A Numerical Model for the Dynamics and Composition of the Venusian Thermosphere

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ABSTRACT

The structure, composition and winds of the mesosphere and thermosphere of Venus are investigated using a nonlinear time-dependent hydrodynamic model. The assumption that all variables depend only on altitude and distance from subsolar point allows a two-dimensional formulation of the problem. Within this framework the model provides an entirely self-consistent treatment of the multi-component fluid. The system solved consists of four time-dependent equations for motion, temperature, and the distributions of O and CO, and two diagnostic equations representing continuity and hydrostatic balance for the total fluid. The model is forced by absorption of solar radiation which provides heating of molecules and dissociation of CO₂ into CO and O. A large-scale circulation is calculated, the gross features of which resemble those derived in an earlier simplified model, consisting of a single cell with rising motion on the dayside, sinking motion on the nightside, and a day-to-night horizontal flow. This circulation in the global mean acts to remove the light gases to balance the photodissociation. Relatively large concentrations of light gases build up on the nightside. The consequent increase of the pressure at a given level acts to block the nightward circulation. Hence little motion occurs in a large region centered around the antisyol point. Instead, most of the downward vertical motion occurs within an internal boundary layer just to the nightside of the terminator. Exospheric temperatures predicted by the model, using an EUV heating efficiency of 0.30, range from greater than 600 K at the subsolar point to less than 300 K at the antisyol point, whereas there are typically 20-40 K horizontal variations of temperature in the mesosphere. Maximum horizontal velocities are order of 300 m s⁻¹ and occur on the dayside near the terminator at the level of the exobase. The model predicts that CO and O will have relative number densities of 4% on the dayside at the level of the F₁ ionospheric peak, provided vertical eddy mixing is negligible.

1. Introduction

A planetary thermosphere depends for its very existence on fluxes of extreme ultraviolet (EUV) solar radiation to provide it with a distribution of heating and molecular photodissociation. These fluxes are of primary importance for the determination of global mean profiles of temperature and composition. One-dimensional models with vertical transfer of heat and constituents by eddy mixing have generally provided the first preliminary derivations of these profiles, often giving results in remarkable agreement with available observations. Such models have not been developed with a quantitative understanding of the motions responsible for transport, but in reproducing observations they show what time scale of transport must be provided by whatever motions are present. Except for this estimate of time scales, the one-dimensional modelers generally make no statement as to the nature of the motions, either as to their spatial scale or as to whether they are in reality turbulent or organized.

The usefulness of one-dimensional models is thus limited since they do not use motions calculated from first principles. Furthermore, they are incapable of providing information as to the horizontal distribution of temperature and composition. Their application to the determination of local values, using local fluxes of solar radiation, has generally been of doubtful applicability because they neglect the redistribution of temperature and composition by motions. The logical basis, then, for a calculation of thermospheric structure in more than one dimension is a hydrodynamic model, that is, a model which derives from first principles the motions of a planetary thermosphere.

It would appear from the theoretical viewpoint that the simplest thermosphere to model hydrodynamically is that of Venus. Time scales for the most important global dynamic processes in the Venusian thermosphere are much less than one earth day, whereas the sidereal and solar rotation periods of the planet are, respectively, 243 and 117 days. Thus it seems appropriate to assume that the dynamic and thermal structure of the Venusian thermosphere depends only on the solar zenith angle, that is, the angular distance from the direction of the sun.

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Several important questions have been raised concerning the thermosphere of Venus in the light of present observational and theoretical studies which can only be properly treated theoretically with a dynamic model. First, we would like to know what are the fields of large-scale horizontal and vertical velocities, perhaps the most basic variables of a dynamic model. A related question is what is the variation of temperature from dayside to nightside. A previous model (Dickinson, 1971) derived wind and temperature variation for the Venusian thermosphere. That calculation assumed that nonlinear hydrodynamic terms could be neglected and that the thermospheric composition could be considered pure carbon dioxide. It was, however, difficult—at least at some levels—to justify the neglect of the nonlinear terms because of the calculated velocities of several hundred meters per second from dayside to nightside and the accompanying strong vertical shears.

The nature of errors in calculating motions resulting from the assumption of pure carbon dioxide composition depends on the distribution of constituents that is consistent with the large-scale circulation of the dynamic model. It is, indeed, important to establish whether the circulation defined by the model will act to remove from the Venusian thermosphere the CO and O produced by photodissociation of CO₂. To examine this question we previously developed a model for the balance in the upper atmosphere of Venus between photoproduction of CO and O and transport by diffusive and hydrodynamic terms (Dickinson and Ridley, 1972). As a first attempt at defining the role of large-scale circulation in determining the composition of the Venusian thermosphere, we integrated this model with winds defined according to the Dickinson (1971) model. That calculation gave a dayside concentration of O and CO at the level of the F-1 peak ionosphere of approximately 2%, consistent with existing observational constraints (e.g., as discussed by Donahue, 1971).

The conclusion that thermospheric motions could transport downward the relatively light gases generated by photodissociation depends on a nightside increase of the concentrations of O and CO so that downward vertical transfer will occur by a negative correlation between vertical motion and light gas mixing ratios. Indeed, the calculation showed the CO and especially O concentrations would attain sufficiently large values on the nightside to significantly elevate constant pressure surfaces and, indeed, to reverse their gradient relative to its value calculated assuming pure CO₂. The resulting pressure force directed toward the subsolar point would tend to reverse the direction of the circulation from that derived assuming pure CO₂ composition, so that it appeared necessary to determine jointly the composition and circulation in order to satisfactorily determine either set of parameters.

