A New Model of Cometary Ionospheres

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The coupled continuity, momentum, and energy equations were solved for ionospheric conditions appropriate for comet Halley at 1 AU. The numerical scheme used is such that any shock transition appears naturally in the solution and no a priori assumptions are necessary. Solutions were obtained for a number of different assumptions concerning electron heating rates, but all showed that the electron temperatures increase rapidly and significantly at a distance from the nucleus where collisional electron-neutral cooling becomes unimportant. This temperature increase is accompanied by a sharp increase in both the plasma pressure and its associated polarization electric field, causing the supersonic plasma flow to go subsonic. It is not clear at this time whether or not this sonic transition is accompanied by a shock.

1. INTRODUCTION

A number of papers dealing with the chemistry, energetics, and dynamics of cometary ionospheres have been published during the last decade. Most of the early models (for references, see Mendis et al. [1985]) concentrated on the chemistry of the atmosphere and ionosphere and used highly simplified assumptions of important parameters such as species velocity and temperature. Marconi and Mendis [1982] solved the coupled hydrodynamic equations for the dominant neutral gas along with a limited set of chemical reactions. At about this same time it was recognized that neutral hydrogen easily decouples from the heavier species (e.g., H2O and CO2), and this led to a number of multifluid models [Marconi and Mendis, 1982; Huebner and Keady, 1983]. It was also realized by Marconi and Mendis [1983] that at distances greater than about 10^3 km, electrons may decouple energetically from the heavier ions and neutrals and significantly elevated electron temperatures may result. Cravens et al. [1984] used the Marconi and Mendis [1983] results to establish the fact that the rapidly increasing electron temperatures beyond about 10^3 km imply a large pressure gradient (and related ambipolar electric field), which in turn is likely to decelerate the supersonically flowing plasma to subsonic velocities by means of a shock. Kőrösmézy [1984] examined this problem in more detail by solving the coupled steady state hydrodynamic equations for the heavy neutrals, hydrogen, dust, ions, and electrons. These calculations did establish that the ion and neutral temperatures decouple even in the collision-dominated inner coma; it also indicated that an inner shock is likely to be present, but the steady state nature of the model did not allow any definitive conclusions to be drawn concerning this question.

The purpose of the calculations presented in this paper is to study the dynamics and the energetics of the ionosphere in more detail than previous models. The calculations presented in this work were obtained by solving the dimensionless, coupled, time-dependent continuity, momentum, and energy equations of the ion and electron gas. In the collision-dominated region the charged particles are believed to have very little influence on the neutral gas and dust behavior; therefore in the present work the neutral gas parameters necessary to solve the ion and electron equations are obtained from the neutral gas/dust model (Table 1) of Gombosi et al. [1986] and are shown in Figure 1. The total gas production rate for this model is Q = 5.1 x 10^29 s^-1, appropriate for comet Halley at about 1 AU.

Section 2.1 describes the actual equations being solved, and the necessary input parameters are described in section 2.3. A brief discussion of the numerical scheme is given in section 2.2. The results of the calculations and the implications of these new results are discussed in sections 3 and 4, respectively.

2. DESCRIPTION OF THE MODEL

2.1. Controlling Equations

The model discussed in this paper assumes a pure water ice comet for the sake of simplicity; therefore the only ions considered were H2O+ and H2O-. When calculating the energy source terms we also consider reactions producing molecules or ions other than H2O, H2O+, or H2O-, but we do not keep track of the minor reaction products. Furthermore, spherical symmetry is assumed in the calculations. Denoting the number density of H2O+ as n_i, where n_i is the total ion density, the continuity equations for the total ion and H2O+ densities can be written as [Schunk and Nagy, 1980; Gombosi et al., 1985]

\[ \frac{\partial n_i}{\partial t} + \frac{1}{r^2} \frac{\partial}{\partial r} (r^2 n_i u_i) = S_2 + S_3 \]  

\[ \frac{\partial c}{\partial t} + u_i \frac{\partial c}{\partial r} = \frac{1}{n_i} [ (1 - c) S_2 - c S_3 ] \]  

where S_2 and S_3 are the net source terms for H2O+ and H2O-, respectively, and u_i and c are the ion and electron velocities, respectively. The momentum equations for these ions and electrons can be written as [cf. Schunk and Nagy, 1980; Gombosi et al., 1986]

\[ m_i \frac{\partial}{\partial t} (n_i u_i) + \frac{m_i}{r^2} \frac{\partial}{\partial r} (r^2 n_i u_i^2) + \frac{\partial p_i}{\partial r} = F_{\text{in}} + e n_i E + m_i u_e S_{\text{chem}} - m_i u_i S_{\text{rec}} \]  

