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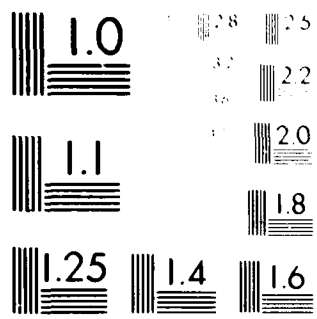

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In the earth's ionosphere the presence of nonthermal electrons causes enhancements of the electron plasma line resonance observed by incoherent backscatter radar. This paper extends previous calculations of the plasma intensity by including resonance broadening due to electron neutral collisions. A BGK collision term is included in the derivation of the electron spectral density function obtained from Fluctuation Dissipation theory. Assuming the non-thermal electron distribution is a small perturbation on a Maxwellian distribution of background electrons,

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20. **ABSTRACT (Continued)**

the expression describing the signal intensity for a plasma line reduces to the form obtained previously when only electron-ion collisions were included. Results of this extended model are compared to recent measurements made with the Chatanika radar.

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THE EFFECTS OF ELECTRON-NEUTRAL COLLISIONS ON THE INTENSITY OF PLASMA LINES

I. INTRODUCTION

The ionosphere is a region of weakly ionized plasma essentially in thermal equilibrium. When incoherent radar is used to probe the ionosphere, VHF radiation is scattered off density inhomogeneities at various altitudes and the spectral character of the scattered wave is examined. If the inhomogeneities have wavelengths longer than the Debye length, the scattered spectra display sharp resonances near the plasma frequency.[1] The intensity of these resonances, called plasma lines, depends on the electron density distribution and is enhanced by the presence of nonthermal electrons.[2] These nonthermal electrons may be photo electrons excited by solar electromagnetic radiation or they may be secondary electrons excited by energetic particles such as electrons, protons or heavier ions deposited in the ionosphere. The shape and intensity of the plasma lines are further altered by Landau damping and collisional broadening.

The intensity of a plasma line has been calculated by Yngvesson and Perkins[3], who included the effect of Coulomb collisions between electrons and ions as well as the effect of Landau damping. Their result may be expressed in terms of an effective temperature for the plasma line.

$$T_p = \frac{f_m + f_p + \chi_{ei}}{f_m + L_p + \chi_{ei}} \tag{1}$$

The terms, f_p and L_p, represent excitation and damping of plasma waves due to nonthermal electrons. The term, f_m, describes Landau damping while χ_{ei} describes effects of electron-ion collisions. The term f_m incorporates an assumption that the distribution of background thermal electrons is basically Maxwellian. T_p of Eq. (1) may be related directly to the power of the returned radar signal.

The measurement of plasma lines is of current interest due to both the availability of accurate observational data[4,5] and the development of sophisticated theoretical and computational techniques[6-10]

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for calculating the ionospheric electron distribution function. For example, comparisons of model predictions with recent Chatanika data[5] have shown good agreement at altitudes high enough to allow the neglect of electron-neutral collisions.

The purpose of the present work is to extend the theoretical predictions of Yngvesson and Perkins[3] to include self-consistently the effects of electron collisions with neutral species which must be included at lower altitudes. The approach we have taken is to apply Fluctuation Dissapation theory [10], representing collisions by a BGK collision term in the Boltzmann equation. Since Fokker-Planck theory is only appropriate for Coulomb collisions, we cannot use the same approach as Perkins and Salpeter.

In the next section we discuss the representation of collisions in the Boltzmann equation. The following sections outline the calculation of the spectral density function, which is integrated across the plasma resonance to yield the intensity of a plasma line. The result is shown to compare well with effective temperatures derived from Chatanika data.[7]

II. THEORETICAL REPRESENTATION OF COLLISIONS

Perkins and Salpeter applied a Fokker-Planck collison term in order to be able to model small-angle deflections for electron-ion collisions. They derived a collision frequency for these long-range Coulomb collisions, which had the same form as a typical momentum transfer collision frequency. This frequency contains kinetic information about local collisional processes, but also contains an implicit assumption that the distribution function for electrons is basically Maxwellian. Since electron-ion collisions in the ionosphere principally involve particles of average energy, these collisions cannot produce an enhancement of the electron density fluctuations, but only damping.