Hence, the primary objective of the present paper is to provide a model which calculates self-consistently the dynamics and composition of the Venusian thermosphere. The various constituents are redistributed by the motions and their concentrations in turn define in part the horizontal pressure gradients which drive the winds. Since this coupling is basically nonlinear, it is most convenient to derive the steady-state solution by means of a time-dependent calculation. Such an approach also allows inclusion of the previously omitted nonlinear hydrodynamic terms without undue further complication. We have carried out these numerical calculations with the finite-difference model described in the next two sections. The results of the model, as described in Section 4, are similar on the dayside to
those previously reported by Dickinson (1971) and Dickinson and Ridley (1972). Some very interesting and significant differences are, however, found for the nightside. The nightside bulge in O and CO concentrations is again obtained, but without the previously determined very large concentrations of atomic oxygen at the antisol point. The reversal of the pressure gradient by the gases has a major effect on both horizontal and vertical velocities, as is later discussed.

2. Formulation

The intent of this section is to present the model equations being integrated and the physical processes being considered. The model assumes the composition to consist of CO$_2$ which is photodissociated by solar radiation into CO and O. The solar radiation also provides a spatially varying heat source. The heat source drives a circulation system that acts, along with vertical molecular diffusion, to transport mass, heat and momentum. At a fixed level in the vertical, all variables are assumed to depend only on the solar zenith angle $\theta$, as indicated in Fig. 1. The vertical coordinate used is defined to be log pressure measured from the $p_0 = 5 \times 10^{-3}$ dyn cm$^{-2}$ (i.e., $\mu b$) pressure level, i.e.,

$$z = \log p(p_0/p),$$

where $p$ is total pressure. This is the same definition used in our previous models. The appropriate hydrostatic equation then is

$$\frac{\partial \Phi}{\partial z} = R^* T/m,$$  \hspace{1cm} (2)

where $\Phi$ is the potential of a constant pressure surface, $T$ temperature, $R^*$ the universal gas constant, and $m$ the number-density-weighted mean molecular mass. Altitude $h$ is related differentially to $\Phi$ according to

$$dh = d\Phi/g(h),$$ \hspace{1cm} (3)

where $g$ is the value of gravity. Horizontal velocities $u$ are determined from the equation of motion,

$$\frac{1}{r \sin \theta} \frac{\partial}{\partial \theta} \left[ \frac{\mu}{\partial \Phi} + \frac{\partial u}{\partial \theta} \right] + \frac{\partial v}{\partial z} = \frac{1}{\rho_0} \frac{\partial \Phi}{\partial z} = \frac{1}{\rho_0} \frac{\partial \Phi}{\partial z} - \frac{1}{r} \frac{\partial}{\partial \theta},$$  \hspace{1cm} (4)

where $t$ is time, $r$ the distance from center of planet approximated by 6200 km, $w = dz/dt$ is the vertical velocity parameter, $\mu$ the viscosity coefficient, and $H = R^* T/m g$. The square brackets contain the nonlinear terms previously neglected by Dickinson (1971). The vertical and horizontal velocities are related by the continuity equation, written

$$\frac{1}{r \sin \theta} \frac{\partial (u \sin \theta)}{\partial \theta} + \frac{\partial}{\partial z} (e^{-2w}) = 0.$$  \hspace{1cm} (5)

The equation for determining CO and O concentrations was derived in Dickinson and Ridley (1972) and is written

$$\frac{\partial}{\partial t} \psi = -e^{-2w} \frac{\partial}{\partial z} m_{\text{CO}} (T_{\text{eq}}) \alpha \psi^2 - \frac{u}{r} \frac{\partial}{\partial \theta} \psi - \frac{w}{\partial z} \psi + S,$$  \hspace{1cm} (6)

where $\psi = (\psi_{\text{CO}}, \psi_{\text{O}})^t$ is the mass mixing ratio for CO and O written as a vector; $\tau$ the diffusion time scale $\tau = (p_0 \rho_0 D_0)/p_0 D_0 = 9.25 \times 10^2$ s, where $p_0 = 10^4$ dyn cm$^{-2}$; $D_0 = 0.2$ cm$^2$ s$^{-1}$ is a characteristic diffusion coefficient at pressure $p_0$, and temperature $T_{\text{eq}} = 273$ K; and

$$S = (J/m_{\text{CO}}) (1 - \psi_{\text{CO}} - \psi_{\text{O}}) \left( \frac{m_{\text{CO}}}{m_{\text{O}}} \right)$$

is the effective mass source per unit mass derived from rate $J$ of CO$_2$ photodissociation and the molecular masses for CO$_2$, CO and O denoted $m_{\text{CO}_2}$, $m_{\text{CO}}$, and $m_{\text{O}}$, respectively. The matrices $\alpha$ and $L$ are defined in the following discussion. The matrix $\alpha$ varies as the inverse of diffusion coefficients, according to

$$\alpha_{ij} = -\left( \frac{\phi_{ij}}{\phi_{ij} + (\phi_{ij} - \phi_{ij}) \psi^2} \right),$$  \hspace{1cm} (7)

where using $i = 1$ denotes CO, $i = 2$ denotes O, $i = 3$ denotes CO$_2$, and

$$\psi_{ij} = \frac{D_{ij} m_j}{D_i}$$  \hspace{1cm} (8)

with the $D_{ij}$ as the mutual diffusion coefficients for gases $i$ and $j$, using the values at STP given in Table 1 and the accompanying discussion of Dickinson and Ridley (1972), and with $D$ a characteristic diffusion coefficient with the same temperature and pressure dependence as the $D_{ij}$, i.e.,

$$D = [D_0 \left( \frac{T}{T_{\text{eq}}} \right)^{1.75}.$$

The matrix operator $L$ is diagonal, with elements

$$L_{ij} = \delta_{ij} \left[ -\epsilon_{ii} \right],$$  \hspace{1cm} (10)

where

$$\epsilon_{ii} = \frac{m_i}{m_i} - \frac{1}{m_i} \frac{\partial m_i}{\partial z}.$$  \hspace{1cm} (11)

The matrix $L$ defines diffusive equilibrium solutions through $L \psi = 0$. Departures from diffusive equilibrium are driven by the hydrodynamic transport terms in
square brackets and by the photodissociation source term $S$.