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Fig. 1. Total neutral density as a function of cometocentric distance, based on the model of Gombosi et al. [1986].

\[
m_e \frac{\partial}{\partial t} \left( n_e u_e \right) + m_e \frac{\partial}{\partial r} \left( r^2 n_e u_e^2 \right) + \frac{\partial p_e}{\partial r} = F_{en} - en_e E + m_e u_e S_{chem} - m_e u_e S_{rec} \tag{4}
\]

where \( p_i \) and \( p_e \) are the ion and electron pressures, respectively; \( u_i \) is the neutral gas velocity; \( m_i \) and \( m_e \) are the ion and electron masses, respectively; and \( S_{chem} \) and \( S_{rec} \) are the total ion number density source and sink terms due to photoionization and recombination, respectively. Note that \( n_e = n_i \) and that \( u_i = u_e \) in order to keep the current equal to zero. Combining the two momentum equations, one obtains the following expression for the polarization electric field:

\[
E = \frac{en_e}{en (m_i + m_e)} \left[ \frac{\partial}{\partial r} \left( \frac{m_i}{m_e} p_i - \frac{m_e}{m_i} p_e \right) + \frac{m_i}{m_e} F_{en} \right] \tag{5}
\]

It is reasonable to assume that in the region of interest for this work the collision frequency among the different ions is sufficiently large to lead to a single ion temperature. The corresponding ion and electron energy equations can be written as [cf. Schunk and Nagy, 1980; Gombosi et al., 1986]

\[
m_i \frac{\partial}{\partial t} \left( \frac{n_i u_i^2}{2} + \frac{n_i}{\gamma_i - 1} k_{Ti} \right) + \frac{m_i}{r^2} \frac{\partial}{\partial r} \left( r^2 n_i u_i^3 \right) + r^2 \frac{\gamma_i}{\gamma_i - 1} \frac{n_i u_i k_{Ti}}{2} = Q_i + Q_{in} + Q_{rec} + r^2 \frac{\gamma_e}{\gamma_e - 1} \frac{n_e u_i k_{Te}}{2} = Q_e + Q_{in} + Q_{rec} \tag{6}
\]

where \( \gamma_i \) and \( \gamma_e \) are the ratios of specific heats; \( Q_i \) is the ion heating rate due to chemical processes; \( Q_e \) is the electron heating rate due to photoelectrons; \( Q_{in} \) and \( Q_{rec} \) are the elastic energy exchange rates between the ions and neutrals and ions and electrons, respectively; and \( L_{ion} \) and \( L_{vib} \) are rotational and vibrational cooling rates for electron collisions with \( \text{H}_2\text{O} \). Finally, \( \kappa_i \) and \( \kappa_e \) are the ion and electron thermal conductivities, respectively.

2.2. Numerical Method

The numerical code solves the system of equations, consisting of the two continuity equations (1) and (2), the ion momentum equation (3), and the two energy equations (6) and (7), for the five unknowns \( n_i, c, u_i, T_i \), and \( T_e \). It was assumed that \( n_i = n_e \) and that \( u_i = u_e \) in order to keep the current equal to zero.

The numerical method used to solve the equations was a modified form of the first-order scheme of Godunov [1959]. The acoustic approximation formulae of Godunov were used except that instead of the sound velocity being used (as in Godunov's original scheme) the ion acoustic velocity is employed:

\[
v_i = \left( k \frac{\gamma_i T_i + \gamma_e T_e}{m_i + m_e} \right)^{1/2} \tag{8}
\]

Another modification to the original scheme is that a Crank-Nicholson type method [cf. Gerald, 1978] was incorporated to solve the second-order equations at each time step; this is required because thermal heat conduction is included, unlike the original Godunov scheme.

