Model Calculations of the Dayside Ionosphere of Venus: Energetics

T. E. CRAVENS, T. I. GOMBOSI, J. KOZYRA, AND A. F. NAGY

Space Physics Research Laboratory, Department of Atmospheric and Oceanic Science, University of Michigan
Ann Arbor, Michigan 48109

L. H. BRACE

NASA Goddard Space Flight Center, Greenbelt, Maryland 20771

W. C. KNUDSEN

Lockheed Palo Alto Research Laboratory, Palo Alto, California 94304

A model of the energy balance of the dayside ionosphere of Venus is presented. Calculations of the dayside electron and ion temperature profiles are carried out and compared with data from experiments on the Pioneer Venus orbiter. The coupled heat conduction equations for electrons and ions are solved for several values of the solar zenith angle. It is shown that thermal conductivities are inhibited by the presence of a horizontal magnetic field. A realistic model of the magnetic field that includes fluctuations is employed in deriving an appropriate expression for the thermal conductivity. The contributions of photoelectrons, ion chemistry, Joule heating, and solar wind heating to the energy balance of the ionosphere are considered.

1. INTRODUCTION

The physical nature of the ionosphere of Venus is now becoming better understood both because of the growing data base from the Pioneer Venus orbiter (PVO) and because of theoretical model calculations utilizing this data base. Prior to Pioneer Venus, numerous models were developed, and these are referenced both by Schunk and Nagy [1980] and by Cravens et al. [1979]. This paper as well as the companion paper on the ion composition [Nagy et al., this issue] are extensions of the calculations of Chen and Nagy [1978] and Cravens et al. [1978] done prior to Pioneer Venus and of Cravens et al. [1979] and Nagy et al. [1979] immediately after the Pioneer Venus encounter.

In this paper the coupled heat conduction (or energy) equations for electrons and ions are solved for several values of the solar zenith angle. In the companion paper [Nagy et al., this issue] the coupled continuity and momentum equations for seven major ionic species are solved. Ideally, the continuity, momentum, and energy equations would all be solved simultaneously; however, until the basic physical and chemical processes relevant to the ionosphere of Venus are better understood, this would be premature. Instead, in this paper the energy equations are solved by using measured values for ion composition. Similarly, in Nagy et al. [this issue], measured values of electron and ion temperatures are adopted, and ion densities are calculated.

Using ion composition and neutral densities measured during the first few orbits of PVO, Cravens et al. [1979] calculated electron and ion temperatures both with and without considering the effects of a quasi-horizontal magnetic field on the thermal conductivities. They found that without a magnetic field the calculated temperatures were much lower than the measured temperatures unless one includes external heat inputs at the ionopause of $5 \times 10^{7}$ eV cm$^{-2}$ s$^{-1}$ for electrons and $7 \times 10^{7}$ eV cm$^{-2}$ s$^{-1}$ for ions. Heat fluxes of the same order of magnitude were found to be necessary to reproduce the observed daytime electron and ion temperatures by Knudsen et al. [1979b]. However, if a quasi-horizontal magnetic field is included, the necessity for external heat input is either eliminated or reduced [Cravens et al., 1979; Knudsen et al., 1979b]. The magnetic field used in previous models was an oversimplified one: a magnetic field uniform both in magnitude and direction. The measured magnetic field is highly variable or fluctuating, even turbulent, with many complicated spatial structures [Russell et al., 1979]. In the present paper a much more realistic model of the magnetic field will be employed and an appropriate expression for the thermal conductivity will be developed. Photoelectron fluxes in the ionosphere of Venus are also calculated. In addition, this paper will explore the variation of calculated electron and ion temperatures with solar zenith angle.

Most ionospheric parameters measured by PV, such as electron temperature and density, exhibit very large orbit to orbit variations on the nightside and moderate variations on the dayside. In this paper the ionospheric energetics for average or typical dayside conditions at four solar zenith angles are explored. The degree of variability of the parameters being considered is small enough so that one can choose between (or put bounds on) the physical mechanisms that we are invoking to explain the energetics of the dayside ionosphere. Until these mechanisms are better understood a detailed orbit by orbit theoretical study is premature.

2. DETAILS OF THE MODEL CALCULATION

The Equations

By using the numerical technique outlined by Hastings and Roble [1977], solutions of the coupled electron and ion energy equations were obtained over an altitude range from 120 to 500 km. The following equation [Banks and Kockarts, 1973] was solved for electrons and ions simultaneously:

$$\frac{3}{2} n_e k \frac{\partial T_e}{\partial t} - \frac{\partial}{\partial z} \left[ K_m \frac{\partial T_m}{\partial z} \right] = Q_m - S_m$$  (1)
where

- $m$: index for electrons or ions;
- $z$: vertical coordinate;
- $t$: time;
- $T_e$: electron or ion temperature;
- $n_e$: electron density;
- $K_e$: electron or ion thermal conductivity;
- $Q_e$: electron or ion heating rate;
- $S_e$: electron or ion cooling rate.