To calculate temperature, we solve the thermodynamic equation, written

$$
\frac{\partial T}{\partial t} = - \frac{k}{\rho C_p} \frac{\partial T}{\partial z} - a(T - \bar{T})
$$

$$
- \left[ \frac{\mu}{r} \frac{\partial T}{\partial \theta} + \frac{\partial^2 T}{\partial \xi^2} + \frac{R^* T^2}{C_{p,\theta}} \right] + \frac{Q}{C_p},
$$

(12)

where $k$ is the coefficient of thermal conductivity, $\bar{T}$ is the global mean equilibrium temperature calculated from the detailed model for nonlocal thermodynamic equilibrium (NLTE) infrared processes described in Dickinson (1972), $a$ is a radiative damping coefficient depending only on $z$ (also calculated from that computer program), $Q$ is the rate of heat addition per unit mass, and $C_p$ is the specific heat at constant pressure. The contribution of solar infrared heating and 15 $\mu$m CO$_2$ emission cooling to $Q$ is calculated at $T = \bar{T}$ according to Dickinson (1973) and depends only on solar zenith angle and pressure level. At levels below $z = -1$, the infrared solar heating and thermal cooling dominate and essentially cancel in the global mean, providing a global mean $Q$ that is close to zero. A small global mean cooling in the mesosphere results from slight differences between the mean 15 $\mu$m cooling from Dickinson (1973) and that calculated according to Dickinson (1972) which was used to determine $\bar{T}$. The greatest differences arise in the regions above $z = -5$ where NLTE 15 $\mu$m cooling is coupled to zenith-dependent solar infrared fluxes in a nonlinear fashion. Dickinson (1972) used the mean solar zenith for specifying solar fluxes whereas Dickinson (1973) calculated the NLTE 15 $\mu$m cooling at different values of solar zenith.

The definition of specific heats and thermal conductivities for CO$_2$ and CO is not entirely straightforward, since the populations of their vibrational energy levels are not in local thermodynamic equilibrium with molecular energy translation at the pressures of our model. At pressures at and above $p_0$, the internal energy carried by vibrations is negligible compared to the translational energy. We assume this to be the case at all levels so that specific heat per mole of CO$_2$ and CO is simply $(7/2)R^*$. The mean specific heat $C_p$ is given by the mass-weighted average of the specific heat per gram of the individual components. For thermal conductivities of individual molecular constituents, we use the definition

$$
k = \mu C_p / Pr,
$$

where $Pr$ is the Prandtl number, and the viscosities $\mu$ for CO$_2$ and CO are obtained from the tables of Hilsenrath et al. (1960). We use the expressions of Dalgarno and Smith (1962) for atomic oxygen viscosity and conductivity and Eq. (14.32) of Banks and Kockarts (1973) to define viscosity and conductivity averaged over the constituents.

Finally, we must specify the rate of heat input, $Q$, and photodissociation per CO$_2$ molecule, $J$. The solar radiation absorbed by CO$_2$ is calculated as was done by Dickinson and Ridley (1972) except that we use the cross sections of Shemansky (1972) and his suggested scattering corrections for wavelengths $> 1900$ Å. Cross sections for CO and O in the EUV are taken from Sun and Weissler (1955) and Dalgarno et al. (1964), respectively. We neglect thermal emission in the 4.7 $\mu$m vibration-rotation system of CO and the 63 $\mu$m lines of O. Furthermore, no allowance is made for an incremental quenching of CO$_2$ vibrations by collisions with CO and O. Significant infrared heating and thermal emission by CO$_2$ occurs only at levels where CO and O are minor constituents. Thus for the spatial variation of infrared heating we use the results of Dickinson (1973), based on the assumption of a pure CO$_2$ atmosphere.

The rate of EUV energy absorption is as before multiplied by an efficiency factor of 0.3 to allow for loss into chemical and airglow energy. The efficiency for $\lambda > 1080$ Å radiation is calculated as described previously. In contrast to our earlier calculations it is not possible to exhibit a priori the heating and molecular dissociation rates, since these depend (in the thermosphere at least) on the calculated composition. These determined rates may be inferred in part from the diagnostic balances shown in Section 4.

3. Numerical procedures

The system to be treated consists of four time-dependent scalar equations [(4), (12) and (6); which represents two scalar equations] and the two diagnostic equations [(2) and (5)]. The horizontal geometry is spherical, as indicated in Fig. 1. Grid points, 37 in number, are placed every 5° in the solar zenith from subsolar to antisolar points. The vertical range of the model is from $z = -10$ to $z = 7$, spanning an altitude range from below 100 km to more than 200 km above the planetary surface. A uniform vertical mesh with eight points per scale height is used, i.e., $\Delta z = 0.125$.

A time step is chosen small enough to insure computational stability for the time differencing used. Evaluation of all terms at the previous time step, referred to as "explicit differencing," is simplest to implement. However, the time step would then be limited by the least time required for a disturbance to travel between any two mesh points. For the assumed vertical increments, time steps less than $\Delta t \sim 10^{-4}$ s would be required by vertical diffusion in the vicinity of the top boundary. It would be impossible to carry out the integration with such short time steps, so we have made implicit the differencing for vertical diffusion. The hydrodynamic transport terms remain explicit since stability
considerations for these only limit time steps to less
than \(\sim 10^3\) s.