2.3. Model Inputs and Parameters

Neutral density. The total neutral gas production rate and gas terminal velocity adopted for this model [from Gombosi et al., 1986] are \( Q_n = 5.1 \times 10^{29} \text{ s}^{-1} \) and \( u_n = 6.6 \times 10^4 \text{ cm s}^{-1} \), giving a ratio of \( Q_n/u_n = 0.77 \times 10^{25} \text{ cm}^{-1} \). The parent molecules were assumed to be all \( \text{H}_2\text{O} \). The neutral gas number densities were obtained from the model of Gombosi et al. [1986], adopted for this work, and are shown in Figure 1. The neutral mass spectrometer carried aboard the Giotto spacecraft measured a total neutral gas density which was within a few percent of the value assumed in this calculation, and it also found that the water vapor abundance was about 80% of the total [Krankowsky et al., 1986]. The mass spectrometer also established that the ratio \( Q_n/u_n \approx 6.9 \times 10^{30} \text{ s}^{-1} \cdot 0.9 \times 10^4 \text{ cm}^{-1} = -0.76 \times 10^{29} \text{ cm}^{-1} \). The neutral gas sensor of the VEGA Plasmag-1 instrument package found a somewhat larger total neutral gas production rate of \( \sim 1.3 \times 10^{30} \text{ s}^{-1} \) [Gringauz et al., 1986].
Photoionization and chemistry. As indicated in section 2.1, only a very simple chemical scheme was adopted for this model, because the emphasis of these calculations is on the energetics and dynamics of the ionosphere. The model includes the following reactions:

\[
\begin{align*}
\text{H}_2\text{O} + h\nu &\rightarrow \text{H}_2\text{O}^+ + e^- \quad (9) \\
\text{H}_2\text{O} + h\nu &\rightarrow \text{H}^+ + \text{OH} + e^- \quad (10) \\
\text{H}_2\text{O} + h\nu &\rightarrow \text{OH}^+ + \text{H} + e^- \quad (11) \\
\text{H}_2\text{O}^+ + \text{H}_2\text{O} &\rightarrow \text{H}_3\text{O}^+ + \text{OH} + 1.9 \text{ eV} \quad (12) \\
\text{H}_2\text{O}^+ + e^- &\rightarrow \text{OH} + \text{H} + 7.5 \text{ eV} \quad (13) \\
\text{H}_2\text{O}^+ + e^- &\rightarrow \text{O} + \text{H} + 7.5 \text{ eV} \quad (14) \\
\text{H}_2\text{O}^+ + e^- &\rightarrow \text{OH} + \text{H} + 6.4 \text{ eV} \quad (15) \\
\text{H}_2\text{O}^+ + e^- &\rightarrow \text{OH} + \text{H} + 1.2 \text{ eV} \quad (16) \\
\text{H}_2\text{O}^+ + e^- &\rightarrow \text{H}_2\text{O} + 6.4 \text{ eV} \quad (17)
\end{align*}
\]

The dissociative recombination reaction rates (13)-(17) and the rearrangement rate (12) were taken from the tabulations of Mendis et al. [1985]. The photoionization rates (9)-(11) were calculated using the ionization cross-section values of Haddad and Samson [1986] and the solar flux values measured by Hinteregger et al. [1973] corresponding to solar cycle minimum and maximum conditions [cf. Gombosi et al., 1986]. The calculations were carried out using 37 wavelength intervals. The ionization frequencies \( f_i \) thus calculated (for a distance of 4650 km) are given in Table 1, for both solar minimum and maximum conditions and for a heliocentric distance of 1 AU, but only the former are used in the model calculations described in this paper. Note that the VEGA and Giotto Halley encounters were at a heliocentric distance of about 0.8–0.9 AU and not 1 AU. The ions in the ionosphere are primarily heated by the excess energy carried off by the ions during some of the photochemical reactions just described (this heating rate is designated \( Q_i \)). Most of the heating comes from (12), with very minor contributions from (9)-(11). Coulomb collisions with electrons are a relatively minor source of heat, which is even negative (cooling) as small values of \( r \).

Photoelectrons. In order to calculate the secondary ionization and electron heating rates due to photoelectrons, we calculated steady state photoelectron fluxes using a modified version of the two-stream approximation, first introduced for the terrestrial ionosphere by Nagy and Banks [1970]. The cross sections for electronic excitation and dissociation and dissociative attachment were taken from Olivero et al. [1972], for rotational excitation from Cravens and Korosmezey [1986], and for vibrational excitation from Seng and Linder [1976], and for elastic scattering from Seng and Linder [cf. Gianturco and Thompson, 1980]. The backscatter probability for inelastic collisions was estimated from references given by Olivero et al. [1972] to be 0.35 for energies below 2 eV and 0.5 above 6 eV, with a linear variation between these values. The elastic backscatter probability was taken to be roughly 0.5 as estimated from the momentum transfer cross section of Seng and Linder [cf. Gianturco and Thompson, 1980]. A knowledge of the photoelectron fluxes also allows the secondary ionization rates to be calculated and included in the appropriate source terms.

