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Structure of a photoionization layer in the solar chromosphere

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Abstract. We investigate the structure of a pure photoionization layer in the solar chromosphere on the basis of a one dimensional model in which an incident flux of EUV photons from above ionizes neutral hydrogen assumed to be flowing up with steady subsonic speeds from below. It is useful to analyze the structure equation in the phase plane of the hydrogen density (or velocity) and the ionizing photon flux. This shows that there is a unique critical solution which links the bottom of the layer to the top. Its structure is analogous to a classical constant pressure weak deflagration. The model is extended to include the ionization of minor species. In contradiction to previous work we note that one dimensional steady models are incapable of giving rise to a FIP (first ionization potential) effect by virtue of the fact that in a collision– dominated situation all minor species enter the layer at the hydrogen speed and exit at the proton speed and their fluxes are conserved.

Key words: Sun: chromosphere – solar wind

1. Introduction

The connection between the lower solar corona and the chromospheric network from which the solar wind is believed to emanate [Axford & McKenzie (1993)] is complicated both by the complex magnetic field topology and the different physics operating in different regions. Elsewhere [McKenzie et al. (1997)] we have analyzed how the solar wind plasma flow in magnetic funnels located at the boundaries of the network, are connected to thin (O(10 km)) ionizing layers situated at the bottom of these funnels. In this case the dominant ionizing agent is electron impact driven by a downward electron heat flux from the lower corona. However in the lower regions of the layer, where the temperature is low (O (few 10^4 K)) the impact ionization length can exceed the photoionization length scale (O(40 km)) and therefore it is relevant to consider the structure of a pure photoionization layer.

As a model we take a one dimensional, steady configuration in which a downward flux of EUV photons ionizes neutral hydrogen flowing upward. The structure equations show that there is a unique critical solution which links the physical variables at the bottom of the layer to those at the top. Both hydrogen atoms and protons are accelerated within the layer so that the protons exit the layer with a number density relative to that of the incoming hydrogen equal to the ratio of the entering hydrogen speed to the exiting proton speed. The structure is completely analogous to a classical weak, constant pressure, deflagration [Courant & Friedrichs (1963)], or a very weak D–type ionization front [Axford (1961)]. For the case in which there are sufficiently frequent hydrogen - proton collisions in a characteristic photoionization length scale to maintain approximately equal speeds for hydrogen atoms and protons we have obtained a simple analytic solution which neatly highlights the properties of the layer. In particular the ratio of the densities of the exiting protons to the entering hydrogen atoms is simply $T_H/(T_p+T_e)$, where $T_j (j = H, p, e)$ is the temperature of each component.

It is of interest to generalize this model to include the photoionization of minor singly ionized elements in the presence of the hydrogen-proton background. In subsonic flow the momentum equations for the neutrals and their ionized counterparts reduce simply to the balance between the partial pressure gradients and collisional friction with the hydrogen–proton background flow. It is evident that the neutrals enter the layer at the hydrogen speed and their ionized counterparts exit at the proton speed. Therefore no fractionation, i.e. an enhancement or depletion of any species relative to another, can occur and hence such one dimensional, steady flows cannot give rise to a FIP effect as, indeed, appears to be confirmed by observations [Geiss et al. (1995)].

Our calculation should be regarded as merely illustrative because it is artificial in that it requires that we have just the right amount ionizing of photon flux for each species. In reality there exists an excess of such photon fluxes in the form of $Ly\quad\alpha$, in which case the ionizing layer for each species would not be stationary and in fact would move down, so that the elements would come in ionized from deep down. We should then regard the metals as coming into the hydrogen-proton layer fully ionized and with the same speed as hydrogen. Since all ions exit the layer at the proton speed again no fractionation can take place. In order to obtain such an effect, namely an excess of metals,

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one must have a procedure for removing hydrogen in some way, or by implanting metals.