Subscripts \(n\) are used to denote variables evaluated
in the \(n\)th step, and \(\bar{\omega}\) is used to denote
\[
\bar{\omega} = \frac{\rho}{\rho_0} \omega.
\] (13)

The semi-implicit, leap frog time-differenced version of
(4) is then written
\[
\frac{\mu}{2\Delta t} \frac{\partial}{\partial \zeta} \frac{\partial^2}{\partial \zeta^2} \frac{\partial}{\partial \theta} \left( \frac{\partial}{\partial \zeta} \frac{\partial}{\partial \theta} \right) + \frac{\partial}{\partial \theta} \frac{\bar{\omega}}{r} \left[ \frac{\partial u_{n-1}}{\partial \theta} + \frac{\partial u_{n-1}}{\partial \zeta} \frac{1}{r} \frac{\partial}{\partial \theta} \right] F_{n-1}.
\] (14)

Likewise, (12) is written
\[
\frac{T_{n} - T_{n-2}}{2\Delta t} = \left( \frac{\partial}{\partial \zeta} \frac{\partial^2 T}{\partial \zeta^2} + \frac{\partial}{\partial \theta} \frac{\partial}{\partial \zeta} \frac{\partial}{\partial \theta} \right) T_{n-1} - a(T_{n} - T_{n-1}) - \left[ \frac{u_{n}}{r} \frac{\partial T_{n-1}}{\partial \theta} + w_{n-1} \right]
\times \left[ \frac{\partial T_{n-1}}{\partial \zeta} + \left( \frac{K}{C_p r^\theta} \right) T_{n-1} \right] \frac{Q}{C_p}.
\] (15)

To write (6) in finite-difference form, we define
\[
A_r = \alpha^{-1} \left( \frac{T_{m}}{T_{mCO_2}} \right)^{0.28} \left( \frac{\bar{\omega}}{\rho_0} \right)^{-1}
\]
\[
E_r = \begin{pmatrix} e_{11} & 0 \\ 0 & e_{22} \end{pmatrix}
\]
\[
D_r = e^r A_{r-1} \left[ -\frac{\partial}{\partial \zeta} \psi - \frac{\partial}{\partial \zeta} \psi_{n-1} - \frac{\partial}{\partial \zeta} \psi_{n-1} \right]
\]
\[
+ e^r \left( \frac{\partial A}{\partial \zeta} \right)_{n-1} \left[ \frac{\partial \psi}{\partial \zeta} - \frac{\partial \psi_{n-1}}{\partial \zeta} \right] + \psi_{n}.
\] (16)

so that (6) is
\[
\frac{\psi_{n} - \psi_{n-2}}{2\Delta t} = -L_n \left[ \frac{u_{n}}{r} \frac{\partial \psi_{n-1}}{\partial \theta} + w_{n-1} \frac{\partial \psi_{n-1}}{\partial \zeta} \right] + S_n.
\] (17)

Away from boundaries, all \(s\) and \(\theta\) derivatives are evaluated using centered second-order differences. Following
each time step, we integrate (2) upward from the bottom boundary to get \(\Phi\) from \(T\) and (5) downward from the upper boundary to get \(w\) from \(u\).

a. Boundary conditions

Vertical mesh points are numbered from bottom to top by \(j = 1, \ldots, J\). At the top we assume no diffusive

transports of heat, momentum, or relative mass concentrations expressed at the level \(j = \frac{1}{2}\) by
\[
\frac{\partial T}{\partial \zeta} - T_{j} - T_{j-1} = 0
\]
\[
\frac{\partial u_{j}}{\partial \zeta} - u_{j} - u_{j-1} = 0
\]
\[
\frac{\partial \psi_{j}}{\partial \zeta} - \psi_{j} - \psi_{j-1} = 0
\]
\[
\frac{\partial \psi_{j}}{\partial \theta} - \psi_{j} + \psi_{j-1} = 0
\]
\[
\frac{\partial w}{\partial \zeta} - \frac{1}{r} \frac{\partial}{\partial \theta} (u \sin \theta) = 0.
\] (21)

At the subsolar and antisolar points we assume from
continuity considerations that \(u\), \(\partial T/\partial \theta\) and \(\partial \psi/\partial \theta\) vanish. At the bottom we assume that at the \(\frac{1}{2}\) level
\[
\Phi = 0.
\] (22)

If \(w > 0\): \(T = T_{j}\), \(u = 0\), \(\psi = 0\); if \(w < 0\): forms of (4), (6) and (12) without the vertical diffusion terms are used. These are obtained by neglecting the terms proportional to \(e^r\), since these are negligible at the lowest levels. The resulting equations are differenced as described above except that a centered two-point difference is used to describe first derivatives with respect to the vertical coordinate. These difference equations then provide lower boundary conditions for \(T\), \(u\) and \(\psi\).

In summary, the three time-dependent equations (4), (6) and (12) have boundary conditions at top and
bottom depending respectively on \(u\), \(\psi\) and \(T\) or their derivatives, whereas the diagnostic equation (2) is integraled upward from an assumed constant \(\Phi\) at the lower boundary and (5) is integrated downward from the top boundary where \(\partial w/\partial \zeta\) is assumed to vanish.

b. Stability considerations

We would like to find a time-independent (steady-state) solution to the above difference approximations
to the system of differential equations. Because of the
nonlinearities of the system, we do not know \(a priori\)
whether a steady-state solution exists, or if it does,
whether integrations starting from any particular initial
conditions will evolve into the steady-state solution.
For solutions to the difference equations to approxi-
mate solutions to the differential equations, they must
be stable. In practice, stability implies the absence of
solutions with rapidly growing amplitudes. To avoid
rapidly growing "linear" instability it was found ne-
necessary to use \(\Delta t\) no larger than \(400 \) s; however, a slow
instability attributed to the nonlinear hydrodynamic
terms remained. Such instabilities are commonly en-
countered in the nonlinear time-dependent hydrody-
namic equations solved by finite differences and without
"horizontal diffusion." This instability is a consequence of the inability of the nonlinear terms in such models to cascade energy to higher wavenumbers than resolved by the mesh spacing. A common cure is the addition of "nonlinear viscosity" that essentially simulates the effect of wavenumbers higher than resolved by the model on the wavenumbers carried by the model. Leith (1971) has derived from two-dimensional turbulence theory an eddy dissipation damping term depending on wavenumber. For large wavenumbers this term becomes proportional to the fourth power of the wavenumber.