In case of a completely magnetic field free ionosphere at comet Halley the collisional mean free path of a 10-eV photoelectron is less than the scale size (which is roughly the radial distance \( r \)) only at radial distances less than about 1000 km; this defines a photoelectron collision zone. Inside this collision zone the assumption of local energy deposition for the photoelectrons is appropriate. Outside this zone the photoelectrons move “freely,” and the spherical geometry must be taken into account. A Monte Carlo type model would be well suited to carry out such calculations, because nonradial transport is also possible. Ashihara [1978] did carry out such calculations for a pure \( \text{H}_2\text{O} \) comet with a production rate \( Q_n \) equal to \( 10^{30} \) \( \text{s}^{-1} \). The photoelectron fluxes calculated by Ashihara [1978] varied approximately as \( 1/r \), as one would expect for purely radial transport. In the present calculations a two-stream model was used, which only allows for radial transport, with an additional constraint of a constant cross-sectional area for the flux; this is equivalent to assuming the presence of a constant, uniform, and radial magnetic field. It will be shown later that the fluxes thus calculated do vary roughly as \( 1/r \) for electron energies greater than about 10 eV (the energy range where electron-electron collisions are negligible) and somewhat more steeply than \( 1/r \) at lower energies. A magnetic field free ionospheric cavity was found by the magnetometer carried aboard the Giotto spacecraft [Neubauer et al., 1986]. The upper bound on the x and y components of the field given by Neubauer et al. [1986] is 0.3 nT. However, the gyroradius of a 10-eV electron is \( \sim 100 \) km even in a magnetic field as small as 0.1 nT, which suggests that even an extremely small “remnant” field can significantly inhibit photoelectron transport. The photoelectron model calculations were carried out considering two limiting cases: (1) no photoelectron transport (local energy deposition) and (2) transport along a steady parallel field. We do not suggest that these photoelectron transport assumptions are correct, but the two limiting cases should indicate how sensitive the calculated fluxes and the heating rates are to the transport assumptions.

Collision terms. The momentum and energy transfer collisions and related processes (e.g., thermal conductivities) involving water molecules have to be considered in some detail.

<table>
<thead>
<tr>
<th>Solar Cycle Conditions</th>
<th>Photoelectron Transport</th>
<th>Direct Photoionization Rate</th>
<th>Secondary Impact Ionization Rate</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>( \text{H}_2\text{O}^+ )</td>
<td>( \text{OH}^+ )</td>
</tr>
<tr>
<td>minimum</td>
<td>no</td>
<td>3.19(-7)</td>
<td>5.48(-8)</td>
</tr>
<tr>
<td>minimum</td>
<td>yes</td>
<td>3.19(-7)</td>
<td>5.48(-8)</td>
</tr>
<tr>
<td>maximum</td>
<td>no</td>
<td>8.70(-7)</td>
<td>1.57(-7)</td>
</tr>
<tr>
<td>maximum</td>
<td>yes</td>
<td>8.70(-7)</td>
<td>1.57(-7)</td>
</tr>
</tbody>
</table>

The values in the parentheses should be read as \((-8) = \times 10^{-8}\). Frequencies are in \( \text{s}^{-1} \).
because H$_2$O is a polar molecule. The H$_2$O$^+$-H$_2$O collision frequency can be written, as shown recently by Cravens and Körömezey [1986], as

$$v_{in} = 1.3 \times 10^3 \left[ 1 + \frac{1}{2(\pi T_u)^{1/2}} + \frac{9(1 + 2T_u)}{8(\pi T_u)^{1/2}} \right] \left( \frac{x_0}{\mu_n} \right)^{1/2} n_n$$

where

$$T_u = 0.7T_e + 0.3T_i$$

and $T_e$ and $T_i$ are the neutral and ion temperatures, respectively; $x_0$ is the polarizability of the neutral H$_2$O molecule in cgs units; and $\mu_n$ is the reduced mass in atomic mass units.