Bulk transport terms are not expected to be too important, at least at altitudes well below the ionopause on the dayside, although the possibility of Joule heating due to bulk motions will be considered briefly. The cooling processes used were outlined in *Chen and Nagy [1978]* and *Chen [1977]*, although some changes were introduced by *Cravens et al. [1979]*. In particular, the O(3P), fine structure cooling rate is from *Hoegy [1976]*, the O(D) cooling rate is from *Banks and Kockarts [1973]*, and the CO and CO$_2$ vibrational cooling rates are adopted from *Porter and Mayr [1979]*. The basic expressions for the standard thermal conductivities are from *Banks and Kockarts [1973]*. The ion conductivity, as well as the electron conductivity, was corrected for damping by the neutral gas, using a procedure discussed in *Banks and Kockarts [1973]*. Later in this section there will be a discussion of how the standard expressions for thermal conductivity can be modified to take into account the effects of a fluctuating magnetic field.

**Neutral Atmosphere and Ionosphere Parameters**

Four neutral models appropriate for average conditions at solar zenith angles of 0°, 60°, 80°, and 90° were constructed by using neutral density values consistent with data from the neutral mass spectrometer aboard the PV orbiter [Niemann et al., this issue]. The orbit to orbit variation of neutral density is considerably less than 50% for the dayside at an altitude of 150 km, and the ‘average’ models used here should be appropriate. The reader is referred to *Niemann et al. [this issue]* for information on these neutral densities. Additional details are considered in the companion paper by *Nagy et al.* [this issue].

In order to solve (1) for electron and ion temperatures, knowledge of ion densities as functions of altitude and solar zenith angle is necessary. For this purpose, ion density values measured by the ion mass spectrometer on PVO [Taylor et al., 1980] were used for the four values of solar zenith angle being considered. The data used were from orbits 126, 134, 147, and 184, and since the orbit to orbit variations of the measured ion densities are relatively small on the dayside, these four orbits should be approximately representative of the average conditions at the respective solar zenith angles. Below the PVO perigee altitude of about 150-160 km, no data were available, and consequently, we used model calculations [Nagy et al., 1979, this issue] as a guide for extrapolating the measured densities. For the type of calculations considered here, only the O$^+$ and O$_2^+$ densities are of major importance.

**Thermal Conductivity**

In the absence of collisions, charged particles in a uniform magnetic field will move in helical paths and stay ‘frozen’ to the magnetic field lines. Collisions enable these charged particles to diffuse perpendicular to the field lines. *Cravens et al. [1979]* incorporated a uniform magnetic field into their temperature model and showed that the subsequent reduction of the thermal conductivities resulted in elevated electron and ion temperatures. The ratio of the thermal conductivity for a strictly horizontal magnetic field relative to the ‘standard’ conductivity is [Hocksim and Massel, 1969; cf. Johnson, 1978]

$$F_D = \left[ \frac{F_{e,D}}{\Omega_e h} \right]^{-2}$$

where $F_{e,D}$ is the average total collision frequency for species $m$, $\Omega_e$ is the gyrofrequency, and $g$ and $h$ are slowly varying correction factors ($g/h$ is about .25 for ions and .14 for electrons). Above an altitude of about 200 km, the ratio of the electron collision frequency to the gyrofrequency becomes extremely small, even for a very small magnetic field, and there is very little electron heat flow perpendicular to the field lines. If only (2) were relevant, then the electron temperature would be at least an order of magnitude higher than what is observed.

Dayside measurements from the magnetometer on PVO [Russell et al., 1979] indicate a magnetic field that is highly variable in both magnitude and direction; sometimes, the ‘typical’ field is larger than 10 γ and at other times it is smaller than 2 γ. The magnetic field usually has a very large fluctuating component with a fluctuation correlation length probably less than 10 km (R. Elphic and C. T. Russell, private communication, 1980). To obtain the thermal conductivity of a charged particle in an irregular or fluctuating magnetic field, the value of the mean free path of the particle in this fluctuating field is required. For fluctuations of the magnetic field of the same order as or smaller than the average magnitude of the field, then the results of a quasi-linear theory of charged particle motion can be used to obtain values for the effective mean free path $\lambda$. This quasi-linear theory is discussed in the appendix. Given the mean free path $\lambda$, the thermal conductivity can be written as [Banks and Kockarts, 1973]

$$K_f = \frac{3}{4} \pi n_e k \bar{V} \lambda$$

where $\bar{V}$ is the average thermal speed, $n_e$ is the electron density, $k$ is Boltzmann’s constant, and $\pi$ is a correction factor of order unity (.66 is used here) that takes into account thermo-electric effects. Equation (3) predicts that as $\lambda$ goes to zero, the conductivity goes to zero. In this case the conductivity should actually be the standard conductivity ($K_s$) multiplied by $F_D$ (equation (2)). As $\lambda$ becomes infinite, $K_f$ becomes infinite, whereas the conductivity should really be just the standard conductivity $K_s$. A phenomenological expression that incorporates all the correct asymptotic limits is

$$K = K_s \left[ \frac{K_f + F_D K_s}{K_f + K_s} \right]$$

where $K$ is the final form of the thermal conductivity. Notice that if there is no magnetic field, $K$ becomes just $K_s$ no matter what the value of $\lambda$ is.