Consequently, for eddy dissipation, we adopted a damping proportional to the fourth power of the wavenumber. After a fixed number of time steps (we used five), the values of all the fields at each level were Fourier-analyzed into wavenumber components and each component was damped back to zero at a rate proportional to the fourth power of wavenumber. For simplicity, we gave all grid points equal weight rather than weighting them according to area. The damping rate was tuned to achieve the minimum value that would stabilize the difference equations. Ordinary linear viscosity depends only on squared wavenumbers and hence is not nearly so scale-selective. Our procedure removes the "aliasing" noise on the grid-interval scale while but little damping the global-scale components. One disadvantage of our damping prescription is that it does not conserve the global mean of variables which would be conserved in a more realistic prescription.

Before final selection of the difference scheme for this model, a large number of variations were tried, some of which appeared as satisfactory as the one described above. One further difficulty, apparently due to the leap-frog time differencing, was a very slow separation between the even and odd time step solutions, leading to eventual instability. This was cured by averaging even and odd time levels after every 70 time steps.

Since the preparation of this manuscript we have replaced the wavenumber-dependent smoothing by a nonlinear eddy diffusion term defined by finite differ-
ences. This has both shortened somewhat the required computer time and decreased numerical errors in the compositional equations in regions of large gradients.

4. Numerical solution

The numerical approximations described in the previous section were integrated in time, starting with the steady-state solutions derived according to the previous linear uncoupled dynamic model (Dickinson, 1971) and the composition model for these given winds (Dickinson and Ridley, 1972). An essentially steady-state solution was reached after carrying this out for something on the order of ten earth days of time (several thousand time steps). No exact statement of the time elapsed can be made because numerous changes were made in the numerical procedures in the course of the integration. These changes involved seeking the maximum time step and minimum Fourier damping required for a stable solution for the various difference schemes that we tried. Since evolution to a steady state required on the order of 10 h of computer time on a Control Data Corporation 7600, we considered it too costly to attempt to redo the calculation ab initio with the numerical procedures finally adopted.

For fixed winds, the composition would be expected to reach essentially a steady state in the time it takes a parcel of fluid to travel in a trajectory from the bottom boundary on the dayside to the bottom boundary on the nightside. This is on the order of an earth day or less above $z = -6$, but becomes several days or more near the bottom boundary. Because of the strong radiative damping, the motions and temperature would be expected to come into equilibrium in less than a day for fixed composition. In agreement with this discussion, the distribution of composition near the lower boundary appeared to take longest to adjust to equilibrium values and still exhibited some slow trends at the termination of the integration. Because CO and O are in concentrations of less than 1 part per thousand at these levels, further integration would not be expected to alter significantly those variables above $z = 0$ or any other variables.

The final solutions obtained are shown in Figs. 2 and 3. It is somewhat difficult to describe simply the relationships between the different variables in a sequential fashion because they are all strongly coupled together. The upper left frame of Fig. 2 shows departures from the global mean of the height of constant pressure surfaces. The gradient gives the pressure force driving horizontal winds according to Eq. (4). Winds on the dayside are pushed in the nightward direction as expected, but in the vicinity of the terminator the height gradient changes sign and acts to retard motion directed toward the antisolar point. The pressure force on the nightside is toward the dayside everywhere except for a region above $z = 5$ where the height gradient again reverses. Motions much above $z = 0$ are more controlled by the viscous coupling with lower layers than by the height gradient. On the dayside the horizontal winds (shown in the lower left frame of Fig. 2) are everywhere directed toward the antisolar point. The reversed height gradient rapidly slows down the horizontal winds above $z = 0$ as they cross the terminator. At the higher levels, a weak wind reversal occurs over the half of the nighttime hemisphere nearest the subsolar point. The sharp convergence of horizontal winds generates a tongue of large downward vertical velocities (as seen in the upper right frame of Fig. 2) and the adiabatic heating generated by these vertical motions produces the belt of warm air nightward of the terminator seen on the lower right frame of Fig. 2.

The reversal of height gradient at upper levels in Fig. 2 must be interpreted in terms of the calculated concentrations of O and CO shown in Fig. 3. The calculation shows that motions build up a large bulge in the light gas concentrations on the nightside relative to their dayside values. This build-up is interpreted in terms of trajectories of air parcels, as represented by mass streamlines in Fig. 4. Parcels rise on the dayside due to net heating and, in the absence of vertical diffusion, increase in relative concentrations of light
Fig. 5. Rates of change given at the $z = -6$ level by various terms in the momentum equation, normalized by dividing by $u$. Positive values are directed toward the antisolar point.

Gases through photodissociation. On the nightside parcels sink due to net cooling and, again assuming no vertical diffusion, conserve their relative concentrations. Thus, at levels in the vicinity of or below $z = 0$, the concentrations of O and CO on the nightside in Fig. 3 are essentially constant along the streamlines of Fig. 4. Higher level concentrations are largely determined by the values in the vicinity of $z = 0$ through diffusive coupling. (Recall that the pressure level where $z = 0$ was originally selected as the level where, for global-scale motions, scaling arguments indicate that vertical diffusive transports are comparable to hydrodynamic transport terms.)

To summarize, concentrations of minor constituents increase on the nightside toward the antisolar point because locations at greater distances from the terminator receive molecules crossing the terminator at higher levels. The lighter gases on the nightside define a smaller $\bar{m}$ which according to (2) is equivalent to a warmer temperature in increasing the heights of constant pressure surfaces. Thus, the large reversal of height gradient above $z = 2$ as seen in Fig. 2 is a consequence of the nightside bulge of light gases.