The polarizability of the neutral H$_2$O molecule is taken to be $1.444 \times 10^{-24}$ cm$^3$ [Eisberger and Kautzmann, 1969]. Note that there are two typographical errors in equation (20) of Cravens and Körömezey [1986]; the expression in brackets should be divided by $2n^{1/2}$ and exp $(T_e^{-1})$ should be exp $(-T_e^{-1})$. Once the ion-neutral collision frequency is known, the momentum and energy exchange rates can be easily obtained by using the following relations [cf. Banks and Kockarts, 1973]:

$$F_{en} = m_n v_{en} u_n (u_n - u_e)$$

$$Q_{en} = \frac{2m_n^2}{(m_i + m_n)^2} v_{en} u_n [\frac{\nu}{3} k(T_n - T_i) + \frac{1}{2} (m_i u_i + m_n u_n) (u_n - u_i)]$$

The momentum transfer cross section for low-energy elastic collisions between electrons and H$_2$O molecules was measured by Pack et al. [1962]; using this cross section, the velocity-dependent collision frequency can be written as

$$v_{en}(\nu) = n_n u/\nu$$

where $a = 10.533$ in cgs units and $\nu$ is the relative velocity between the particles. Using this velocity-dependent collision frequency, the Maxwellian-averaged elastic electron collision frequency can be calculated [cf. Banks and Kockarts, 1973] and is

$$v_{en} = \frac{4a}{3} \left( \frac{m_e}{\pi kT_e} \right)^{1/2} n_n$$

The corresponding momentum and energy exchange rates, in terms of collision frequency, are [cf. Banks and Kockarts, 1973]

$$F_{en} = m_n v_{en} n_n (u_n - u_e)$$

$$Q_{en} = \frac{2m_n^2}{m_e} v_{en} n_n [\frac{\nu}{3} k(T_n - T_i) + \frac{1}{2} m_n u_n (u_n - u_e)]$$

The electron-ion energy exchange rates were calculated from the classical expression as given by Banks and Kockarts [1973] or Schunk and Nagy [1980]. Inelastic collisions between the thermal electrons and the H$_2$O molecules can result in rotational and vibrational excitation of the molecule; the rate for this interaction is especially large because of the polar nature of the H$_2$O molecule. The rates for these rotational and vibrational cooling processes were discussed in detail and presented in a recent paper by Cravens and Körömezey [1986]; the present model uses the cooling rates derived by these authors.

The expression for the electron thermal conductivity, including the effects of ion-neutral collisions, is [Banks and Kockarts, 1973]

$$\kappa_e = \left[ \frac{1}{\kappa_{el}} + \frac{1}{\kappa_{en}} \right]^{-1}$$

where $\kappa_{el}$ and $\kappa_{en}$ are the thermal conductivities corresponding only to charged particle and electron-ion collisions, respectively. The thermal conductivity for the fully ionized case, $\kappa_{el}$, is

$$\kappa_{el} = 1.23 \times 10^{-6} T_e^{5/2} \text{ erg cm}^{-1} \text{ s}^{-1} \text{ K}^{-1}$$

Assuming zero current and using the electron-neutral collision frequency as given by (21), $\kappa_{en}$ can be calculated from the
The integral relations given in the work by Banks and Kockarts [1973], which leads to the following expression:

$$\kappa_{\text{ia}} = 4.92 \frac{n_e}{n_i} T_e^{3/2} \text{erg cm}^{-1} \text{s}^{-1} \text{K}^{-1}$$  \hspace{1cm} (27)

A similar expression for $\kappa_{\text{ia}}$ was also obtained by Marconi and Mendis [1984]. Finally, combining (26) and (27), the following relation for the electron thermal conductivity is obtained:

$$\kappa_e = 1.23 \times 10^{-6} T_e^{5/2} \frac{1}{1 + 2.5 \times 10^{-7} T_e(n_e/n_i)} \text{erg cm}^{-1} \text{s}^{-1} \text{K}^{-1}$$  \hspace{1cm} (28)

For almost all values of $r$, Coulomb collisions dominate the conductivity for values of $T_e$ less than about 1000 K, and electron-neutral collisions dominate for $T_e > 1000$ K. The electron mean free path at distances less than about 10 km is much less than the value of $r$, which is the generally accepted scale length; therefore the concept of thermal conductivity is valid in this region. If there are significant magnetic field fluctuations, even in the case of very weak fields ($B < 0.1$ nT), the effective electron thermal conductivity may be smaller than the collisional value given by (26) [cf. Cravens et al., 1980].