In order to understand what kind of approximations can be made in representing different kinds of collisions, it is important to know the length scales and timescales which dominate physical processes in various regions of the ionosphere. The collision time, $\tau = 1/\nu$, represents the amount of time needed for an electron to experience a significant deflection. That is, such a deflection results from many small angle collisons in the case of electron-ion collisions, whereas a single electron-neutral collision would produce on average the same deflection. Both electron-neutral and electron-ion collision times are much larger than the reciprocal electron plasma frequency, ω_p^{-1}.

$$\nu_{ei} << \nu_{en} << \omega_p \qquad z < 150 \text{ km}$$
$$\nu_{en} << \nu_{ei} << \omega_p \qquad z > 200 \text{ km}.$$

The radar only detects waves of a particular wavelength, typically 23cm for recent Chatanika experiments. This is long compared to typical electron Debye lengths which are of the order of a few centimeters. Typical length scales compare with the collisional mean free paths $l = \left(\frac{\gamma kT}{m}\right)^{1/2} \frac{1}{\nu}$, as

$$\lambda_{De} < l_{en} < l_{ei} < \lambda, \qquad 100 \text{ km} < z < 110 \text{ km}$$
$$\lambda_{De} < \lambda < l_{en} < l_{ei}, \qquad z > 120 \text{ km}.$$

This relationship suggests that one could even use fluid theory to describe electron-neutral collisions in the E-region since the E-region plasma is not sensitive to details of individual collisions. Kinetic theory is required for the F-region where collisional mean free paths are greater than the radar wavelength. Hence, we have chosen a BGK representation of the collision term in the Boltzmann equation. The BGK collision term is particle conservative and is designed for electron-neutral or ion-neutral collisions.[11] It requires that a collision frequency be provided which contains information about individual collisions and represents the frequency at which significant deflections occur in the direction of particle motion.

Note that the Perkins and Salpeter electron-ion collision frequency provides an adequate representation which can sensibly be applied in this model. Recall, their value does not indicate the frequency of small-deflection collisions but rather an averaged value based on a collection of small angle collisions sufficient to produce a significant deflection. Since we will be assuming that the electron distribution function is basically Maxwellian (as was assumed in the derivation of Eq. (1))[2], we can represent the effect of electron-ion collisions by a BGK term. This treatment is not typical but may be used for calculating the intensity of a plasma line for which collisions are a minor effect. We show later that our results reduce to the Fokker-Planck results when electron-neutral collisons are neglegible.

The BGK representation of the Boltzmann collision term has the form

$$\left(\frac{\partial f}{\partial t}\right)_c = -\nu\left[f_1(\mathbf{x},\mathbf{v},t) - \frac{N_1(\mathbf{x},t)}{N_0} f_0(\mathbf{v})\right], \tag{2}$$

where f_0 is the ambient election distribution in velocity space, f_1 is a perturbation on this distribution, and N_0 and N_1 are integrals of f_0 and f_1 respectively, taken over velocity space. A Boltzmann collision term must be provided to represent each kind of collision of electrons with other species. The appropriate collision frequencies have been derived by Banks[13] and Itikawa[14] for electron-neutral collisions. The collision frequencies do not depend on the electron density and are so slowly varying in space and time that they can be assumed to be constant except for a weak altitude dependence. [The quantity $\nu(\partial\nu/\partial z)^{-1}$ has a value of 10 km at an altitude of 105 km for the dominant electron-neutral collisions. This value is much larger than the wavelengths of interest to us.]

A total collision term described by Eq. (2) can hence be written using a total collision frequency

$$\nu \equiv \sum_j \nu_{en}(n_j) + \nu_{ei} \tag{3}$$

where the summation is taken over all relevant electron-neutral collisional interactions.

