upon a plasma slab (a) cold plasma theory (b) linearized fluid model (c) model based on the linearized Vlasov equation.
The reflection and absorption coefficients are shown in figure 4 for various values of $v_T/v_{\text{phase}}$, where the abcissa has been grossly exaggerated. For $v_T/v_{\text{phase}} = 0.157$, little Landau damping occurs, as evidenced by a narrow resonance having a peak reflection coefficient near unity. However, absorption becomes very pronounced with only a small increase in $v_T/v_{\text{phase}}$. We note that the resonance is heavily Landau damped at a ratio of $v_T/v_{\text{phase}}$ of less than 0.2. It is somewhat surprising to see so much damping at such a low value of $v_T/v_{\text{phase}}$, but this occurs because the width of the resonance is so small. The pertinence of the line width will be discussed later when we consider collisional damping. The resonant peak also occurs at smaller values of $(\omega_p/\omega)^2$ as $v_T/v_{\text{phase}}$ increases, as is evidenced in figure 5, where the abcissa is again exaggerated. The reflection and absorption coefficients are plotted as functions of $(\omega_p/\omega)^2$, for various values of the angle of incidence, in figure 6 with the following parameters fixed:

\begin{align*}
  k_0L &= 0.1 \\
  v/\omega &= 0 \\
  v_T/c &= 1.16 \times 10^{-3} \\
  (v_T/v_{\text{phase}} &= 0.182)
\end{align*}

As the angle of incidence increases from $\theta = 5^\circ$ to $\theta = 15^\circ$, the peak value of the absorption coefficient increases, in large part, because the longitudinal component of the electric field increases in proportion to $\sin \theta$. As $\theta$ further increases, the transverse electromagnetic waves become evanescent within the slab. Since the transverse waves and longitudinal waves are coupled, this leads to a washing out of the
Figure 4.- Reflection and absorption coefficients as a function of the plasma electron density for various electron thermal velocities.
Figure 5. - Resonant frequency as a function of electron density and electron thermal velocity.
Figure 6.- Reflection and absorption coefficients as a function of electron density for various angles of incidence.
resonance. Since electromagnetic waves in a plasma propagate as
\[ e^{\pm i k^p x} = e^{\pm i k_0 x} \sqrt{1 - \frac{\omega^2_0}{\omega^2} - \sin^2 \theta}, \]
the waves become evanescent when
\[ \theta = \cos^{-1} \frac{\omega_p}{\omega}, \]
which is about 19° for the case considered here. We further note in figure 6 that there is very little shift in the position of the resonance with increasing values of the angle of incidence.

The reflection and absorption coefficients plotted as a function of \((\omega_p/\omega)^2\) and \(v/\omega\) are shown in figure 7 for the model based on the Vlasov equation, with the following parameters fixed:
\[ \theta = 15^\circ \]
\[ k_0 L = 0.1 \]
\[ v_T/c = 10^{-3} \]
\[ (v_T/v_{\text{phase}} = 0.157) \]

For values of \(v/\omega\) less than about \(10^{-6}\), the collisions do not appreciably influence the reflection and absorption coefficients; however, as \(v/\omega\) increases a couple of orders of magnitude, the damping becomes pronounced. While it may, at first glance, seem surprising that such a large effect occurs for \(v/\omega\) as low as \(10^{-4}\), we can see from figure 7 that the line width is of the order of \(10^{-4}\). From our general knowledge of resonance phenomena, we expect damping to be appreciable whenever the collision frequency is of the order of the line width. This also accounts for the degree of Landau damping observed in figure 4.

Similar conclusions can be drawn by inspecting the fluid results, shown in figure 8.