The ion thermal conductivity $\kappa_i$ is obtained in a similar way by assuming zero current and using the ion-neutral collision frequency given by (18), including both ion-ion and ion-neutral collisions, and is

$$\kappa_i = \frac{9.48 \times 10^{-20} m_i^{-1/2} T_i^{5/2}}{1 + 5.21 \times 10^{-5} T_i(n_e/n_i)} \text{erg cm}^{-1} \text{s}^{-1} \text{K}^{-1}$$  \hspace{1cm} (29)

The value of $\kappa_i$ used to obtain (29) was taken from Banks and Kockarts [1973], while the expression for $\kappa_{\text{ia}}$ used was [cf. Banks and Kockarts, 1973]

$$\kappa_{\text{ia}} = \frac{4}{3} \frac{k}{n_i} \langle \nu \rangle$$  \hspace{1cm} (30)

where $\nu_0$ is given by (18) and $\langle \nu \rangle$ is a mean ion velocity.

3. RESULTS

3.1. Photoionization

The primary and secondary photoionization rates were calculated as outlined in section 2.3. The secondary ionization rates presented here assume no photoelectron transport (this case will be referred to as local). The calculated ionization frequencies for both solar maximum and minimum conditions are given in Table 1; as can be seen, secondary ionization makes a significant contribution to the total ionization rates, especially for $H^+$. The variations with radial distance of the calculated ionization rates for solar minimum conditions are shown in Figure 2. It should be pointed out that what is actually plotted in Figure 2 is $r^2 S$, resulting in a relatively flat variation with distance.

3.2. Photoelectron Fluxes and Electron Heating Rates

The steady state photoelectron fluxes were calculated using the two-stream method, as indicated in section 2.3, considering the two limiting cases of no transport (local energy deposition) and radial transport. These calculations require an a priori knowledge of the electron density profile. This was handled by using an iterative approach, in which an estimated
The calculated energy distributions of the photoelectron production rates for both solar cycle maximum and minimum conditions at a heliocentric distance of 1 AU and a radial distance of 4650 km are plotted in Figure 3. As expected, the main change with solar cycle is in magnitude and not the spectral shape. The energy spectra of the photoelectron fluxes depend both on the structure of the initial electron production rate, which in turn depends mainly on the solar EUV spectrum, and on the electron energy losses, which are determined by the electron impact cross sections and the cometary electron and neutral gas densities. The calculated photoelectron flux in the case of local energy deposition assumptions (no transport) is almost independent of $r$ for energies greater than about 10 eV, because both the production and loss rates are directly proportional to the neutral density; however, the flux does decrease with increasing $r$ at lower energies. The photoelectron fluxes calculated assuming no transport (local energy deposition) are shown in Figure 4 for two representative altitudes. For the case of no transport the upward flux ($\phi^+$) is the same as the downward directed one ($\phi^-$). The photoelectron fluxes calculated assuming radial transport (parallel magnetic field) and no inward flux at the upper boundary altitude of 10$^4$ km are shown in Figure 5 at the same two representative altitudes as for Figure 4. At low altitudes, below about 10$^3$ km, in the collision-dominated region, the upward and downward fluxes are still approximately equal, but at high altitudes the upward flux exceeds the downward directed flux by a significant factor. As discussed earlier, the total flux ($\phi^+ + \phi^-$), calculated with transport considerations included, varies roughly as $1/r$, as do the ones calculated by Ashihara [1978]. There is also a general agreement between the energy spectrum of the fluxes presented here and those calculated by Ashihara [1978]. The variation with radial distance of the...
photoelectron flux, for 20 eV, is shown in Figure 6. The calculated heating rates of the ionospheric thermal electrons by these photoelectrons are shown in Figure 7. The difference in these heating rates reaches a factor of about 2 at large distances.

There are indications in the electron flux data from ICE [Bame et al., 1986; Zwickl et al., 1986] and VEGA [Gringauz et al., 1986] that in the vicinity of comets, but outside the classical ionosphere [Cravens et al., 1987], a non-solar wind component is present at energies between about 10 eV and 100 eV. Zwickl et al. [1986] suggested that these electrons are photoelectrons; the overall shape of the spectrum presented here supports this suggestion.