III. CALCULATION OF THE SPECTRAL DENSITY, $n(k,\omega)$

Given a suitable total collision frequency ν_l, the first order Boltzmann equation for species l is

$$\frac{\partial F_{1l}(\mathbf{r},\mathbf{v},t)}{\partial t} + \mathbf{v}\cdot\frac{\partial F_{1l}(\mathbf{r},\mathbf{v},t)}{\partial r} + \frac{q}{m}\mathbf{E}_1\cdot\frac{\partial F_{0l}(\mathbf{v})}{\partial \mathbf{v}} = -\nu[F_{1l}(\mathbf{r},\mathbf{v},t) - n_l F_{0l}(\mathbf{v})] \tag{4}$$

with

$$n_{1l}(r,t) = \int d\mathbf{v} F_{1l}(\mathbf{r},\mathbf{v},t). \tag{5}$$

We assume $F_{0l}(\mathbf{v})$ is a zeroth order distribution function normalized so that the integral of F_{0l} taken over velocity space is unity. F_{1l} $(\mathbf{r},\mathbf{v},t)$ is a first-order spatial perturbation normalized by the same constant. Similarly, we define a perturbation electric field $\mathbf{E}_1 \equiv -\nabla\phi$, where ϕ is the electrostatic potential, so that

$$\nabla\cdot\mathbf{E}_1 = \sum_l 4\pi q_l \int d\mathbf{v} F_{1l} \tag{6}$$

The summation is taken over species l, representing electrons and ions, of charge q. We then have

$$\nabla^2\phi(\mathbf{r},t) = -4\pi(Zen_{1i} - en_{1e}) \tag{7}$$

In order to determine the spectral densities, $n_{1l}(\mathbf{k},\omega)$, we calculate the Fourier-Laplace transforms of (7) and (4):

$$-k^2\phi = -4\pi[Zen_{1i}(\mathbf{k},\omega) - en_{1e}(\mathbf{k},\omega)] \tag{8}$$

$$-F_{1l}(\mathbf{k},\mathbf{v},0) + [i\omega + \nu_l - i(\mathbf{k}\cdot\mathbf{v})]F_{1l}(\mathbf{k},\mathbf{v},\omega) =$$

$$-\frac{i4\pi q}{mk^2}[Zen_{1i}(\mathbf{k},\omega) - en_{1e}(\mathbf{k},\omega)]\,\mathbf{k}\cdot\frac{\partial F_{0l}}{\partial \mathbf{v}} + \nu_l n_{1l}(\mathbf{k},\omega)F_{0l} - n_{1l}(\mathbf{k},0)\,F_{0l} \tag{9}$$

The first term in Eq. (9) appears as a constant of integration because we are performing a Laplace transform rather than a Fourier transform in time. When we divide Eq. (9) by $(\omega - \mathbf{k}\cdot\mathbf{v} - i\nu_l)$ and integrate over $\mathbf{v}$, the first term is expressed in terms of

$$n_{1e}(\mathbf{k},0) = \sum_{j=1}^{N} e^{i\mathbf{k}\cdot\mathbf{r}_j(0)}\,\delta(\mathbf{v} - \mathbf{v}(0)), \tag{10}$$

and (9) becomes[12]:

$$-\frac{n_{1e}(\mathbf{k},0)}{\omega - \mathbf{k}\cdot\mathbf{v} - i\nu_e} + in_{1e}(\mathbf{k},\omega) = +\frac{i4\pi e}{mk^2}[Zen_{1i}(\mathbf{k},\omega)] - en_{1e}(\mathbf{k},\omega)]\int\frac{d\mathbf{v}\,\mathbf{k}\cdot\dfrac{\partial F_{0e}}{\partial \mathbf{v}}}{\omega - \mathbf{k}\cdot\mathbf{v} - i\nu}$$

$$+\nu_e n_{1e}(\mathbf{k},\omega)\int d\mathbf{v}\frac{F_{0e}}{\omega - \mathbf{k}\cdot\mathbf{v} - i\nu_e}$$

$$-n_{1e}(k,0)\int d\mathbf{v}\frac{F_{0e}}{\omega - \mathbf{k}\cdot\mathbf{v} - i\nu_e} \tag{11}$$