Quasi-linear theory is an approximate guide to the transport properties of charged particles for fluctuations of about the same magnitude as the average magnetic field. If the gyroradius of a particle is less than the correlation length of the fluctuations, then the theory predicts that the mean free path of the particle is approximately equal to the correlation length [(A13) in the appendix]. On the other hand, if the gyroradius is considerably larger than the correlation length, one can show with quasi-linear theory (using an equation similar to (A10) in the appendix) that the mean free path is of the same...
order as or larger than the gyroradius. If the magnitude of the fluctuations is much larger than the magnitude of the average magnetic field, then quasi-linear theory is not valid. In this case, intuition suggests that if the gyroradius is smaller than the correlation length, then the mean free path is still about equal to the correlation length. However, in this case, if the gyroradius is much larger than the correlation length, the mean free path will be extremely large and the particles do not 'feel' the magnetic field.

The gyroradius of a thermal electron ($\sim \frac{1}{2}$ eV) in a magnetic field of $2\gamma$ is only about 1 km, and if the magnitude of the magnetic fluctuations is the same order as the value of the average magnetic field strength, then the mean free path of a thermal electron is approximately the same as the correlation length if this is larger than 1 km. At this moment our knowledge of the exact nature of the magnetic fluctuations is very limited, therefore a range of possible $\lambda$ values will be used in this paper. Values between 1 km and 10 km are favored at this time.

The gyroradius of a thermal O$^+$ ion ($\sim 15$ eV) in a $2\gamma$ field is 110 km and in a $10\gamma$ field is 22 km. For an average horizontal magnetic field strength of $2\gamma$, quasi-linear theory gives a mean free path of about 100 km, which is quite large. Quasi-linear theory assumes that the plasma is collisionless. However, for a $2\gamma$ field the collision frequency everywhere in the ionosphere is large enough to decouple the ions from the field, and consequently, the magnetic field (with or without fluctuations) will have little influence on ion transport. For a $10\gamma$ field the mean free path is about 20 km, but even for $B = 10\gamma$ ion collision frequencies for zenith angles less than $80^\circ$ are large enough to decouple effectively the ions from the magnetic field because of the rather large ion densities at these zenith angles. The results of Cravens et al. [1979] for a uniform magnetic field indicate that calculated ion temperatures are only marginally affected by magnetic field strengths less than $10\gamma$. Later in this paper we will employ a range of $\lambda$ values for ions, but one should bear in mind that for ions magnetic fluctuations are not too important.

**Photoelectrons and the Electron Heating Rate**

The rate at which photoelectrons heat the ambient thermal electron population is obtained from the calculated photoelectron flux. The two-stream photoelectron transport method of Nagy and Banks [1970] was modified for Venus and was used to calculate the photoelectron fluxes. The photoelectron production rates were calculated by using the measured solar EUV flux appropriate for 1979 solar conditions, using measured photoionization and photoabsorption cross sections, and using measured neutral densities for each of the solar zenith angles being considered. This material is discussed in more detail in Nagy et al. [1980].

Electron impact cross sections for the excitation and ionization of CO$_2$ and CO were taken from Sawada et al. [1972a, b]. The cross sections for atomic oxygen were taken from Green and Stolarski [1972] and for ionization and dissociative ionization of N$_2$ from Banks and Kockarts [1973]. The expression for the electron-electron energy transfer rate was given by Swartz et al. [1971]. The electron density and temperature profiles required to calculate electron heating rates are those provided by the empirical model of Theis et al. [this issue], which is based upon measurements from the Langmuir probe instrument (OETP) aboard PVO. Below about 150 km, this data-based model does not provide information, therefore the electron densities and temperatures were extrapolated. For the photoelectron flux and heating rate calculations, the energy bin width was chosen to be 0.5 eV for energies less than 10 eV and 1.0 eV for energies greater than 10 eV.

The work of Butler and Stolarski [1978] established that even very small quasi-horizontal magnetic fields will inhibit the vertical transport of photoelectrons and force them to deposit their energy locally up to fairly high altitudes. For this reason, most of the calculations presented in Cravens et al. [1979] assumed local energy deposition. In this paper the effects of magnetic field fluctuations on photoelectron transport is incorporated in a very approximate and phenomenological manner. Photoelectrons will be 'scattered' and will at least partially decouple from the average horizontal magnetic field if the amplitude of the fluctuations is large enough. If the gyroradius of a typical photoelectron (about 3 km for a $10\gamma$ field and 15 km for a $2\gamma$ field) is equal to or smaller than the correlation length of the magnetic fluctuations, then photoelectron transport can be described as a random walk. If the mean free path due to this fluctuating magnetic field is about 10 km, then for a random walk the total path length for a photoelectron is about thirty times larger than the vertical extent of the ionosphere.