On the basis of the assumptions made in these calculations, elements with photoionization cross-sections greater than that of hydrogen (e.g. Ar, AL, Si, C) are ionized deeper in the layer than hydrogen, whereas those with smaller cross-sections (e.g. Fe, O, Mg) are ionized higher up, with the notable exception of helium. This is rather misleading however: we have not allowed for the fact that the low FIP species are fully ionized long before entering the layer since there is a great excess of ionizing photons. The ionizing time determined by convolving the cross-sections with the corresponding ionizing flux is probably a more useful means of distinguishing the behaviour of different species (e.g. Geiss & Bochsler (1986)). In any case the scope of the analysis must be altered somewhat since photo-ionization is not a significant cause of photon loss in these circumstances and it is therefore necessary to describe the distribution of $Ly \quad \alpha$ photons in particular in a different way.

2. Model equations for a photoionization layer

We adopt a simple one dimensional model of a photoionization layer in the solar chromosphere in which a flux of EUV ($\sim$ 900Å) photons incident from above ionizes neutral hydrogen flowing up with subsonic speeds from below. The photon flux J decays into the layer according to

$$\frac{dJ}{dz} = \alpha N J, \tag{1}$$

in which z is directed upwards and J decays downwards, α is the cross section ($\sim 6.3 \cdot 10^{-18}$ cm^2, see Allen (1976)) for photoionization by $Ly \quad \alpha$ and N is the neutral hydrogen density in the layer. The decay of the photon flux is accompanied by the production of protons with density n and speed v according to

$$\frac{d(nv)}{dz} = \alpha N J = -\frac{d(NV)}{dz}. \tag{2}$$

The total particle flux of hydrogen and protons is conserved throughout the layer, that is

$$NV + nv = const, \tag{3}$$

where V is the neutral hydrogen speed within the layer. At the bottom of the layer ($z = -\infty$) there are no protons ($n = 0$) whereas at the top ($z = +\infty$) there are no neutral hydrogen atoms ($N = 0$). The constant in (3) is then just the hydrogen flux at the bottom, $N_b V_b$, say. Eq. (1) and (2) yield the simple integral

$$J = nv, \tag{4}$$

since the ionizing photon flux is assumed to be completely extingiushed at the bottom of the layer. Therefore if J_∞ denotes the total incident ionizing photon flux, particle conservation can be written

$$NV = J_\infty - J, \tag{5}$$

because $J_\infty = N_b V_b$.

Since the ionizing length scale ($(\alpha N_b)^{-1} \sim 40$ km) is much smaller than the gravitational scale height (300 km) and also because we assume the hydrogen is flowing upwards at low subsonic speeds, the momentum equations for neutral hydrogen and protons can be reasonably approximated by a balance between their pressure gradients and friction arising from collisions between protons and hydrogen atoms. Thus we can write

$$\frac{d}{dz}(nk(T_p + T_e)) = m_p n \nu_{pH}(V - v), \tag{6}$$

$$\frac{d}{dz}(NkT_H) = m_p N \nu_{Hp}(v - V). \tag{7}$$

Because total momentum is conserved, which in the subsonic approximation reduces to the constancy of total pressure. i.e.,

$$nk(T_p + T_e) + NkT_H = const, \tag{8}$$

the collision frequencies must obey the relation,

$$n_p \nu_{pH} = N \nu_{Hp}, \tag{9a}$$

$$\nu_{pH} = \frac{N}{(n+N)\tau_{pH}}, \quad \nu_{Hp} = \frac{n}{(n+N)\tau_{Hp}}. \tag{9b}$$

Thus the collision times are equal and can be expressed as

$$\tau_{Hp} = \tau_{pH} = \tau = \frac{1}{\sigma N V_t}, \tag{9c}$$

where V_t is thermal speed of hydrogen ($\sqrt{kT_h/m_p}$) and σ is the cross section ($\sim 5 \cdot 10^{-15}$ cm^2 , see Halsted (1972) for charge exchange collisions between hydrogen atoms and protons).

For simplicity we assume the temperatures T_p, T_e and T_H are given throughout the layer (for example isothermal). Thus we have five Eqs. (1), (4), (5), (6) and (7) for the five physical variables n, N, V, and J, which determine the structure of the layer.