Computer results for the impedance of the plasma capacitor are shown in figure 9, when the resistance and reactance (normalized to \(X_0 = 1/\omega C\))
Figure 7.- Reflection and absorption coefficients as a function of electron density for various values of collision frequency for a plasma model based on the linearized Vlasov equation with a B.G.K. collision term.
Figure 8. Reflection and absorption coefficient as a function of electron density with collision frequency as a parameter for a linearized fluid equation.
Figure 9.- The impedance of a plasma capacitor as a function of electron density for several values of the thermal velocity.
\( l = 5 \)
\[
\frac{v_T}{c} = 1.20 \times 10^{-3}
\]
\[
\frac{v_T}{v_{\text{phase}}} = 0.188
\]

(b) \( \frac{v_T}{c} = 1.20 \times 10^{-3} \)

Figure 9.- Continued.
\( l = 5 \)

\[
\frac{v_T}{c} = 1.30 \times 10^{-3}
\]

\[
\frac{v_T}{v_{\text{phase}}} = 0.203
\]

\[
\left( \frac{\omega_p}{\omega} \right)^2
\]

\( (c) \quad \frac{v_T}{c} = 1.30 \times 10^{-3} \)

Figure 9.- Concluded.
are plotted as a function of $(\omega_p/\omega)^2$. Figure 9(a) gives results for
$v_T/c = 10^{-3}$ ($v_T/v_{\text{phase}} = 0.157$), figure 9(b) gives than for
$v_T/c = 1.20 \times 10^{-3}$ ($v_T/v_{\text{phase}} = 0.188$), and figure 9(b) gives than for
$v_T/c = 1.30 \times 10^{-3}$ ($v_T/v_{\text{phase}} = 0.203$). In each case the normalized
slab thickness $k_0L$ is fixed at 0.1. Therefore, except for the angle
of incidence, which does not appear in the capacitor expression, all
pertinent capacitor parameters are the same as those for the $l = 5$ plane
wave resonance. We see from figure 9 that the resonance is very sharp
for low values of $v_T/c$, and the width at half-maximum noticeably broadens
as the thermal velocity increases, as does the reflection coefficient
of the plane wave. We further note that if figure 9 is compared with
figures 4 and 5, the peak of the resonance for the capacitor and the
obliquely incident p-polarized plane wave occurs at precisely the same
values of the plasma parameters.
CHAPTER III. DISCUSSION AND CONCLUDING REMARKS

Our prime objective in the work described herein was to analytically and computationally examine the details of the coupling phenomena that occur between electromagnetic waves and longitudinal plasma oscillations. We did this in order to relate this problem to the Tonks - Dattner resonances, which are known to occur when an electromagnetic field is applied to an inhomogeneous plasma in such a way as to couple to longitudinal plasma waves. Our model of the inhomogeneous plasma was of the simplest type: a thin uniform plasma slab. The specific boundary-value problem considered an electromagnetic plane wave obliquely incident upon the slab, and assumed specular reflection of plasma electrons from the faces of the slab. By solving for the reflection, transmission, and absorption coefficients, we could therefore examine the detailed behavior of the resonances; i.e., the shift in resonant frequency and changes in the width at half-maximum of the reflection coefficient as a function of the plasma parameters. The problem was approached in such a way as to delineate the differences, both from a physical and computational viewpoint, between a cold plasma model, a fluid model, and a model based on the Vlasov equation. The case where the electric vector was polarized parallel to the plane of incidence considered in detail. Only this polarization excites longitudinal plasma oscillation, for which the ratio of the thermal velocity of the particles to the phase velocity of the wave is not so small as to diminish the kinetic effects. Our analysis further showed that, from the practical viewpoint, the lower order longitudinal plasma resonances (i.e., \( l = 1, 3, \) and \( 5 \)) can be supported; these are
well separated only if the plasma slab is thin compared to the wavelength of the incident wave. The cold plasma model predicts, for such a slab, that only one resonance can be supported, and that occurs when the signal frequency equals the plasma frequency. The fluid model was found to support a series of resonances, which become narrower as the order increased, and which are similar in nature to the Tonks - Dattner resonances in the sense that the secondary resonances occur at successively lower values of the electron density than does the main resonance. Similar resonances were observed when a kinetic analysis was undertaken, using the Valslov equation, except that electron densities for resonance were shifted, and the effects of Landau damping became evident as the thermal velocity increased. This is a manifestation of the fact that the ratio $v_T/\nu_{\text{phase}}$ can no longer be considered negligible as the order increases. This damping with increasing order is qualitatively consistent with the experimental observations of the behavior of the Tonks-Dattner resonances. Such collisionless damping cannot be anticipated from the fluid equations. The fluid equations, however, have the advantage of presenting a simple physical picture of the standing wave processes that occur within the slab.