The height gradient reversal, which below $z = 0$ occurs at larger solar zeniths at lower levels, is determined in part by the belt of warm air produced by the strong downward motions. The axis of maximum horizontal and vertical velocities likewise moves toward the antisolar point with decreasing altitude.

Further interpretation of the solutions and comparison with our previous simplified calculations is given in the next section. The remainder of this section is devoted to analysis of the balances between various terms in the governing equations, which were examined every 70 time steps as a diagnostic device to determine degree of convergence to a steady state. Fig. 5 shows final values obtained for the various terms in the momentum equation evaluated at $z = -6$. These terms have been normalized by dividing by $u$ and so give inverse time scales for change of $u$. The horizontal pressure force nearly balances the nonlinear hydrodynamic-inertial terms. Drag by molecular viscosity is negligible. There is a residual tendency term smaller than the dominant terms by an order of magnitude or more. This imbalance is inserted by the Fourier damping procedure where the largest gradients occur and does not disappear with further integration in time. The balances at other levels are qualitatively similar on the dayside except that above $z = 2$ viscous drag assists the inertial terms in counteracting the pressure forces. On the nightside above $z = 0$ and within $60^\circ$ of the antisolar point viscous drag becomes more important than inertial terms in balancing the pressure forces. Over small regions near the location of maximum horizontal wind decrease around $z = 2$ and $z = -4$, the net tendency due to Fourier damping gets nearly as large as the other terms. This imbalance indicates that the $5^\circ$ grid is too coarse at these points to adequately resolve the sharp changes in the velocity pattern. An overall view of the momentum equation balances is given by taking area-weighted averages of the terms at each level. Table 1 presents this calculation for the momentum equation at each scale height but without dividing by $u$. Above $z = 0$, the viscous, height gradient and inertial terms are

<table>
<thead>
<tr>
<th>$z$</th>
<th>$\frac{g}{\rho u^2} \frac{\partial \mu}{\partial z}$</th>
<th>$-\frac{u}{\partial z}$</th>
<th>$-\frac{u}{\partial \theta}$</th>
<th>$\frac{1}{\partial \theta}$</th>
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Fig. 6. Rates of temperature change at \( z = -4 \) and \( z = 2 \) given by the various adiabatic transport terms in the thermodynamic equations and normalized by dividing by \( T \). Positive sign indicates heating.

Generally of comparable importance. From \( z = 1 \) to \( z = 5 \) the mean height gradient drives motions in the dayward direction, and it is the inertial terms that drive motions in a nightward direction. A large nightward-directed pressure force develops in the vicinity of \( z = -3 \), which is resisted by the inertial terms. Vertical motions are especially important in the mesosphere for the vertical redistribution of momentum, smoothing out possible vertical variations in the horizontal winds. The global mean net tendency due to the Fourier filtering is comfortably small but not negligible.

Figs. 6 and 7 display the balances between various terms in the thermodynamic equation at \( z = -4 \) and \( z = 2 \). For these levels Fig. 6 shows the contribution of the various hydrodynamic transport terms to the net adiabatic heating by motions. At and below \( z = -4 \), the adiabatic cooling can be closely approximated by the product of vertical motion and gravity, i.e., the expansion cooling of an isothermal atmosphere. As seen in Fig. 6, the maxima contributions of the horizontal heat transport term, \( (u/v)(\partial T/\partial \theta) \), occur on either side of the temperature minimum that occurs at higher levels near the terminator and shifts nightward at lower levels. Heat is extracted on the nightward side and added on the dayward side of the temperature minimum, helping to shift the minimum in the nightward direction. Around \( z = -2 \) this cooling on the nightward side is comparable in magnitude to the heating due to the vertical motion terms. The \( w \partial T/\partial z \) advection term becomes as large as the expansion cooling term near the \( z = 0 \) level where thermospheric temperatures are rapidly increasing upward and hence there are large vertical temperature gradients. Above \( z = 2 \), where temperatures vary little in the vertical, the expansion cooling term is again dominant.

Fig. 7 compares the sum of the terms in Fig. 6 with the radiation and conduction terms. Molecular conduction removes a significant fraction of the dayside heating at and above \( z = 2 \). Near \( z = 0 \) it acts to supplement the solar heating in balancing the adiabatic cooling, and below \( z = -2 \) it becomes negligible. The relative contribution of the Newtonian cooling term maximizes near \( z = -4 \) where the Newtonian cooling coefficient has its largest value. It becomes relatively unimportant near \( z = -2 \) and below \( z = -6 \) due to smaller values of the coefficient with comparable magnitudes of horizontal temperature variation. Between \( z = 0 \) and \( z = 2 \) it becomes significant again due to the order-of-magnitude increase in horizontal temperature contrast above \( z = 0 \) but at higher levels it becomes negligible with the exponential decrease of the cooling coefficient.

The net infrared cooling on the nightside is given by the sum of the specified heating term and the Newtonian cooling term. It is, of course, desirable that the sum of these two terms closely approximate the infrared cooling that would be derived from exact calculations of cooling for the model-determined vertical temperature profiles. Below \( z = 0 \) the departures of temperature form global mean values is sufficiently small that the Newtonian cooling approximation should be adequate. For the \( z = -4 \) level Fig. 7 shows that within \( 45^\circ \) of the antisolar point the net infrared cooling is an order of magnitude smaller than either of these separate terms. Reference to the temperature contours in Fig. 2 shows...
minimum temperature in this region. It is plausible that warming due to radiative transfer from lower levels could nearly cancel the "cool-to-space" contribution at this level. At lower levels as well as at higher levels up to $z \sim 0$, the Newtonian cooling term is generally less in magnitude than the mean cooling term so that large net infrared cooling on the nightside results. However, as seen in Fig. 7, at $z = 2$ the Newtonian cooling term determines a net infrared radiative heating. At this level the local temperature is 120 K colder than the global value used to derive the Newtonian cooling term. Past study of the validity of the Newtonian cooling approximation has indicated it would still be useful for positive perturbations of this magnitude but would underestimate net cooling for such a negative perturbation.