3. Solution of the structure equations

Using the concept of a phase plane it is useful to analyze the structure in terms of how v, V, n and N vary with J throughout the layer. This is accomplished by dividing the stress balance Eqs. (6) and (7) by the photon flux decay Eq. (1) to obtain the following two differential equations for V and v, in which the spatial coordinate z is eliminated in favour of J,

$$\frac{dV}{dJ} = \frac{V}{J_\infty - J} \frac{c_H V(V - v) - (J + \frac{v}{V}(J_\infty - J))}{J + \frac{v}{V}(J_\infty - J)}, \tag{10}$$

$$\frac{dv}{dJ} = \frac{v}{J} \frac{(J + \frac{v}{V}(J_\infty - J)) - c_p v(V - v)}{J + \frac{v}{V}(J_\infty - J)}, \tag{11}$$

in which we have used Eqs. (4) and (5) to eliminate n and N in terms of J, and V. The constants c_p and c_H are given by,

$$c_p = \frac{m_p}{k(T_p + T_e)\alpha \tau_{pH}}, \tag{12a}$$

$$c_H = \frac{m_p}{kT_H \alpha \tau_{pH}}. \quad (12b)$$

In fact it is only necessary to solve one of the above differential equations, say Eq. (10) for the hydrogen speed V as a function of the ionizing photon flux J, since constancy of total pressure, Eq. (8), along with flux conservation Eqs. (4) and (5), yield v in terms of V and J, namely

$$v = \frac{((T_p + T_e)/T_H)J}{J_\infty/V_b - (J_\infty - J)/V}. \quad (13)$$

In this approach the structure of the layer is determined by Eq. (10) with the photon flux decay Eq. (1) merely serving as a reference which can be consulted for J as a measure of z. (For an explicit expression for $J(z)$ see the asymptotic solution given in Appendix (B)). In order to avoid a singularity at the upper boundary, where $J = J_\infty$, the numerator of the expression on the right hand side of Eq. (10), must vanish, so that, although there are no hydrogen atoms left, they must exit asymptotically at a speed given by

$$V(J_\infty) = \frac{v_p(J_\infty)}{2}\left[1 + \sqrt{1 + 4J_\infty/c_H v_p^2}\right]. \quad (14a)$$

The protons exit the layer at the speed $v_p(J_\infty)$ which follows from (13) as

$$v_p(J_\infty) = V_b(T_p + T_e)/T_H. \quad (14b)$$

Similarly at the lower boundary, where $J = 0$, Eq. (11) tells us that the protons must enter the layer at the speed,

$$v(J = 0) = V_b(1 - J_\infty/V_b^2 c_p). \quad (15)$$

The dimensionless parameter which characterizes this problem is,

$$\overline{c}_H = c_H V_b^2/J_\infty = \frac{V_b}{V_t^2}\left(\frac{L}{\tau_{Hp}}\right) = \frac{V_b}{V_t}\left(\frac{\sigma}{\alpha}\right). \quad (16)$$

The ratio of the cross sections is large since,

$$\frac{\sigma}{\alpha} \simeq \frac{5 \cdot 10^{-15}}{6 \cdot 10^{-18}} \sim 10^3. \quad (17)$$

Therefore even if hydrogen enters the layer very subsonically, say $V_b/V_t = 10^{-1}$ (corresponding to $V_b = 1$ km/sec and $V_t \simeq 10$ km/sec), the parameter given by (16) is indeed very much greater than unity. In these circumstances the asymptotic solution, given in Appendix B, provides a simple, analytic description of the structure of the photoionization layer.

The results of the numerical integration of the solution of Eq. (10) for the illustrative case in which $\overline{c}_H = 4$ are shown in Fig. (1).

Hydrogen atoms enter at speed V_b and density N_b and are depleted on exit. Protons enter the layer at speed $v(0)$ (given by (15)) and exit at twice V_b with half the density of the entering hydrogen atoms. The asymptotic solution, which is shown by

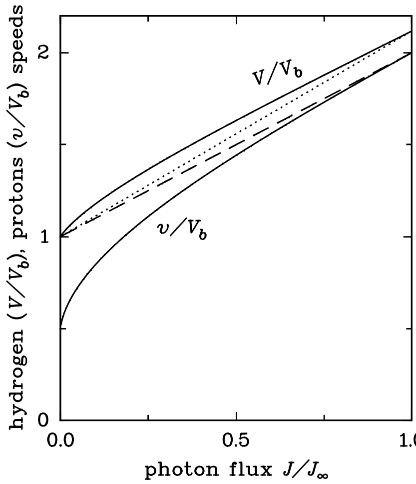

Fig. 1. The critical solutions for the hydrogen (V) and proton (v) speeds as a function of the photon flux (J) in the layer for the illustrative case of $\overline{c}_H = 4$. The upper dotted curve (V_+) corresponds to the locus on which the phase trajectories $dV/dJ = 0$ and the lower broken curve is the asymptotic solution (Appendix A) for large $\overline{c}_H$.