The pertinence of the parallel-plate, plasma filled capacitor problem was demonstrated by calculating the impedance of the capacitor as a function of slab thickness, plasma frequency, propagating frequency, and thermal velocity. It was found that when the slab is driven either by the capacitor or by a plane wave, resonances occur at the same values of the above parameters, and that Landau damping commences in both problems at the same values of the thermal velocity.
Most of the computational effort pertaining to the Vlasov equation consisted of evaluating the dispersion function, and using it to computationally determine the series given by (68) and (69) (see appendix C for details) this would, in turn, be used to determine reflection, transmission, and absorption coefficients. One of the more surprising aspects of the computations was that, for the cases considered, the infinite series given by (68) and (69) can accurately be represented by only two terms in the series. These terms consist of the first \( l = 0 \) term and that odd term \( l = 1, 3, 5, \ldots \) which corresponds most nearly to resonance for the parameters under consideration. This is to be contrasted with the half space solution (to which our results reduce as the slab dimension goes to infinity), which usually requires more involved computational techniques.

We then considered the influence of collisional damping of the resonances. For the particular parameters considered, it was observed that unless the collision frequency-to-signal frequency ratio is less than about \( 10^{-4} \), the third odd resonance is washed out. Since a collision frequency ratio of less that \( 10^{-4} \) for gaseous plasmas, is not easy to achieve in the laboratory, it is difficult to conclude that the uniform slab can support the experimentally observed Tonks-Dattner resonances.

We are tempted to conclude that the inhomogeneity of the plasma does more than merely control the spacing of the resonances\(^7\). Our results lead us to believe that the inhomogeneity may also tend to broaden the width of the higher-order resonances at half-maximum—which would make them less sensitive to collisional and Landau damping. This broadening could occur either because the inhomogeneity provides a gradual transition in the impedance between the plasma and the air interface, or because the
inhomogeneity induces a background field which causes resonant trapping of the electrons. An investigation of the latter problem (i.e., consideration of the background field) is a very formidable task, indeed. The description of even the simplest problems have developed into enormous and frustrating computer programming projects\(^{(25)} \, (26)\), and illuminating results are difficult to achieve unless many simplifying assumptions are made\(^{(27)}\). However, such results should present some new and rather interesting kinetic effects, and the problem is therefore worthwhile to pursue. If we can ignore the background field\(^{(28)}\), the problem becomes more tractable (but still computationally much more difficult than the uniform plasma). Such a solution may be valuable in order to determine whether or not a gradual reduction in electron density at the boundaries will broaden the resonances.

The techniques described here can easily be extended to other problems such as: (1) the study of the impedance characteristics of antennas under plasmas\(^{(9)}\), (2) the effects of nonspecular reflecting boundaries and (3) electromagnetic waves obliquely incident upon a plasma slab in which a static magnetic field is applied normal to the boundaries. All of these problems are of interest, and may be approached by extending the techniques of the current work.
APPENDIX A

FURTHER PROPERTIES OF THE DISPERSION INTEGRALS

If the collision frequency \( \nu \) is zero, we may define the various dispersion integrals as follows:

\[
\begin{bmatrix}
J_0 \\
J_1 \\
J_2 \\
J_3 \\
J_4 \\
J_5
\end{bmatrix} = \int\int_{-\infty}^{\infty} \frac{d\nu_x \, d\nu_z \, F_0}{\omega - k_0 v_z \sin \theta - \frac{\omega \pi v_x}{L}}
\]

where the integration over \( \nu_y \) has been performed, to give

\[
\int_{-\infty}^{\infty} d\nu_y \, F_0(v_x, \nu_y, \nu_z) = F_0(v_x, \nu_z),
\]

which is the two-dimensional Maxwellian distribution function. From the integral relationships:

\[
\int\int_{-\infty}^{\infty} \nu_x F_0 \, d\nu_x \, d\nu_z = \int\int_{-\infty}^{\infty} \nu_z F_0 \, d\nu_x \, d\nu_z = 0
\]

and

\[
\int\int_{-\infty}^{\infty} F_0 \, d\nu_x \, d\nu_z = 1
\]