We define:

$$D_l(\mathbf{k},\omega) \equiv i\nu_l\int_{-\infty}^{\infty}\frac{d\mathbf{v}\,F_{0l}}{\omega - \mathbf{k}\cdot\mathbf{v} - i\nu_l},\ \text{and} \tag{12}$$

$$C_l(\mathbf{k},\omega) = \frac{1}{(1 + D_l)}\int_{-\infty}^{+\infty} d\mathbf{v}\,\frac{4\pi Z_l e^2 n_0}{m_l k^2}\,\frac{\mathbf{k}\cdot f_{0l}/\partial\mathbf{v}}{\omega - \mathbf{k}\cdot\mathbf{v} - i\nu_l}. \tag{13}$$

Then Eq. (11) for electrons and an analogous equation for ions may be combined to eliminate n_{1i}, yielding:

$$n_{1e}(\mathbf{k},\omega)$$
$$= -i\left[\left[1-\frac{C_e}{\epsilon(\mathbf{k},\omega)}\right]\times\left[\frac{1}{1+D_e}\sum_{j=1}^{N}e^{i\mathbf{k}\cdot\mathbf{r}_j(0)}\left\{\frac{1}{\omega-k\cdot\mathbf{v}_j(0)-i\nu_e}-\frac{D_e}{i\nu_e}\right\}\right]\right.$$
$$\left.+\frac{ZC_e}{\epsilon(\mathbf{k},\omega)}\left[\frac{1}{(1+D_i)}\sum_{j=1}^{N}e^{i\mathbf{k}\cdot\mathbf{r}_j(0)}\left\{\frac{1}{\omega-k\cdot\mathbf{v}_j(0)-i\nu_i}-\frac{D_i}{i\nu_i}\right\}\right]\right]. \tag{14}$$

with

$$\epsilon(\mathbf{k},\omega) = 1 + C_e(\mathbf{k},\omega) + C_i(\mathbf{k},\omega) \tag{15}$$

This value will be used in the next section to calculate the effective temperature associated with a plasma line.

IV. CALCULATION OF THE EFFECTIVE TEMPERATURE, T_p

As in the treatment by Yngvesson and Perkins[3], we define a spectral density function $S(\mathbf{k},\omega)$ which describes the spectral amplitude as a function of wavenumber and frequency,

$$S(\mathbf{k},\omega) = \lim_{\substack{V\to\infty \\ \nu_i<<\omega}} \frac{2\nu_l}{V}\frac{<|n_{1e}(\mathbf{k},\omega)|^2>}{n_{e0}}. \tag{16}$$

The intensity associated with a resonant peak is proportional to the effective temperature, T_p of Eq. (1), and is simply $S(\mathbf{k},\omega)$ integrated over the frequency range about the plasma resonance.

When the value of $n_{1e}(\mathbf{k},\omega)$ is substituted from Eq. (13), the function $S(\mathbf{k},\omega)$ becomes

$$S(\mathbf{k},\omega) = 2|(1 + C_i)/\epsilon|^2 B_e + 2Z|C_e/\epsilon|^2 B_i, \tag{17}$$

where

$$B_l = \frac{\nu_l}{|1 + D_l|^2}\left\{\int_{-\infty}^{+\infty}\frac{d\mathbf{v}f_{0l}(\mathbf{v})}{(\omega - \mathbf{k}\cdot\mathbf{v})^2 + \nu_l^2} - \frac{|D_l|^2}{\nu_l^2}\right\}. \tag{18}$$

The distribution function $f_{0l}(\mathbf{v})$ describes species l in the thermal plasma and it is best represented as a Maxwellian[3].

Choosing the Maxwellian distributions:

$$f_{e0}(\mathbf{v}) = \exp(-v^2/a^2)/(\pi a^2)^{3/2}, \quad a = (2\kappa T_e/m_e)^{1/2}, \tag{19}$$

$$f_{i0}(\mathbf{v}) = \exp(-v^2/b^2)/(\pi b^2)^{3/2}, \quad b = (2\kappa T_i/m_i)^{1/2}, \tag{20}$$

the previous coefficients become,

$$D_e(\mathbf{k},\omega) = \frac{i\nu_e}{ka}\left[2\exp(-y_e^2)\int_0^{y_e}\exp(p^2)\,dp + i\pi^{1/2}\exp(-y_e^2)\right] \tag{21}$$