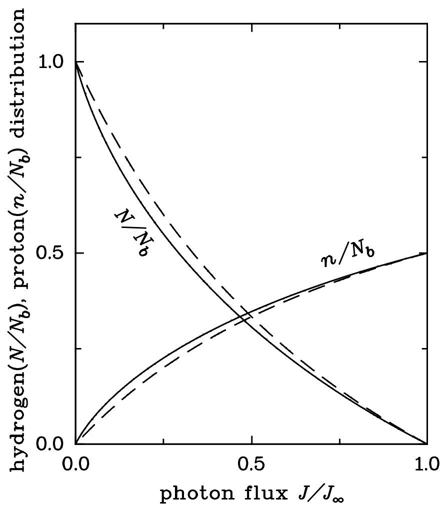

Fig. 2. The full curves show the hydrogen (N) and proton (n) densities in the layer as a function of J for the illustrative case $\overline{c}_H = 4$. The broken curves correspond to the asymptotic solution ($\overline{c}_H \gg 1$).

the broken curve for comparison, is indeed an excellent approximation even for the moderate value of 4 for the characteristic dimensionless parameter of the problem, Eq. (16). The curves of V and v as functions of J represent the critical solution of Eq. (10) which uniquely links the bottom of the layer to the top (but see also the discussion given in Appendix A, in which the structure is analyzed in the (N, J) plane). The corresponding distributions of hydrogen and protons are shown in Fig. (2).

4. Ionization of minor species

It is of some interest to consider the photoionization of minor singly ionized species in the presence of the hydrogen-proton background whose layer structure we have just analyzed. Since the minor species is ionized by its own photon flux J_m and also by the photon flux, J_h say, which ionizes hydrogen, the decay of J_m and the production of ions can be described by the equations

$$\frac{dJ_m}{dz} = \alpha_m(J_m + J_h)N_m, \tag{18a}$$

$$\frac{d}{dz}(N_m V_m) = -\alpha_m(J_m + J_h)N_m, \tag{18b}$$

$$n_m v_m + N_m V_m = const. \tag{18c}$$

in which α_m is the cross section for photoionization of species "m", and where $N_m(n_m)$, $V_m(v_m)$ are respectively the density and speed of the neutrals (and their ionized counterparts). The hydrogen–proton background is described by the previous set of equations in which J_h assumes the role of J. In a similar way to the hydrogen–proton layer we assume that the incoming subsonic neutrals are very subsonic so the pressure gradients of the neutrals and their ion counterparts are balanced by frictional forces with the hydrogen-proton background enabling us to write,

$$\frac{d}{dz}(N_m kT_m) = -\overline{m}\frac{NN_m}{\tau_{mN}}(V_m - V) - \overline{m}\frac{nN_m}{\tau_{mp}}(V_m - v), \tag{19a}$$

$$\frac{d}{dz}(n_m kT_m) = -\overline{m}\frac{Nn_m}{\tau_{+mN}}(v_m - V) - \overline{m}\frac{nn_m}{\tau_{+mp}}(v_m - v), \tag{19b}$$

where

$$\frac{1}{\overline{m}} = \frac{1}{m_p} + \frac{1}{m_m}.$$

Again we may specify the temperatures and proceed to solve these equations for the structure of a layer of minor species and their ion counterparts superimposed on the hydrogen–proton layer which is unaffected by the minor constituents. However it is clear from Eqs. (19) that the requirement of homogeneity on either side of the layer (i.e. no gradients at $z = \pm\infty$) implies that the neutrals enter the layer at the hydrogen speed V_b and exit the layer at the proton speed v_p (given by Eq. (14b)). Therefore from continuity (Eq. 18c) the ratio of the density of ions at the top to the density of their neutral counterparts at the bottom is simply