It is possible to show, by multiplying the numerators and denominators of the integrands of (A-2) and (A-3) by \( \omega - k_0 v_z \sin \theta - \frac{\omega \pi v_x}{L} \) that

\[
J_1 = \frac{J_{2e}}{c} \sin \theta + \frac{\omega \pi}{\omega L} J_5
\]

\[
J_2 = \frac{J_{4e}}{c} \sin \theta + \frac{\omega \pi}{\omega L} J_2
\]
\[ 1 = \omega J_{0} - \left( \frac{L \pi}{L} J_{1l} + k_{0} \sin \theta J_{3l} \right) \]  
(A-6)

\[ J_{3l} = \frac{k_{0}L}{L \pi} \sin \theta J_{1l} \]  
(A-7)

And, as shown in the text, it is possible to use Fourier transforms in velocity space to re-express (A-1) in terms of single integrals over the transform variable \( \Lambda \). The results are:

\[
\begin{bmatrix}
J_{0} \\
J_{1l} \\
J_{3l} \\
J_{4l} \\
J_{5l}
\end{bmatrix} =
-\frac{i}{L \pi} \int d\Lambda \phi \Lambda^{2} \left[ e^{i \Lambda \phi / L \pi} - e^{-i \Lambda \phi / L \pi} \right] e^{i \Lambda \phi / L \pi}
\]

(A-8)

In the limit as \( v_{T} \rightarrow 0 \), the exponential involving \( v_{T}^{2} \) may be expanded as a Taylor series to give approximations for the real parts of the dispersion integrals as the thermal velocity of the plasma approaches zero. The results are:

\[ J_{0} \approx \frac{1}{\omega} \left[ 1 + \frac{v_{T}^{2}}{\omega^{2}} \left( k_{0}^{2} \sin^{2} \theta + \frac{L^{2} \pi^{2}}{L^{2}} \right) \right] \]  
(A-9)

\[ J_{1l} \approx \frac{v_{T}^{2}}{\omega^{2}} \frac{L \pi}{L} \left[ 1 + \frac{3 v_{T}^{2}}{\omega^{2}} \left( \frac{L^{2} \pi^{2}}{L^{2}} + k_{0}^{2} \sin^{2} \theta \right) \right] \]  
(A-10)
\[ J_{2l} \cong 2 k_0 \sin \theta \frac{l \pi}{L} \frac{V_T^4}{\omega^5} \] (A-11)

\[ J_{3l} \cong \frac{k_0 L}{\ell \pi} \sin \theta \overline{J}_{2l} \] (A-12)

\[ J_{4l} \cong \frac{V_T^2}{\omega} \left[ 1 + \frac{V_T^2}{\omega^2} \left( \frac{l^2 \pi^2}{L^2} + 3 k_0^2 \sin^2 \theta \right) \right] \] (A-13)

\[ J_{5l} \cong \frac{V_T^2}{\omega} \left[ 1 + \frac{V_T^2}{\omega^2} \left( \frac{3 l^2 \pi^2}{L^2} + k_0^2 \sin^2 \theta \right) \right] \] (A-14)

We now relate the above integrals to the tabulated plasma dispersion function, as given by Fried and Conte\(^{(23)}\). Fried and Conte define the dispersion function as follows:

\[ \Xi(\xi) = \frac{1}{\sqrt{\pi \xi}} \int_{-\infty}^{\infty} \frac{dx e^{-x^2}}{x - \xi} \] (A-15)

which may be recognized as the Hilbert transform of the Gaussian. If we perform a change of variables so that \( x = \frac{v \xi}{\sqrt{2}a} \), with

\[ a = v_T \sqrt{1 + \frac{k_0^2 L^2 \sin^2 \theta}{\frac{L^2 \pi^2}{2}}} \],

and define \( \zeta \) as

\[ \zeta = \frac{\omega L}{\sqrt{2}v_T} \frac{1}{\sqrt{1 + \frac{k_0^2 L^2 \sin^2 \theta}{\frac{L^2 \pi^2}{2}}}}, \]

the velocity transforms may be used to show that

\[ \Xi(\xi) = -\frac{\omega}{\xi} J_\ell (\zeta) \] (A-16)
Further manipulations may be performed to show that:

\[ J_{1l}(\xi) = \frac{2l \pi}{k_0 L} \left( \frac{v_T}{C} \right)^2 \frac{1}{k_0} \xi^2 \left[ 1 + \xi Z(\xi) \right] \]  
\[ (A-17) \]

\[ J_{2l}(\xi) = 2 \left( \frac{v_T}{C} \right)^4 \sin \theta \left( \frac{l \pi}{k_0 L} \right) \frac{c}{k_0} \xi^3 \left[ Z(\xi) - 2 \xi^2 Z(\xi) - 2 \xi \right] \]  
\[ (A-18) \]

\[ J_{3l}(\xi) = 2 \sin \theta \left( \frac{v_T}{C} \right)^2 \frac{1}{k_0} \xi^2 \left[ 1 + \xi Z(\xi) \right] \]  
\[ (A-19) \]