Since the preparation of this manuscript we have revised our radiative cooling formulation to allow properly for the nonlinear temperature dependence of the infrared cooling. This change reduces the magnitude of the perturbation cooling term at $z = 2$ to less than 80% of that of the mean cooling so that the total cooling rate at that level is everywhere positive and order of $3 \times 10^{-4}$ s$^{-1}$ near the antisolar point. The effect

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of this change on other results is interesting. In particular, the temperatures at the antisolari point above \( z = 2 \) drop from \( \sim 270 \) to \( \sim 230 \) K. The maximum high level thermospheric winds are increased by \( \sim 50 \) m s\(^{-1}\), the region of dayward winds above \( z = 0 \) on the night-side is replaced by \( \sim 25 \) m s\(^{-1}\) dayward winds, and the region of maximum convergence of the horizontal winds shifts \( 15^\circ \) toward the antisolari point. Several other features in Fig. 2 similarly shift. In particular, the axis of minimum heights, the axis of maximum downward vertical motions, and the tongue of warm air. With this shift, the region of \( T < 160 \) K at the antisolari point disappears but a region of \( T < 160 \) K appears near \( z = -4 \), and \( \theta = 110^\circ \) in the cold tongue on the dayside of the warm tongue. These changes are not regarded as sufficiently significant to warrant revision of the figures, but do serve to indicate the sensitivity of our results to details of the radiative heating prescription.

Table 2 compares the area-weighted global means of temperature rates of change for the various terms in the thermodynamic equation. Above \( z = 0 \), the primary balance is between the net temperature-independent radiative heating and the cooling by downward molecular conduction. At the highest levels, down-gradient horizontal transport of heat contributes an order of \( 10\% \) of the EUV heating to the net heat input. At \( z = 0 \) there is net radiative cooling, which balances a net heating by conduction on the order of \( 200 \) K per earth day. Maxima in the Newtonian cooling term contribution to global mean heating occur at \( z = 1 \) and \( z = -6 \). The first of these is required to balance net adiabatic cooling on the order of \( 100 \) K per earth day, largely due to correlation of downward motions with regions of warmer temperature. It is interesting to note that this transport mechanism for cooling is analogous to the mechanism of mean light gas downward transport by correlation between high mixing ratios and downward motion. Whereas the distribution of light gases is largely determined by this mechanism, its only effect on the mean temperature profile is to decrease it by a few degrees from what it would be in radiative-conductive equilibrium. Again at \( z = -6 \), vertical motions produce net cooling. This cooling is partially cancelled by horizontal transports and is only on the order of \( 20 \) K per earth day. The Newtonian cooling and vertical motion terms essentially balance. The decrease of global mean temperature from its radiative value at this level is also as much as a few degrees due to the smaller cooling coefficient. The magnitude of the other radiative component is only half as large and at this level derives from differences on the order of \( 1 \) K between the assumed temperature profile and the actual global mean radiative equilibrium profile appropriate to net radiative heating calculated according to Dickinson (1973). Finally, the last column of Table 2 confirms the conclusion from Fig. 7 that the \( \partial T / \partial t \) term introduced by smoothing is negligibly small at all levels.

To illustrate the balances between various rate terms in the compositional equations, we show their values at \( z = 2 \) in Fig. 8. The dominant dayside balance at essentially all levels is between photoproduction and removal by hydrodynamic transport. The only exception occurs over the top few scale heights of the model where photoproduction is reduced because of low CO\(_2\) concentration and the O and CO removed by transport.
Table 3. Area-weighted global mean of various terms in the atomic oxygen transport equations. Units 10^{-8} s^{-1}.

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are restored primarily by upward diffusion. In general, above $z=0$ diffusion serves the important role of maintaining vertical distributions close to diffusive equilibrium. At the $z=2$ level shown here, transport largely tends to increase the light gases on the nightside and vertical diffusion again removes them. The rate of removal by transport on the dayside is much greater than the rate of addition on the nightside, resulting in a net global mean removal. Tendency terms introduced by Fourier smoothing are large primarily on the nightside of the terminator. These terms are non-negligible at all levels and become uncomfortably large near $z=-4$. They seem to result from application of our spectral smoothing procedure to compositional fields where there are sharp changes in horizontal gradient. Reference to Fig. 3 shows that for atomic oxygen this occurs in the vicinity of the terminator above $z=0$ and approximately 20° to the nightward side of the terminator below $z=0$. The 0.5% contour for O in Fig. 3 falls near the $z=-4$ level of greatest smoothing error. This contour shows a sharp dip at $\theta=110^\circ$ and a sharp peak at $\theta=125^\circ$. The net tendency due to transports is to fill the compositional minimum at the peak with O from the dip on the order of 20 time steps, but the smoothing procedure continues to reinsert the non-equilibrium values. In a new version of the model with nonlinear diffusion prescribed by a finite-difference diffusion term rather than by spectral smoothing, we have significantly reduced but not entirely eliminated this difficulty.

Examination of area-weighted global mean values in Table 3 indicates that the mean rate of imbalance due to smoothing is an order of magnitude less than the photoproduction term between $z=-4$ and $z=4$. Below $z=-4$ it acts as an additional mass source of magnitude comparable to photoproduction, whereas above $z=4$ it acts to remove mass at a rate which is less than the transport rate of removal. Elimination of the smoothing error could only be expected to significantly affect the calculated nightside concentrations of O and CO below $z=-4$. A calculation with higher horizontal resolution could conceivably reduce nightside concentrations at some levels below $z=-4$ by as much as a factor of 2 but the dayside values should remain essentially unchanged.