To simulate crudely the effects of such an increase in path length, the two-stream transport method with a dip angle of $2^\circ$ was employed. For randomly scattered particles, the pitch angle distribution will be approximately isotropic, and consequently, the upward and downward photoelectron fluxes ($\phi^+$ and $\phi^-$, respectively) will be approximately equal. If the magnetic field lines on which the photoelectrons are traveling in the $2^\circ$ dip case are 'scrambled' in order to create a random magnetic field (keeping the total length of these field lines the same), then the $\phi^+$ ($\phi^-$) calculated for $2^\circ$ dip will now be seen as both an up flux and a down flux with magnitude $\phi^+/2$ ($\phi^-/2$). Actually, even for this random walk problem, the upward flux will still be slightly larger than the downward flux by a couple percent; this small anisotropy is neglected here. The photoelectron flux calculated for inhibited vertical transport is shown in Figure 1 for an energy of 25.5 eV and for a solar zenith angle of 0°. The same photoelectron flux is shown in Figure 2 as a function of energy for an altitude of 304 km. Notice that in Figure 2 the flux units are 'per steradian.'

If the magnitude of the fluctuations is quite large and if the photoelectron gyroradius is considerably larger than the correlation length of those fluctuations, then the photoelectrons...
will not 'feel' the magnetic field and their vertical transport will be uninhibited. In this case the two-stream method with a dip angle of 90° can be used. The calculated upward and downward fluxes are shown in Figures 1 and 2. From Figure 1 it is clear that for uninhibited photoelectrons the flux is very anisotropic at high altitudes and that the total flux is larger than the total flux for inhibited photoelectrons. Figure 2 illustrates that at lower energies there is a large difference between the fluxes calculated with the two different assumptions about vertical transport. At higher altitudes and lower energies, energy loss to thermal electrons is the dominant process. The column density of thermal electrons seen by the inhibited photoelectrons is thirty times longer than that seen by the uninhibited ones and therefore the former lose energy more rapidly than the latter, especially at low energies. The photoelectron flux measured near the subsolar point by the retarding potential analyzer (RPA) on PVO [Knudsen et al., this issue] agrees fairly well with the flux calculated, assuming no inhibition of transport. This implies that at least this one occasion the total magnetic field was small enough or the fluctuations were large enough and the correlation length small enough so that the vertical transport of photoelectrons was not inhibited.

The heating rate of thermal electrons due to photoelectrons is shown in Figure 3 for four zenith angles and for no inhibition of vertical transport. The typical orbit to orbit variability is also shown as are some RPA electron temperature measurements [Miller et al., this issue]. The electron heating rate for uninhibited photoelectrons is used here, and the magnetic field strength is set at 10 nT. Values of the mean free path range from 3 km to infinity—an infinite \( \lambda \) can result from either large fluctuations with very small correlation length or from a zero field strength. The electron temperatures calculated here for infinite \( \lambda \) (or \( B = 0 \)) are clearly much smaller than the measured OETP values, which are about 5000°K at 400 km. As \( \lambda \) decreases, the thermal conductivity decreases and heat is less readily transported from higher altitudes to lower altitudes where the electrons are cooled most effectively. Consequently, as \( \lambda \) decreases, the electron temperature increases. If \( \lambda \) is chosen to be 3 km, then

\[
\begin{align*}
\text{O}^+ + \text{CO}_2 & \rightarrow \text{O}_2^+ + \text{CO} + 1.2 \text{ eV} \\
\text{CO}_2^+ + \text{O} & \rightarrow \text{O}_2^+ + \text{CO} + 1.4 \text{ eV}
\end{align*}
\]

Rohrbaugh et al. [1979] suggested these reactions as a major heat source for the ionosphere of Mars. Their application to the ionosphere of Venus was discussed in Cravens et al. [1979]. Since the ion densities measured by PVO show considerable variability, the chemical heating rates calculated at each zenith angle are subject to some uncertainty.

Taylor et al. [this issue] and Knudsen et al. [this issue] have made preliminary measurements of ion drift velocities that are the order of kilometers per second at altitudes near the ionopause. Depending upon how rapidly these velocities decrease as one goes deeper in the atmosphere, these drifts might be a significant source of 'Joule' heat for the ions. At this time we can only make a very crude estimate of the potential effects of this Joule heating on the ion temperature. For this purpose we postulate a drift velocity that is 2 km/s at 300 km and decreases to 322 m/s at 200 km and 22 m/s at 120 km. Later in this paper the ion temperature that results from this postulated Joule heat source is discussed.

3. RESULTS

Electron and ion temperatures that are calculated for the subsolar point are shown in Figure 4. Measured values of electron temperature from the empirical model of Theis et al. [this issue] are shown in Figure 4 and in other figures. The typical orbit to orbit variability is also shown as are some RPA electron temperature measurements [Miller et al., this issue]. The electron heating rate for uninhibited photoelectrons is used here, and the magnetic field strength is set at 10 nT. Values of the mean free path \( \lambda \) range from 3 km to infinity—an infinite \( \lambda \) can result from either large fluctuations with very small correlation length or from a zero field strength. The electron temperatures calculated here for infinite \( \lambda \) (or \( B = 0 \)) are clearly much smaller than the measured OETP values, which are about 5000°K at 400 km. As \( \lambda \) decreases, the thermal conductivity decreases and heat is less readily transported from higher altitudes to lower altitudes where the electrons are cooled most effectively. Consequently, as \( \lambda \) decreases, the electron temperature increases. If \( \lambda \) is chosen to be 3 km, then

\[
\begin{align*}
\text{O}^+ + \text{CO}_2 & \rightarrow \text{O}_2^+ + \text{CO} + 1.2 \text{ eV} \\
\text{CO}_2^+ + \text{O} & \rightarrow \text{O}_2^+ + \text{CO} + 1.4 \text{ eV}
\end{align*}
\]

Rohrbaugh et al. [1979] suggested these reactions as a major heat source for the ionosphere of Mars. Their application to the ionosphere of Venus was discussed in Cravens et al. [1979]. Since the ion densities measured by PVO show considerable variability, the chemical heating rates calculated at each zenith angle are subject to some uncertainty.