The identities (A-4) and (A-5), in connection with (A-17) - (A-19) may be used to express \( J_{4l}(\xi) \) and \( J_{5l}(\xi) \) in terms of the dispersion function. For \( l = 0 \), the dispersion integrals become:

\[ J_0 = \frac{J_{30}}{v_T^2} = - \frac{Z(\xi_0)}{k_0} \omega \]  
\[ (A-20) \]

\[ J_{10} = J_{20} = 0 \]  
\[ (A-21) \]

\[ J_{30} = - \frac{\left[ 1 + \xi_0 Z(\xi_0) \right]}{k_0 \sin \theta} \]  
\[ (A-22) \]

\[ J_{40} = - \frac{2 v_T \xi_0}{\sqrt{2} k_0 \sin \theta} \left[ \xi_0 Z(\xi) + 1 \right] \]  
\[ (A-23) \]

where

\[ \xi_0 = \frac{\omega}{k_0 \sqrt{2} v_T \sin \theta} \]  
\[ (A-24) \]
Finally, had we eliminated $H_{2y}$ in (43), the Fourier components of the current density could have been expressed in the following form:

\[ j_{l_x} = \frac{i \omega_p^2 \varepsilon_0}{v_T^2} \left[ J_{l_S} E_{l_x} + i J_{l_L} E_{l_z} \right] \]  \hspace{1cm} (A-25)

\[ j_{l_z} = \frac{\omega_p^2 \varepsilon_0}{v_T^2} \left[ i J_{l_L} E_{l_z} + J_{l_L} E_{l_x} \right] \]  \hspace{1cm} (A-26)

If we define the modes of the conductivity tensor ($\sigma$'s) through ohm's law:

\[ \begin{align*}
  j_{l_x} &= \sigma_{x xl} E_{l_x} + \sigma_{x zl} E_{l_z} \\
  j_{l_z} &= \sigma_{z xl} E_{l_x} + \sigma_{z zl} E_{l_z}
\end{align*} \]  \hspace{1cm} (A-27)

Our expressions agree with those of Hinton, who solved this problem by integrating over particle trajectories.

For a finite collision frequency ratio $\nu/\omega$, the pertinent dispersion integrals are given by (56), (57), (59), and (60), which are also expressible in terms of the dispersion function, giving

\[ J_{l_S} = -(1 + i \nu/\omega) \frac{2\pi}{k_0 \omega} \frac{v_T^2}{c} \frac{1}{k_0^2} \varepsilon^2 \left[ l + \xi Z(\xi) \right] \]  \hspace{1cm} (A-28)

\[ J_{l_L} = 2 \left( \frac{v_T}{c} \right)^4 \sin \theta \left( \frac{v_T}{k_0 \omega} \right) \frac{c}{k_0} \frac{1}{k_0} \varepsilon^3 \left[ Z(\xi) - 2\xi Z(\xi) - 2\xi \right] \]  \hspace{1cm} (A-29)
\[ J_{3l} = -2(1 + i \nu / \omega) \sin \theta \left( \frac{v_T}{c} \right)^2 \frac{l}{k_0^2} \xi^2 \left[ 1 + \xi Z(\xi) \right] \]  \hfill (A-30)

\[ J_{4l} = \frac{c}{\sin \theta} \left[ J_{3l} - \frac{\ell \pi}{k_0 L} \frac{1}{c} J_{2l} \right] \]  \hfill (A-31)

where, for a finite collision frequency

\[ \xi = \xi^R (1 + i \nu / \omega) \]  \hfill (A-32)

where

\[ \xi^R = \frac{1}{\sqrt{2} \frac{v_T}{c} \sqrt{\frac{l^2 \pi^2}{k_0^2 L^2} + \sin^2 \theta}} \]  \hfill (A-33)
APPENDIX B