$$D_i(\mathbf{k},\omega) = \frac{i\nu_i}{kb}\left[2\exp(-y_i^2)\int_0^{y_i}\exp(p^2)\,dp + i\pi^{1/2}\exp(-y_i^2)\right], \tag{22}$$

where

$$y_e = (\omega - i\nu_e)/ka, \quad y_i = (\omega - i\nu_i)/kb, \tag{23}$$

$$B_e(\mathbf{k},\omega) = \frac{1}{\pi^{1/2}ka|1+D_e|^2}\,\mathrm{Im}\left\{\int_{-\infty}^{+\infty} d\left(\frac{v}{a}\right)\frac{\exp[-(v/a)^2]}{(y_e - v/a)}\right\} - \frac{|D_e|^2}{\nu_e|1+D_e|^2}$$

$$= \frac{1}{ka|1+D_e|^2}\,\mathrm{Im}\left\{2\exp(-y_e^2)\int_0^{y_e}\exp(p^2)\,dp + i\pi^{1/2}\exp(-y_e^2)\right\}$$

$$- \frac{|D_e|^2}{\nu_e|1+D_e|^2} \tag{24}$$

$$B_i(\mathbf{k},\omega) = \frac{1}{\pi^{1/2}kb|1+D_i|^2}\,\mathrm{Im}\left\{\int_{-\infty}^{+\infty} d\left(\frac{v}{b}\right)\frac{\exp[-(v/b)^2]}{(y_i - v/b)}\right\} - \frac{|D_i|^2}{\nu_i|1+D_i|^2} \tag{25}$$

$$C_e(\mathbf{k},\omega) = \frac{\alpha^2}{(1+D_e)}\left[1 - 2y_e\exp(-y_e^2)\int_0^{y_e}\exp(p^2)\,dp - i\pi^{1/2}y_e\exp(-y_e^2)\right] \tag{26}$$

$$C_i(\mathbf{k},\omega) = \frac{ZT_e}{T_i}\frac{\alpha^2}{(1+D_i)}\left[1 - 2y_i\exp(-y_i^2)\int_0^{y_i}\exp(p^2)\,dp - i\pi^{1/2}y_i\exp(-y_i^2)\right], \tag{27}$$

where $\alpha = (k\lambda_D)^{-1}$ and λ_D is the Debye length. Incorporating these values, Eq. (17) reduces to

$$S_e(\mathbf{k},\omega,\nu_e) = \frac{2}{ka}\,\frac{\mathrm{Im}\left[i\pi^{1/2}\exp(-y_e^2) + 2\exp(-y_e^2)\int_0^{y_e}\exp(p^2)\,dp\right] - \frac{ka}{\nu_e}|D_e|^2}{|1 + D_e + \alpha^2(R\mathrm{w}(y_e) - iI\mathrm{w}(y_e))|^2}. \tag{28}$$

Since $\omega \cong \omega_{pe} >> ka$, that is $\alpha^2 > 1$, we define w by

$$R\,\mathrm{w}(x_e) \cong -\frac{1}{2x_e^2} - \frac{i\nu_e}{x_e^2\omega} - \frac{3}{4x_e^4} \tag{29}$$

and

$$I\,\mathrm{w}(x_e) = \pi^{1/2}\, x_e \exp(-x_e^2) \tag{30}$$

with

$$x_e = \omega/ka >> 1$$

Hence,

$$D_e \cong + \frac{i\nu_e}{\omega}\left[1 + \frac{1}{2x_e^2}\right]. \tag{31}$$

Defining $\psi_e \equiv \nu_e/ka << 1$. Eq. (28) reduces to[12]

$$S_e(\mathbf{k},\omega,\nu_e) \cong \frac{x_{er}^2\left[\pi^{1/2}\exp(-x_{er}^2) + \dfrac{\psi_e}{2x_{er}^4}\right]}{2ka\left\{(x_e - x_{er})^2 + \dfrac{x_{er}^2\alpha^4}{4}\left[\dfrac{\psi_e}{2x_{er}^3} + \pi^{1/2}x_{er}\exp(-x_{er}^2)\right]^2\right\}}$$

where subscript r indicates a calculation at the resonance frequency. In terms of an arbitrary function.