$$\frac{n_m(top)}{N_m(bottom)} = \frac{V_b}{v_p(J_\infty)} = \frac{T_H}{(T_p + T_e)}, \tag{20}$$

which exactly mimics the hydrogen-proton layer. Hence no fractionation can occur contrary to the argument made by Marsch et al. (1995) and Peter (1996). Their results appear to be an artefact of the application of an unsuitable boundary condition applied at some arbitrarily chosen lower boundary and an inappropriate description of the depletion of the photon flux. It is clear that in a one dimensional steady model, such as is described by the above equations, no enhancements or depletions of one type of ion over another can occur because continuity demands what goes in must come out. Therefore the density ratios are given by (20) for all species, because the partial pressures are constant on either side of the layer.

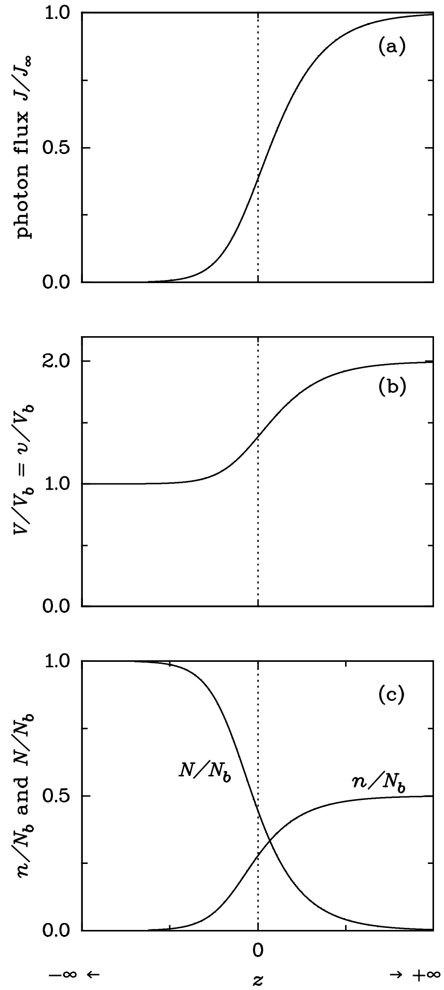

Fig. 3a–c. Distributions of J, $V = v$, N and n as functions of z as given by the asymptotic solution for the case $v_1 = 2$ (corresponds to $T_e = T_p = T_H$).

To solve the above system of equations one could proceed as before by making use of the v_m, V_m, J_m phase space. For the present purposes it is sufficient to employ the asymptotic solution given in Appendix B and assume that collisions are sufficiently frequent to maintain equal speeds for all particles. Thus with $v_m = V_m = v_p = V$, where V is given by Eq. (B2), the decay and continuity Eqs. (18) tell us that

$$n_m = J_m/V, \tag{21a}$$

$$N_m = (J_{m\infty} - J_m)/V. \tag{21b}$$

Substituting (21b) for Nm into (18a) yields the decay equation for J_m in the form

$$\frac{dJ_m}{dz} = \frac{\alpha_m(J_{m\infty} - J_m)(J_m + J_h)}{V_b + (v_p(J_\infty) - V_b)J_h/J_{h\infty}}, \tag{22}$$

in which J_h is given by the asymptotic solution provided by Eqs. (B4). The results of numerical integration of (22) and (21)

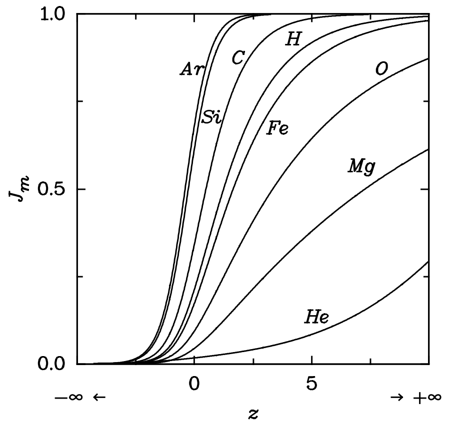

Fig. 4. The normalized photon fluxes J_m as functions of z using the asymptotic solution for the photon flux J_h which ionizes hydrogen.

are shown in Figs. (4 and 5) for J_m, n_m and N_m as functions of z within the layer, for various species.