INTERACTION OF A PLANE WAVE WITH A PLASMA HALF SPACE

For the case where the slab thickness $L$ approaches infinity, we can solve the problem in the original way that Landau and others have done. We write the Vlasov equation in the form:

$$-i\omega f_1 + v_x \frac{\partial f_1}{\partial x} + i k_0 v_Z \sin \Theta - \frac{e}{m} E_Z \frac{\partial f_0}{\partial v_z} - \frac{e}{m} E_x \frac{\partial f_0}{\partial v_x} = 0 \quad (B-1)$$

We again Fourier analyze the Vlasov equation in configuration space, except that the half space requires a superposition of continuous modes rather than discrete modes. In other words, $f_1$ is expanded as a Fourier integral instead of a Fourier series. We have:

$$f_1(\mathbf{v}, x) = \int_{-\infty}^{\infty} f_1(\mathbf{v}, k_x) e^{ik_x x} \, dk_x \quad (B-2)$$

A straightforward Fourier analysis of $(B-1)$, in connection with the convolution theorem gives:

$$f_1 = -\frac{e}{mv_x} \int_x^\infty dx' e^{ib(x-x')} \left[ \frac{\partial f_0}{\partial v_z} E_z(x') + \frac{\partial f_0}{\partial v_x} E_x(x') \right] : \quad v_x < 0 \quad (B-3)$$

$$f_1 = \frac{e}{mv_x} \int_{-\infty}^x dx' e^{ib(x-x')} \left[ \frac{\partial f_0}{\partial v_z} E_z(x') + \frac{\partial f_0}{\partial v_x} E_x(x') \right] : \quad v_x > 0$$
with

\[ b = \frac{\omega - k_0 v_z \sin \theta}{\sqrt{\chi}} \]  \hspace{1cm} (B-4)

If we express the field components of (B-3) in terms of their Fourier transforms,

\[
\tilde{f}_i = \frac{n \varepsilon_0 c}{\text{im} \nu_i^2} \frac{F_0}{\omega - k_0 v_z \sin \theta - k_\chi \nu_x} \left[ v_x \tilde{E}_z + v_x \tilde{E}_x \right]
\]  \hspace{1cm} (B-5)

for all \( \nu_x \) where a 'bar' denotes the Fourier transform. From our definition of the current components, we have:

\[
\tilde{j}_x = \frac{i \varepsilon_0 \omega^2}{\nu_x^2} \left[ \tilde{E}_z J_2(k_\chi) + \tilde{E}_x J_5(k_\chi) \right]
\]  \hspace{1cm} (B-6)

\[
\tilde{j}_z = \frac{i \varepsilon_0 \omega^2}{\nu_x^2} \left[ \tilde{E}_z J_4(k_\chi) + \tilde{E}_x J_1(k_\chi) \right]
\]  \hspace{1cm} (B-7)

where the \( J \)'s in (B-6) and (B-7) are identical to the dispersion integrals given in appendix A, with \( \lambda \pi/L \) replaced by the continuous wave number, \( k_\chi \). The currents explicitly appear in the wave equation as follows:

\[
\frac{d^2 H_y}{dx^2} + (k_0^2 - k_\chi^2 \sin^2 \theta) H_y = -i k_\chi j_x + \frac{d j_z}{dx}
\]  \hspace{1cm} (B-8)
The form of (B-3) suggests the symmetry properties $H_y(x) = -H_y(-x)$, $j_x(x) = -j_x(-x)$, and $j_z(x) = j_z(-x)$. Using these properties to Fourier transform (B-3) gives:

$$
(k_0^2 - k_x^2 \sin^2 \theta - k_x^2) \overline{H_y} - ik_x \overline{H_y(0)} = 
\frac{\varepsilon_0 \omega^2}{V_f} E_x \left[-k_x \overline{J_2(k_x)} + k_0 \sin \theta \overline{J_5(k_x)} \right] 
+ \frac{\varepsilon_0 \omega^2}{V_f} E_z \left[-k_x \overline{J_4(k_x)} + k_0 \sin \theta \overline{J_7(k_x)} \right]
$$

(B-9)