Taylor et al. [this issue] and Knudsen et al. [this issue] have made preliminary measurements of ion drift velocities that are the order of kilometers per second at altitudes near the ionopause. Depending upon how rapidly these velocities decrease as one goes deeper in the atmosphere, these drifts might be a significant source of 'Joule' heat for the ions. At this time we can only make a very crude estimate of the potential effects of this Joule heating on the ion temperature. For this purpose we postulate a drift velocity that is 2 km/s at 300 km and decreases to 322 m/s at 200 km and 22 m/s at 120 km. Later in this paper the ion temperature that results from this postulated Joule heat source is discussed.

3. RESULTS

Electron and ion temperatures that are calculated for the subsolar point are shown in Figure 4. Measured values of electron temperature from the empirical model of Theis et al. [this issue] are shown in Figure 4 and in other figures. The typical orbit to orbit variability is also shown as are some RPA electron temperature measurements [Miller et al., this issue]. The electron heating rate for uninhibited photoelectrons is used here, and the magnetic field strength is set at 10 nT. Values of the mean free path \( \lambda \) range from 3 km to infinity—an infinite \( \lambda \) can result from either large fluctuations with very small correlation length or from a zero field strength. The electron temperatures calculated here for infinite \( \lambda \) (or \( B = 0 \)) are clearly much smaller than the measured OETP values, which are about 5000°K at 400 km. As \( \lambda \) decreases, the thermal conductivity decreases and heat is less readily transported from higher altitudes to lower altitudes where the electrons are cooled most effectively. Consequently, as \( \lambda \) decreases, the electron temperature increases. If \( \lambda \) is chosen to be 3 km, then

\[
\begin{align*}
\text{O}^+ + \text{CO}_2 & \rightarrow \text{O}_2^+ + \text{CO} + 1.2 \text{ eV} \\
\text{CO}_2^+ + \text{O} & \rightarrow \text{O}_2^+ + \text{CO} + 1.4 \text{ eV}
\end{align*}
\]

Rohrbaugh et al. [1979] suggested these reactions as a major heat source for the ionosphere of Mars. Their application to the ionosphere of Venus was discussed in Cravens et al. [1979]. Since the ion densities measured by PVO show considerable variability, the chemical heating rates calculated at each zenith angle are subject to some uncertainty.

Taylor et al. [this issue] and Knudsen et al. [this issue] have made preliminary measurements of ion drift velocities that are the order of kilometers per second at altitudes near the ionopause. Depending upon how rapidly these velocities decrease as one goes deeper in the atmosphere, these drifts might be a significant source of 'Joule' heat for the ions. At this time we can only make a very crude estimate of the potential effects of this Joule heating on the ion temperature. For this purpose we postulate a drift velocity that is 2 km/s at 300 km and decreases to 322 m/s at 200 km and 22 m/s at 120 km. Later in this paper the ion temperature that results from this postulated Joule heat source is discussed.
the calculated electron temperatures are close to the measured values, although the shape of the calculated electron temperature profile is not quite correct.

The calculated ion temperatures, which are shown in Figure 4, are almost independent of the value of \( \lambda \) chosen. For a 10 \( \gamma \) field the ion collision frequency near the subsolar point is of the same order as the ion gyrofrequency, and consequently, ion heat can be conducted perpendicular to the field lines rather readily. As is indicated in Cravens et al. [1979], only when the magnetic field strength is considerably larger than 10 \( \gamma \) is the ion thermal conductivity significantly reduced. Notice that the ion temperature measured between 160 km and 200 km is fairly constant. Without chemical heating the calculated ion temperatures are significantly lower than the measured values. Ion heating, using the assumed chemical processes, is sufficient to reproduce the measured temperatures at altitudes less than 200 km. However, it should be remembered that there are uncertainties associated with the chemical heat source; the potential contribution of Joule heating is discussed later.

On the dayside the measured ion temperature [Knudsen et al., 1979a; Miller et al., this issue] shows considerable orbit to orbit variability, an approximate indication of which is given in the figures in this paper. This variability is still much less than on the nightside, and an average measured temperature is a meaningful quantity. The data shown in this paper are appropriate for all dayside zenith angles since the measured variation of ion temperature with zenith angle is small [Miller et al., this issue]. Since the measured ion temperature is rather variable and since the calculated ion temperature below 200 km depends on the adopted neutral and ion densities, which are also somewhat variable one cannot be certain at this moment whether the chemical heat source below 200 km by itself is sufficient without an additional heat source such as Joule heating.