$$f_{e0}(\omega_r/k) = (1/a\sqrt{\pi})\exp[-(\omega_r/ka)^2],$$

we obtain

$$S_e(\mathbf{k},\omega,\nu_e) \cong \frac{\pi}{2k}\,\frac{\omega_r^2\left\{f_{e0}\left(\dfrac{\omega}{k}\right) + \dfrac{k^3\kappa T_e\nu}{\pi m_e\omega_r^4}\right\}}{\left\{(\omega-\omega_r)^2 + \dfrac{\omega_{pe}^4}{2\omega_r^4}\left[\nu_e - \dfrac{\pi\omega_r^3}{k^2}\left[\dfrac{\partial}{\partial \mathrm{v}}f_{e0}(\mathrm{v})\right]_{\mathrm{v}=\omega_r/k}\right]^2\right\}}. \tag{32}$$

The intensity of a single plasma line is defined as

$$I_p \equiv \lim_{\delta\to 0}\int_{-\delta}^{\delta} S_e(\mathbf{k},\omega)\,d\theta \tag{33}$$

with

$$\theta = \omega - \omega_r \text{ and } \omega \cong \omega_r.$$

Recall

$$\lim_{\delta\to 0}\int_{-\delta}^{\delta}\frac{d\theta}{\theta^2 + B^2} \cong \frac{1}{B}\int_{-\delta}^{\delta}\left[1 - \frac{\theta^2}{B^2}\right]d\left(\frac{\theta}{B}\right) \cong \frac{1}{B}$$

for θ/B much less than one. Hence, from Eq. (31) we derive:

$$I_p = \frac{k\omega_r f_{e0}\left(\frac{\omega_r}{k}\right) + \frac{k^4\omega_r \kappa T_e \nu}{\pi m_e \omega_r^4}}{-2\omega_p^2 \frac{\partial f_{e0}(\mathrm{v})}{\partial \mathrm{v}}\bigg|_{\mathrm{v}=\frac{\omega_r}{k}} + \frac{2\omega_p^2 \nu k^2}{\pi \omega_r^3}}. \tag{34}$$

Since nonthermal electrons constitute only a minor proportion of available electrons, we represent them as a perturbation, $f_p(\mathrm{v})$, on the Maxwellian distribution of thermal electrons, $f_m(\mathrm{v})$. That is,

$$f_{e0}(\mathrm{v}) \Rightarrow f_m(\mathrm{v}) + f_p(\mathrm{v}). \tag{35}$$

We describe the intensity in terms of an effective temperature, T_p, analogous to the value, $T_{\parallel}$ defined by Perkins and Salpeter[1],

$$m\mathrm{v}/kT_{\parallel} \equiv -d \ln f_{e0}(\mathrm{v})/d\mathrm{v} \tag{36}$$

Note Eq. (34) reduces to $I_p = \dfrac{T_{\parallel}}{2\alpha^2 T_e}$ when $\nu = 0$ and $T_{\parallel}$ reduces to T_e for $f_{e0} = f_m$. We define an effective temperature, T_p, such that

$$I_p = \frac{T_p}{2\alpha^2 T_e} \text{ for } \nu \neq 0. \tag{37}$$

Also we define

$$\chi \equiv \frac{k^3 \kappa T_e}{\pi m_e \omega_r^4}\nu \tag{38}$$

Incorporating Eqs. (35), (37) and (38) into Eq. (34) we may write

$$2\alpha^2 I_p = \frac{T_p}{T_e} = \frac{f_m(\omega_r/\kappa) + f_p(\omega_r/\kappa) + \chi}{f_m(\omega_r/\kappa) + L_p(\omega_r/\kappa) + \chi} \tag{39}$$

where

$$L_p \equiv \left(\frac{\kappa}{\omega_r}\right)\frac{\kappa T_e}{m_e}\,\partial f_p(\mathrm{v})/\partial \mathrm{v}\Big|_{\mathrm{v}=\omega_r/\kappa}. \tag{40}$$

The term L_p represents a loss of energy due to changes in the nonthermal electron distribution.