Where a bar over a field component denotes a Fourier transform. The Fourier transforms of the Maxwell curl equations are:

$$
i k_x \overline{E_z} = i k_0 \sin \theta \overline{E_x} - i \omega \mu_0 \overline{H_y}
$$

(B-10)

$$
\overline{E_x} = \frac{k_0 \sin \theta}{\omega \varepsilon_0} \overline{H_y} + \frac{1}{i \omega \varepsilon_0} \left\{ \frac{i \omega \varepsilon_0}{V_f} \left[ \overline{E_z J_2(k_x)} + \overline{E_x J_5(k_x)} \right] \right\}
$$

(B-11)

Equations (B-9), (B-10), and (B-11) are all that are needed to solve for all the field components. After several ponderous algebraic manipulations, one can obtain the following results:

$$
\overline{H_y} = \frac{ik_x \overline{H_y(0)}}{\pi \left\{ k_0^2 - k_0^2 \sin^2 \theta - k_x^2 \frac{\omega^2 k_0^2}{\omega V_f^2} \left[ \overline{J_4(k_x)} - \frac{k_0 \sin \theta \overline{J_7(k_x)}}{k_x} \right] \right\}}
$$

(B-12)
Now, the field expressions for the plasma slab (equations (64)-(66)) are of the following form:

\[
H_y = \sum_{l=1}^{\infty} H_{ly} \sin \frac{l \pi x}{L} \quad \text{(B-15)}
\]

\[
E_x = \sum_{l=1}^{\infty} E_{lx} \sin \frac{l \pi x}{L} \quad \text{(B-16)}
\]

\[
E_z = \sum_{l=0}^{\infty} \frac{E_{lz}}{1 + b_0} \cos \frac{l \pi x}{L} \quad \text{(B-17)}
\]

But, it also follows from (64)-(66) that \( H_{ly} = -H_{-ly}, \quad E_{lx} = -E_{-lx} \) and \( E_{lz} = E_{-lz} \), so that the summation index \( l \) can be extended from \( -\infty \) to \( +\infty \) giving \( H_y = \frac{1}{2\imath} \sum_{-\infty}^{\infty} H_{ly} e^{1\imath\pi x/L}, \) and similar expressions for \( E_x \).
and $E_z$. In the limit as $L \to \infty$, the summations convert to integrals such that $\frac{2\pi}{L} \to k_x$ and $dk_x \to \pi/L$. In this limit:

$$H_y = \frac{1}{2i} \sum_{k_x} H_y e^{\frac{L \pi x}{L}} \int_{-\infty}^{\infty} H_y(k_x) e^{i k_x x} dk_x$$  \hspace{1cm} (B-18)

In which case, the sum of the terms defining $H_y$ in (64) becomes:

$$H_y \to \sum_{k_x} \frac{dk_x e^{i k_x x} k_x H_y(0)}{\pi i \left\{ k_x^2 - k_0^2 \sin^2 \theta - k_e^2 - \frac{\omega^2 \omega_p^2}{v e^2 k_e^2} \right\}}$$  \hspace{1cm} (B-19)

Where $H_y(L) \to 0$ as $L \to \infty$. Equation (B-19) is the Fourier transform of (B-12). Since (B-13) and (B-14) follow in a similar manner, we have proven that our results for the slab reduce to the half space results as $L \to \infty$, which was to be shown.
APPENDIX C

SOME COMMENTS ON THE COMPUTATIONAL PROCEDURE

The programs for all numerical computations were written in Fortran IV for use on the CDC-6600 computer facility at the Langley Research Center.

The programs applicable to the cold plasma and hydrodynamic models involve only the elementary transcendental functions of complex arguments, and therefore do not require any detailed explanation. The programs relating to the Vlasov equation should, however, be discussed in a little more detail. The reflection, transmission, and absorption coefficients were determined by evaluating the series defining $G_1$ and $G_2$ (equations (68) and (69)), which are given again below, for reference:

\[
G_1 = \frac{1}{(k_0 L)^2} \sum_{l=0}^{\infty} \frac{2}{1+\delta_0^2} \frac{\left[ 1 - \sin^2 \theta - \frac{\omega_p^2}{\nu_T^2 \omega_x^2} J_{5l} \right]}{\left[ 1 - \frac{\omega_p^2}{\nu_T^2} \frac{\nu_T}{\omega_x} J_{5l} \right]} \left[ 1 - \sin^2 \theta - \left( \frac{l \pi}{k_0 L} \right)^2 \frac{\omega_p^2}{\nu_T^2 \omega_x^2} \left( J_{4l} - \frac{k_0}{2 \pi} \sin \theta J_{2l} \right) \right]
\]