Figure 5 also shows measured and calculated subsolar temperatures; the difference from Figure 4 is that the electron heating rate for inhibited rather than uninhibited photoelectron transport was employed. Even though the thermal conductivity for a given \( \lambda \) is the same as in Figure 4, because the heating rate at higher altitudes used for Figure 5 is much smaller than for Figure 4, the altitude integrated heating rate (or heat flux) is smaller for Figure 4, and therefore the temperature gradients are smaller. The heat flux for the \( m \)th species \( q_m \) is

\[
q_m = -K_{en} \frac{dT_m}{dz}
\]

where \( K_{en} \) is the electron conductivity. If the conductivity \( K_{en} \) is reduced, then a larger temperature gradient is required to support the same heat flux. If photoelectrons are inhibited, then \( \lambda \) must be .3 km in order for the calculated \( T_e \) to be as large as measured values. An electron mean free path value of .3 km is a factor of 10 smaller than the \( \lambda \) value required for the calculations with the uninhibited photoelectron heating rate. The degree to which photoelectron transport is inhibited is an important factor in the energy balance of thermal electrons. At the moment, the preferred choice is uninhibited photoelectron transport (at least for some orbits) since the measured photoelectron fluxes indicate this [Knudsen et al., this issue].

Figure 6 shows how the calculated electron temperature varies with zenith angle for zero magnetic field strength (or \( \lambda = \infty \)). The variation of the average measured electron temperature [Theis et al., this issue] with solar zenith angle is indicated, whereas in previous figures the typical day to day

Fig. 4. Electron, ion, and neutral temperature profiles are shown for the subsolar point. The magnetic field is 10 \( \gamma \) and temperature profiles for a range of \( \lambda \) are shown. Uninhibited photoelectron transport is assumed. Electron temperature data from the Langmuir probe (OETP) and ion temperature data from the retarding potential analyzer (ORPA) are shown. Some ORPA electron temperatures are also shown. The same symbols will be used in other figures.

Fig. 5. This figure is similar to Figure 4, except that inhibited photoelectron transport is assumed.

Fig. 6. Calculated electron temperature profiles are shown for 0°, 60°, 80°, and 90° solar zenith angles. The magnetic field is assumed to be zero, and uninhibited photoelectron transport is assumed. The range of the average OETP temperature measurements with respect to zenith angle is indicated, rather than the orbit to orbit variation at a particular zenith angle as in Figures 4 and 5.
temperatures are much larger than the calculated temperatures assuming zero magnetic field strength, which decreases from 6000°K at the subsolar point to 3350°K at the terminator. The measured temperatures are much larger than the calculated temperatures at all zenith angles.

Figure 7 shows how the calculated electron temperature varies as a function of zenith angle for a magnetic field strength of 10 γ and an electron mean free path of 3 km. The calculated temperatures shown in Figure 7 are all larger than those calculated assuming zero magnetic field strength, which are shown in Figure 6. At 400 km, the calculated temperature decreases from 6000°K at the subsolar point to 3350°K at the terminator.

The variation of the calculated ion temperature as a function of solar zenith angle is shown in Figure 8, both for no magnetic field (or \( \lambda = \infty \)) and for \( B = 10 \gamma \) with a 10-km ion mean free path. Only the envelope or range of the calculated ion temperatures is shown rather than individual profiles. The reason for this is that both neutral and ion densities are smaller near the terminator so that the ratio of the collision frequency to gyrofrequency is also smaller and the ions are more closely 'tied' to magnetic field lines. When calculations were performed (but not shown) in which \( B = 2 \gamma \) rather than \( 10 \gamma \), the resulting ion temperatures were reduced almost to their \( B = 0 \) values at all zenith angles.

In section 2 it was stated that the electron mean free path in the fluctuating ionospheric magnetic field on Venus is probably larger than 10 km. However, if the value of \( \lambda \) is really much larger than 10 km, then the electron conductivity is relatively unaffected and a significant additional input of energy is required to produce electron temperatures as large as the measured ones. The electric field experiment on PVO [Taylor et al., 1979] has detected whistler wave signals in the 'magnetosheath' region, which are certainly damped at the ionopause. At that time the fraction of wave energy reflected at the ionopause could not be established, and they could only set an upper limit of about \( 3 \times 10^{10} \text{ eV cm}^{-2} \text{s}^{-1} \) on the heat flux deposited into the electron gas via Landau damping at the ionopause. Calculations were carried out that include heat inputs such as this.

If \( B \) is taken to be zero (or \( \lambda = \infty \)), then in order for the electron temperatures calculated at 60° zenith angle to agree with the measured values, one must include a heat flux of \( 3 \times 10^{10} \text{ eV cm}^{-2} \text{s}^{-1} \) at the upper boundary. The required heat flux for ions is \( 5 \times 10^9 \text{ eV cm}^{-2} \text{s}^{-1} \). The calculated \( T_e \) and \( T_i \) profiles for this case are shown in Figure 9. Since the required electron heat flux is close to the upper limit for whistler wave heating, a smaller heat flux appears to be more realistic. The heat flux requirements can be reduced by including a nonzero magnetic strength and a finite mean free path. For \( B = 10 \gamma \) and \( \lambda = 10 \text{ km} \) the heat flux required for electrons is reduced by almost an order of magnitude to \( 5 \times 10^9 \text{ eV cm}^{-2} \text{s}^{-1} \), and the heat flux required for ions is reduced by almost a factor of 2. The results for this case are shown in Figure 10. The shape of the \( T_e \) profile shown in Figure 10 could be better, however.