The effective collison frequency, ν, given in Eq. 3, is the sum of electron-neutral and ion-neutral collision frequencies. Our calculation of the effective electron temperature associated with a plasma line, T_p in Eq. (38), reduces to Eq. (1) if the electron-neutral collisions are negligible.

V. CONCLUSION

Electron-neutral collisons are only important in the lower ionosphere, where the terms included in Eq. (38) vary as shown in Fig. 1. Although the electron-neutral collision terms are small, they can have a significant effect on the effective temperature at low altitudes. This is seen in Fig. 2. The solid curves are calculated using Eq. (38), while the dashed curves have been calculated neglecting electron-neutral collisions. Experimental data from Chatanika[8], plotted on the same graphs, reflect a much better correspondence to the calculations which include electron-neutral collisions than to the collison-less calculations.

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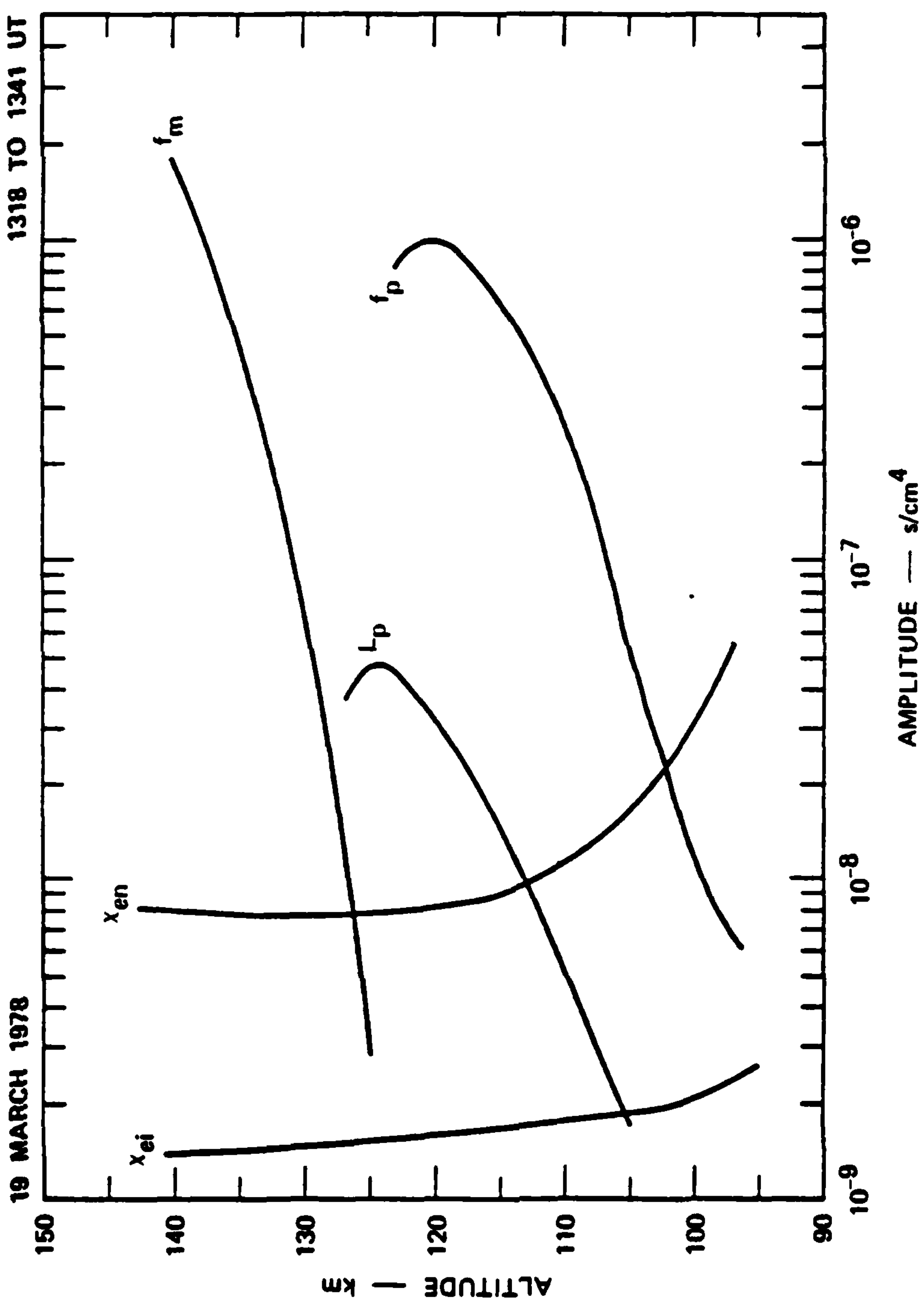