\[
G_2 = \frac{1}{(k_0 L)^2} \sum_{l=0}^{\infty} \frac{2}{1+\delta_0^2} \frac{\left[ 1 - \sin^2 \theta - \frac{\omega_p^2}{\nu_T^2 \omega_x^2} J_{5l} \right] (1)^l}{\left[ 1 - \frac{\omega_p^2}{\nu_T^2} \frac{\nu_T}{\omega_x} J_{5l} \right]} \left[ 1 - \sin^2 \theta - \left( \frac{l \pi}{k_0 L} \right)^2 \frac{\omega_p^2}{\nu_T^2 \omega_x^2} \left( J_{4l} - \frac{k_0}{2 \pi} \sin \theta J_{2l} \right) \right]
\]
Equations (68) and (69) are then used in a trivial manner in connection with (15), (16), (17), and (18) to solve for the reflection, transmission, and absorption coefficients. The dispersion integrals (the $J$'s) in (68) and (69) were then expressed in terms of the dispersion function $Z(\xi)$, which for $\xi < 4$ was computed from the differential equation:

$$Z'(\xi) = -2 \left[ 1 + \xi Z(\xi) \right] \tag{C-1}$$

Under the initial condition that $Z(0) = i\sqrt{\pi}$. For $\xi > 4$, the asymptotic series was used. Calculations were performed only for real

$$\xi_R = \frac{1}{\sqrt{2} \frac{V_T}{c} \sqrt{\frac{e^2 \mu^2}{k^2 L^2} + \sin^2 \Theta}}$$

For small $v/\omega$, the dispersion function was evaluated using a Taylor series expansion:

$$Z(\xi) \approx Z(\xi) \bigg|_{v/\omega = 0} + \frac{d Z(\xi)}{d (v/\omega)} \bigg|_{v/\omega = 0} \cdot \frac{v}{\omega} \tag{C-2}$$

Where, from (A-32) $\xi = \xi^R (1 + iv/\omega)$. Equation (C-2) therefore reduces to

$$Z(\xi) \approx Z(\xi^R) - 2i \left[ 1 + \xi^R Z(\xi^R) \right] \xi^R \frac{v}{\omega} \tag{C-3}$$

for $v/\omega << 1$.

For plasma slabs thin compared to a wavelength, the numerical results indicated that $G_1$ and $G_2$ converged so rapidly that only the
\( \ell = 0 \) term and the most nearly resonant term contribute to the series for \( G_1 \) and \( G_2 \). The resonant term is that term in odd \( \ell \) for which 
\[ 1 - \frac{a_p^2}{\nu_1^2} L/\ell \pi J_{1\ell} \]
in the denominator of (68) and (69) is a minimum.

The dominant \( \ell = 0 \) term is simply

\[
G_1(\ell = 0) = G_2(\ell = 0) = \frac{1}{k_o \nu_1^2 L^2 (1 - \frac{\omega_p^2}{\omega^2})} \quad (C-4)
\]
V. REFERENCES


Calvin T. Swift was born in Quantico, Virginia, February 6, 1937. He graduated from E. C. Glass High School in Lynchburg, Virginia, June 1955. After receiving the B.S. degree in physics from the Massachusetts Institute of Technology in June 1959, he accepted employment as a research engineer with North American Rockwell in Downey, California. In 1962, he came to the NASA, Langley Research Center, where his research specialty is the study of electromagnetic wave interaction with plasmas. The purpose of this research has been to diagnose reentry plasmas and the solar corona. Mr. Swift has published several NASA reports and journal articles in this field, and was recently awarded certificates of achievement for two of the best twelve articles published in an IEEE Transactions during the year, 1968.

In 1962, Mr. Swift began graduate work, and was awarded the degree of Master of Science in physics from the Virginia Polytechnic Institute in June, 1965.

He is married to the former Joanne Taylor of Bucksport, Maine, and they have two children, Pamela and Janet.