The shape of the calculated \( T_e \) profile shown in Figure 10 can be improved if \( \lambda \) varies with altitude instead of being constant. All that was stated about \( \lambda \) in section 2 (and the Appendix) is that magnetometer results indicate that it is probably in the range of 1–10 km. This range of \( \lambda \) is still rather speculative, and it is quite likely that \( \lambda \) really is a function of altitude. We will now assume that \( \lambda \) increases with altitude as follows:

\[
\lambda(z) = \frac{1}{\lambda} \text{ km} \text{ e}^{(-z/120 \text{ km})/30 \text{ km}}
\]

Fig. 7. This figure is similar to Figure 6, except that the magnetic field is \( 10 \lambda \) and the mean free path is 3 km.

Fig. 8. The complete range or envelope of calculated dayside ion temperatures are shown by means of cross-hatching both for zero magnetic field and for \( B = 10 \gamma \) and \( \lambda = 3 \text{ km} \).

Fig. 9. The calculated electron and ion temperatures are shown for zero magnetic field and 60° zenith angle. Heat inputs at the upper boundary are included for both the electrons and ions.
Using this $\lambda(\mu)$ for electrons, $\lambda = 10$ km for ions, $B = 10^7$, an electron heat input of $4 \times 10^9$ eV cm$^{-2}$ s$^{-1}$, and no ion heat input, we obtain the temperature profiles shown in Figure 11.

Also shown in Figure 11 is an ion temperature profile calculated without any external heat flux but with a Joule heating contribution calculated by using an expression from Banks and Kockarts [1973]. The ion drift velocities chosen were described in section 2 and reached values of several km/s near 400 km. Although this choice of drift velocity is rather arbitrary, the values used here are roughly consistent with the small amount of ion drift data now available [Knudsen et al., this issue]. It is clear from the ion temperature results shown in Figure 11 that Joule heating can make a significant contribution to the energy balance of the thermal ions.

4. DISCUSSION

Detailed comparisons of calculated and measured electron and ion temperatures showed reasonable agreement when a suitable set of parameters was chosen. Included in the calculations were photoelectron heating, ion-neutral chemical heating, Joule heating, and heat inputs attributed to the solar wind-ionosphere interaction. A fairly realistic model of the ionospheric magnetic field was used in the construction of thermal conductivity coefficients that reflect the highly variable nature of the magnetic field. The chief obstacle to understanding the energy balance of the dayside ionosphere at this point in the future is the dearth of detailed information on the nature of the magnetic fluctuations—such as the amplitude of the fluctuations and their correlation length. When accurate power spectra of the field are eventually determined from magnetometer data, then our understanding of the transport properties of various particle populations in the ionosphere (such as thermal electrons, photoelectrons, and thermal ions) will improve dramatically.

It is quite likely that the transport properties of various particles will be highly variable, because the nature of the magnetic field appears to change from orbit to orbit. For instance, when the field is quite small and/or very highly variable, then the transport of photoelectrons and ions, and possibly even thermal electrons, will be relatively uninhibited. On the other hand, on those occasions when the field is large and relatively coherent, the vertical transport of all types of particles will probably be inhibited. In this paper we have made a general study appropriate for average daytime conditions in order to sort out the importance of some basic processes. At some point in the future it will be advantageous to do detailed studies of individual orbits, utilizing simultaneous data from many PV instruments at once such as the magnetometer, the Langmuir probe, the retarding potential analyzer, the ion mass spectrometer, the neutral mass spectrometer, and the electric field detector.

APPENDIX

In this appendix a quasi-linear theory of particle transport in a fluctuating magnetic field is used to derive an expression for the effective mean free path of these particles. For some conditions this mean free path can then be used in the expression for thermal conductivity which is discussed in section 2. The details of this theory can be found in the review paper on the propagation of cosmic rays in the solar wind by Jokipii [1971].

One starts with the Fokker-Planck equation in pitch angle and position. The particle density as a function of pitch angle $\mu$, position r, and time t, can be represented in terms of an 'average' nonfluctuating part $\langle U(r, t) \rangle$ and an expansion over Legendre functions of the pitch angle. By assuming that the fluctuating part of the density is small, one can retain only the first term of the expansion. The density is then represented by

$$n(\mu, \chi, t) = \langle U(\mu, r, t) \rangle + \delta n(\mu, r, t)$$

where we assume that $n_t = \langle U \rangle$.