5. Discussion and conclusions

The present calculation was intended, in part, to eliminate uncertainties in the predictions of our previous models as introduced by various simplifications. Comparison of the present Fig. 2 with Fig. 2 of Dickinson and Ridley (1972) shows there to be little difference in the dayside winds above $z=0$ except for a reduction in wind strength by 25%. A mesospheric jet of 500 m s^{-1} was previously calculated on the dayside of the terminator centered at $z=-3.5$. The present calculation shows a smooth decrease of horizontal wind velocities with decreasing elevation to values of only 150 m s^{-1} in the vicinity of the previous jet. We earlier raised the question as to whether the mesospheric jet would be unstable due to small local Richardson numbers. Inclusion of nonlinear terms in the dynamic equations now acts to smear out the jet sufficiently to increase its smallest Richardson number by an order of magnitude and so definitely to preclude the possibility of shear instability.

The calculated horizontal variations in composition now elevate constant pressure surfaces on the nightside to block the flow crossing the terminator. Consequently a weak velocity stagnation zone centered at the terminator is predicted to exist in the thermosphere over much of the nightside. Downward velocities of up to several meters per second are predicted near the terminator in front of this stagnation zone. The net adiabatic heating by the downward motions produces a pronounced tongue of warm air below $z=1$ which tends to shift the height minimum from near the terminator at higher levels toward the antisolar point. Consequently, the field of maximum horizontal and vertical velocities and maximum temperature all shift together toward the antisolar point with decreasing elevation in the mesosphere.

The dayside temperatures in the thermosphere above $z=0$ also resemble those derived with the previous simpler model. Exospheric dayside temperatures of around 600 K are again predicted. There is some indication from previous observations (Kumar and Hunt, 1974) and from the UV spectrometer on the recent Mariner 10 flyby of Venus (Broadfoot et al., 1974) that dayside exosphere temperatures are as low as 400 K. More detailed comparison of calculated and "observed" temperatures awaits further analysis and publication of these observations. It should be pointed out, however, that exospheric temperatures as low as 400 K on the dayside would require drastic changes in the model input, e.g., a factor of 2 decrease in the rate of EUV heating or a greatly enhanced rate of infrared cooling. There is evidence to be discussed elsewhere that the
EUV heating efficiency is closer to 10% than 30%, but on the other hand, the assumed solar fluxes could already be low by a factor of 2 (e.g., Robe and Dickinson, 1973). Any large increase of the day-to-night transport of heat appears extremely unlikely. Large eddy mixing might remove enough heat, but as Hunten (1974) argues, the frictional dissipation of the waves producing the mixing might add more heat than would be removed by the mixing.

The calculated nightside exosphere is quite cold, \( \sim 250 \text{ K} \), as a consequence of the effect of compositional blocking in minimizing hydrodynamic warming.

Temperatures in the mesosphere show much larger horizontal contrasts (on the order of 20–40 K) than obtained previously. These temperature contrasts are required to develop constant pressure surface height gradients sufficiently large to drive the circulation against nonlinear inertial terms and with enough strength to provide the redistribution of thermal energy required by the thermodynamic equation.

The calculated relative concentrations of O and CO shown in Fig. 3 are qualitatively similar to those given by Fig. 5 of Dickinson and Ridley (1972). Two differences worth emphasizing are (i) the small horizontal variation of nightside concentrations near the antisolar point due in part to the weakened and reverse circulation reported here and (ii) an increase by a factor of 2 in dayside concentrations relative to previous values. Predicted concentrations of O and CO at the level of the F1 peak \( (z \approx 1) \) are now 4%. This increase in O and CO results from a reduction in the strength of the dayside circulation below \( z = 0 \) by a factor of 2. Previously, vertical motions at these levels were essentially derived from near balance between net radiative heating and adiabatic cooling as given by the product of vertical motion and global mean static stability. In the present calculation, the local dayside static stability near \( z = 0 \) is significantly larger than its global mean value so the same adiabatic cooling is produced by weaker vertical motions. Furthermore, some of the net radiative heating is balanced by the horizontal transport term and the increased temperature contrast from day to night also allows significant dayside cooling by the Newtonian cooling term. All these factors contribute to the observed reduction of vertical motions and hence to the increased values of O and CO.

Finally, we note that our calculations would be more conveniently related to observations and used to derive "model atmospheres" if they were presented relative to an altitude scale. We have postponed this step until after we have had an opportunity to compare in detail our results with Mariner 10 observations. Rough estimates of altitude can be made by referring our \( z = 0 \) level to a distance from planetary center of 6185 km.

The three most significant modeling assumptions we have made out of ignorance and on the basis of simplicity are that 1) mixing by eddy motions is negligible, 2) the observed "4-day rotation" of the stratosphere does not extend up to the levels of our model, and 3) the heating efficiency of the EUV radiation is 0.30. Eddy mixing would generally transport O and CO downward but would reduce the large-scale transport by weakening the circulation. Extension of the 4-day circulation upward into our model region could induce changes as great as a factor of 2 in our predicted results for the mesosphere but would have lesser effect in the thermosphere where the day-to-night winds would still be much larger than the zonal wind. The largest change would probably be an increase of the O and CO concentrations due to a "short-circuiting" of the transport mechanism. Further exploration of this question awaits a quantitatively valid model of the 4-day circulation. The fraction of solar EUV that goes into heating should become more certain with analysis of UV spectrometer data from Mariner 10 and the future Venus Orbiter. The 0.3 used here was adopted as a likely value on the basis of Stewart's (1972) analysis of Mariner 6 and 7 data but a lower value may be appropriate.

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REFERENCES