3.3. Temperatures

The coupled continuity, momentum, and energy equations were solved as outlined in section 2.2. The results of three sets of calculations are discussed here, namely ones carried out for solar cycle minimum conditions (appropriate for the comet Halley encounters) using (1) electron heating rates calculated assuming radial photoelectron transport and no heat inflow at the upper boundary of 10⁴ km; (2) electron heating rates calculated assuming no photoelectron transport and no heat input at the upper boundary (a variation of this case was also tried, with no heat conduction); and (3) electron heating rates calculated assuming no photoelectron transport and a topside heat inflow of 3 x 10⁹ eV cm⁻² s⁻¹ (this heat inflow is presumably associated with the solar wind interaction by analogy with Venus [Cravens et al., 1979, 1980]; the value actually chosen is about 10% of that found at Venus [Scarf et al., 1979] and of the same order as the solar wind electron energy flux (nₑνₑkTₑ) at 1 AU).

Figure 8 shows the calculated electron and ion temperatures for these three cases considered (the neutral temperatures used for the calculations are also shown for comparison). The results obtained are as expected; smaller volume heating rates result in lower temperatures (i.e., $T_e$ for the case of photoelectron transport is somewhat less than $T_e$ for the local photoelectron case), and the additional topside electron heat inflow raises the electron temperatures at large distances. The electron temperatures follow closely the neutral gas temperature in the high-density region, where the electron-neutral cooling rate is large, and then rises rapidly in the 2000- to 4000-km region as the electron-neutral collision frequency decreases and as the neutral temperature also increases. This very large electron temperature gradient results in a corresponding increase in plasma pressure and associated polarization electric field, which in turn affects the dynamics of the plasma flow which will be discussed next. The electron temperatures reach about 1000 K at $r \sim 10^4$ km for no external heat input. It was found that heat conduction is not very important for this case, because the calculated $T_e$ only increased by about 200 K at the topside boundary when the conductivity was set equal to zero. A significant external (topside) heat input does produce a large gradient in $T_e$, and temperatures close to 10⁴ K are generated near $r = 10^4$ km. Heat conduction is the dominant physical process at large values of $r$ for this case. The ion temperatures for all three cases are always noticeably greater than the neutral temperatures at distances greater than 150 km. This indicates that heating of ions due to photochemistry is not entirely balanced by cooling due to ion-neutral collisions ($Q_{in}$).
3.4. Dynamics

The total ion density, plasma velocity, and Mach number profiles calculated for the three cases detailed in the previous section are shown in Figures 9, 10, and 11, respectively. The neutral gas velocity values from the model calculations of Gombosi et al. [1986], adopted as input parameters for our model, are also shown for comparison in Figure 10. The neutral and ion velocities are approximately the same for the smaller values of $r$ due to the large ion-neutral collision frequency. There are rapid drops in the calculated plasma velocities and the corresponding Mach numbers in the 2000- to 4000-km region where the rapid rise in the electron temperature is located with its corresponding rise in plasma pressure; the plasma pressure gradient decelerates the flow. There is an apparent weak shock transition ("inner shock") in the plasma velocities only for the external heat input case (case 3), while the transition through the sonic point is smooth for case 1 and 2. These drops in velocities are accompanied, as expected, with increases in the plasma densities as shown in Figure 9. The small ion temperature bump near 2500 km, visible in one of the profiles, is also associated with the rapidly changing velocities. The calculated velocities begin to increase sharply again beyond about 5000 km, initially because of outward ion-neutral drag, and later at larger $r$ because of the input of thermal energy, but the velocities remain subsonic (in relation, to the ion acoustic speed) because of the rapidly increasing temperatures.

It should be noted that this model does not include the effects associated with the ionopause/contact discontinuity and this means that the results are unreliable at large distances from the nucleus. The plasma number density is controlled about equally by dissociative recombination and by radial transport at almost all values of $r$. That is, the transport time is roughly comparable to the recombinatiol time. For photochemically controlled situations the plasma density varies approximately as

$$n_e \approx \frac{S}{\alpha_3} \approx \frac{Q_e R}{4n u_s (7 \times 10^{-7})/300} \frac{1/2}{r^{1/4}} T_e^{1/4},$$

where $\alpha_3$ is the dissociative recombination coefficient for $H_2O^+$ (see (15)-(17)); $\alpha_3 = 7 \times 10^{-7} (T_e/300)^{-0.5}$ [Mendis et al., 1985]. Therefore the electron density, in the photochemically controlled region, varies as $1/r$. Near the "shock" where the plasma velocity drops, $n_e$ jumps because of both the sharp increase of $T_e$ and the decrease in $u_e$. At large $r$, $n_e$ decreases faster than $1/r$, because of the rapid increase of $u_e$.