What causes these density fluctuations are the fluctuations in the magnetic field. The magnetic field $B(\mu, r, t)$ can be represented in terms of an average part $B_0$ and a fluctuating part $B_1(\mu, r, t)$:

$$B(\mu, r, t) = B_0 + B_1(\mu, r, t)$$

One can describe this fluctuating field by means of a power spectrum $P(k)$, where k is the wave number. If $\zeta$ is taken to be along the direction of the average magnetic field, one can now write a Fokker-Planck equation for $U(r, t)$:

$$\frac{\partial U}{\partial t} = D_\perp \frac{\partial^2 U}{\partial \chi^2} + D_\parallel \left[ \frac{\partial^2 U}{\partial x^2} + \frac{\partial^2 U}{\partial y^2} \right]$$

**Fig. 11.** This figure is similar to the last figure. For the electrons a variable $\lambda$ is used and a heat input at the upper boundary of $4 \times 10^9$ eV cm$^{-2}$ s$^{-1}$ is included. For the ions, $\lambda$ is infinity and there is no heat input. For the ions the solid and dashed lines indicate results with and without Joule heating, respectively.

**Fig. 10.** This figure is similar to Figure 9 except that the magnetic field is $10^7$ and the mean free path is 10 km, and different heat inputs are used.
where $D_i$, and $D_\perp$ are parallel and perpendicular diffusion coefficients, respectively. $D_i$ and $D_\perp$ are related to changes caused by the fluctuating part of the field by [Earl, 1974]

$$D_i = \frac{\rho}{2} \int_0^1 \frac{(1 - \mu^2)}{((\Delta \mu)^2/\Delta \mu)} \, d\mu$$

(A4)

$$D_\perp = \frac{1}{2} \int_0^1 \frac{(\Delta \mu)^2}{\Delta \mu} \, d\mu$$

(A5)

$\rho$ is the average thermal velocity of the particles (actually the velocity of the particles relative to the frame of reference of the fluctuations). $(\Delta \mu)^2/\Delta \mu$ and $((\Delta \mu)^2)/\Delta \mu$ can be derived from the power spectrum of the magnetic field. If the gyroradius $r_g$ is much less than the correlation length of the fluctuations $l$, then

$$\frac{((\Delta \mu)^2)}{\Delta \mu} = \frac{1 - \mu^2}{|\mu|} \frac{e^2}{m^2 c^2} P \left( k = 1/mr_g \right)$$

(A6)

$$\frac{((\Delta \mu)^2)}{\Delta \mu} = \frac{\rho e^2}{B_0} P(k = 0) + \frac{1 - \mu^2}{2|\mu|/B_0^2} P \left( k = 1/m\mu \right)$$

(A7)

where $e$ is the particle charge, $m$ is the particle mass, and $c$ the speed of light.

If the power spectrum in the ionosphere at all resembles the power spectrum in the solar wind, then the following approximation is useful:

$$P \approx P_0(k \cdot l)^{-q} \frac{1}{k < l}$$

$$P \approx P_0 \frac{1}{k \gg l}$$

(A8)

where $q$ is generally between 1 and 2. Using (A8) for the case when $r_g \ll l$, the diffusion coefficients become

$$D_i = \frac{\rho^2 m^2 e^2}{(2 - q)(4 - q) c^2 P_0 r_g}$$

(A9)

$$D_\perp = \frac{\rho}{B_0} P_0 \left( 1 + \frac{r_g^2}{N} \frac{q}{q + 2} \right)$$

(A10)

Consider perpendicular diffusion. This diffusion coefficient can be written in terms of an effective mean free path for vertical diffusion by $D_\perp = \frac{1}{2} \lambda$. Thus for $r_g \ll l$ and using (A10), $\lambda$ can be expressed as

$$\lambda \approx \frac{2}{B_0^2 P_0}$$

(A11)

At the moment very little is known about the exact nature of the power spectrum in the ionosphere of Venus. R. Elphic (private communication, 1980) indicates that $l \approx 10$ km and $B_i/B_0$ is about 1. A crude approximation for $P_0$ is

$$P_0 \approx \frac{B_i^2}{B_0} \cdot B_0^2 \cdot l$$

(A12)

Using (A12) in (A11) one finds that the vertical mean free path predicted by quasi-linear theory is

$$\lambda \approx 2 \cdot \frac{B_i^2}{B_0} \cdot l$$

(A13)

If $B_i/B_0 = 1$ (the applicability of quasi-linear theory thus being borderline) and $l = 3$ km, then $\lambda = 6$ km. For $B_i/B_0 = 1$ and $l = 3$ km, $\lambda = 1.5$ km.

When the fluctuations are very large and the gyroradius is very large, then it is clear that $\lambda$ will be relatively large; however, quasi-linear theory will not be applicable in this case.

Acknowledgments. The authors wish to thank A. Hedin and H. B. Niemann for providing us with neutral gas composition data prior to publication and H. A. Taylor, Jr., for providing us with ion composition data prior to publication. We also want to acknowledge useful discussions with R. Elphic and C. T. Russell. The work described in this paper was supported by the National Aeronautics and Space Administration through contract NASA-9130, grant NGR23-005-015, grant NAGW-15, and NASA/GSFC contract NAS-25939. The major fraction of the computations were carried out using the computing facilities of the National Center for Atmospheric Research, which is sponsored by the National Science Foundation.

The Editor thanks D. F. Strobel and W. E. Swartz for their assistance in evaluating this paper.

REFERENCES


(Received December 20, 1979; revised April 15, 1980; accepted April 16, 1980.)