Fig. 1 — Calculated excitation and damping terms for the plasma wave temperatures

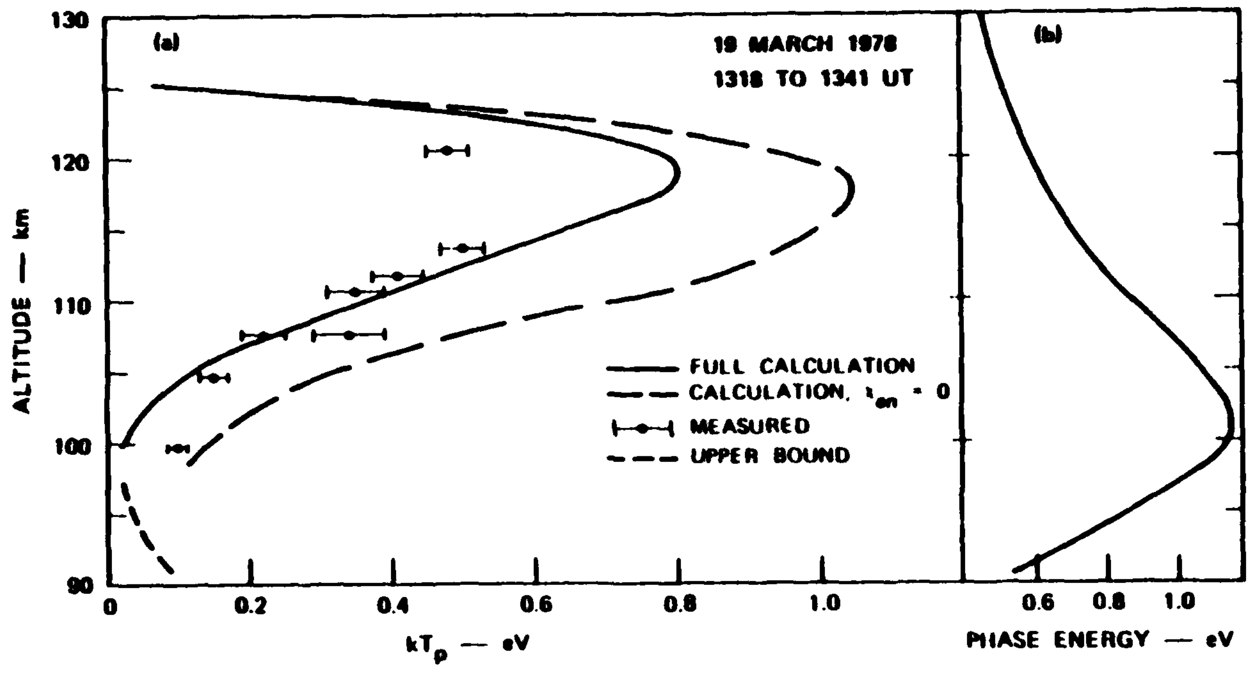

Fig. 2 — Plasma wave temperatures for 1318-1341 UT on 19 March 1978. The circles show the measured kT_p values in the topside of the E layer. The dashed curve in the bottomside shows the lower temperature limit for detecting plasma lines or an equivalent upper bound for the temperatures. The heavy dashed line shows the theoretically calculated kT_p values without electron-neutral collisions, the heavy solid line shows the temperatures with collisions.