4. Discussion

Marconi and Mendis [1986] recently published some results of their model calculations of electron temperatures and densities. They did not solve the ion and electron momentum equations but assumed that the plasma velocity is the same as the neutral gas velocity, and they made some simplifying assumptions about the electron heating rate profile and assumed a gas production rate of $4.0 \times 10^{28} \text{s}^{-1}$, appropriate for comet Giacobini-Zinner. Because of this latter assumption, detailed comparisons are not appropriate, but the predicted rapid rise in the calculated temperatures, in the region where electron-neutral collisions become negligible, indicates a qualitative agreement between the two models. However, their calculated electron temperatures at very large $r$ ($T_e \sim 10^5 \text{K}$) are much larger than those obtained by this model ($T_e \sim 10^3 \text{K}$) because of the probable differences in the electron heating rates. The calculations presented here have established that a large portion of the photoelectron energy is lost to the neutral gas via inelastic collisions rather than to the ambient thermal electrons.

The ion mass spectrometer experiment [Balsiger et al., 1986] and the neutral mass spectrometer operating in its ion mode [Krankowsky et al., 1986], carried by the Giotto spacecraft, measured a number of ionospheric parameters at comet Halley. The ions were observed to consist mostly of $H_2O^+$ with some $H_3O^+$, $OH^+$, and other minor ions present. The total ion density (or more accurately the ion current counting rate) was found to vary approximately as $1/r$ in the iono-
sphere. The data of Balsiger et al. [1986] also show a region at around a few thousand kilometers where the ion counting rate deviates from the 1/r behavior, possibly suggesting the existence of an inner shock; however, this is far from being firm evidence for the existence of such a shock. The ion temperature measured by the Giotto ion mass spectrometer [Balsiger et al., 1986] increases from a few hundred degrees Kelvin near 2000 km to over a thousand degrees Kelvin at the contact discontinuity (~4500 km); the model temperature values presented in this paper increase from over 100 K to just below 1000 K over the same distance. The lack of a large jump in the measured \( n_e \) profile for \( r < 4500 \) km suggests that the gradient of \( T_e \) cannot be extremely large over this same region; our model without heat inflow gives the least amount of structure in the \( n_e \) profile. Therefore, on the basis of this indirect evidence from the Giotto ion density and temperature measurements, the model with no heat inflow is the most consistent in the sense of profile. Therefore, on the basis of this indirect evidence from the Giotto ion density and temperature measurements, the model with no heat inflow is the most consistent with the observations, but there is no information available at this time to choose among the "no heat inflow" models. Large enhancements of \( n_e \) were observed [Balsiger et al., 1986] just outside the contact discontinuity and appear similar to our calculated jump in \( n_e \) when heat input is included, which is closely connected to enhancements in \( T_e \). It is unclear at this point what the connection is, if any, between the modeled and this observed density jump.

The very limited experimental data available at this time are not inconsistent with the basic features of the theoretical results presented in this paper, although the existence of an inner shock has not been established. Any meaningful study of the behavior of the ionosphere in the immediate vicinity of the contact discontinuity will require the development of a more sophisticated model that the one presented here. It is interesting to note that the ion counting rates [Balsiger et al., 1986] continue to vary approximately as 1/r right through the contact surface. This 1/r variation, along with the observed low plasma velocities [Balsiger et al., 1986], suggests that the plasma just outside the contact surface is still photochemically controlled. This outside plasma appears to be ionospheric, certainly by planetary standards (e.g., Venus), even though it is magnetized [Cravens, 1986; Cravens et al., 1987]. This points to the existence of two regions of cometary ionospheric plasma, an inner unmagnetized ionosphere, to which the current theoretical model applies, and an outer magnetized ionosphere. The data available clearly indicate that the internal plasma pressure cannot balance the sum of the external plasma pressure and the external magnetic pressure across the contact surface. In fact, no clear drop in ion density (that is, an ionopause) was observed across the contact surface in a manner analogous to Venus; therefore ion-neutral drag must play a dominant role in the momentum balance of this region [Cravens, 1986; Ip and Axford, 1987].

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