The height z has been normalized to the hydrogen photoionization length and the curves show that species with larger photoionization cross sections than hydrogen (Ar, Al, Si, C) stand to left of hydrogen (H), whereas those with lower cross sections stand to the right (Fe, O, Mg).

Helium is an exception since the the photon flux (with $\lambda \sim 450$Å) which ionizes helium also ionizes hydrogen because its first ionization potential (FIP) is about double that of hydrogen. Therefore the roles of J_m (for helium) and J_h are reversed in Eqs. (18). However since we assume He is still a minor constituent and therefore its effect on the hydrogen equations is neglible this implies that the depletion of J_m for He takes place on its own photoionization length scale $(\alpha_{He} N_{He})^{-1}$ which is much greater than that of hydrogen. This effect is shown in Figs. 4 and 5 in which the He curves are shown standing well to the right of those of hydrogen. A maximum in the ion density n_m (for those ions with the larger photoionization cross sections) can arise because J_m attains its asymptotic value before J_h does, after which the acceleration of the flow takes over to reduce n_m to its final value ($N_b/2$ in the isothermal case $T_e = T_p = T_H$). This calculation should be regarded as merely illustrative because the situation is artificial in that it requires that we have just the right amount of photon flux J_m for all species. In reality there exists an excess of J_m in the form of $Ly \quad \alpha$ so the layer would not be stationary and in fact would move down (to the left), so that the metals would come in ionized from deep down. Furthermore since the ionizing of the metals is done by an enormous $Ly \quad \alpha$ flux it means we should regard the metals as coming into the layer fully ionized with the same speed as hydrogen and hence no fractionation can take place.

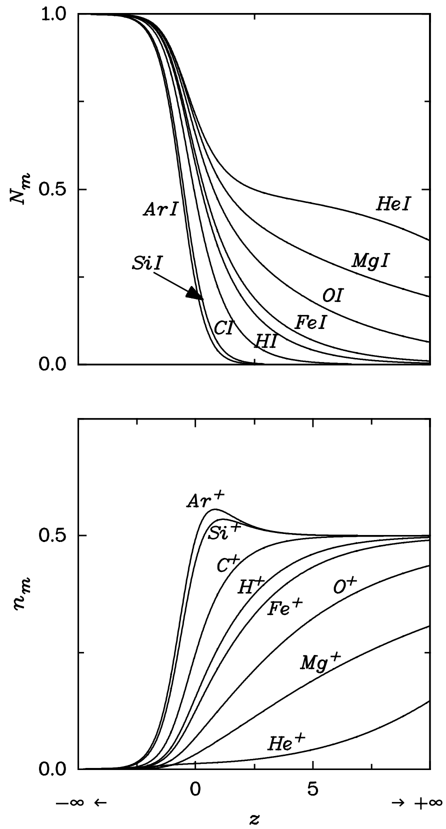

Fig. 5. The neutral (N_m) and ion (n_m) distributions with z for various species. Note that those ions with larger photoionization cross sections stand to the left of hydrogen whereas those with lower cross sections stand to the right. The exception is He since although its cross section is slightly greater than that of hydrogen its ionizing energy is about twice that of hydrogen with the result that it decays according to its own photoionization length which is much greater than of hydrogen.

5. Discussion

It is evident that the FIP effect, regardless of the ionization process, cannot be achieved in a steady-state situation in which collisions play a dominant role in locking species together. In the case of the fast solar wind we have demonstrated that steady flow with photo-ionization will not produce the required separation since there is no obvious way of removing or adding particles of a given species, whether neutral or ionized. There is however a possibility of accounting for the depression of the helium flux if helium remains substantially neutral well above the ionizing layer and is able to escape from the sides as neutrals, without replacement (e.g. von Steiger & Geiss (1989)) This is perhaps sufficient to account for the abundance variations observed in the fast wind which appear to involve helium only. In the slow solar wind the situation is altogether different since it is inherently non-steady, even as far as abundance variations are concerned. The slow wind appears to have its origin in the mainly closed coronal streamer belt and is a consequence of transient opening and closing of the magnetic field lines. The plasma which escapes to form the slow wind may therefore come from almost any altitude (allowing gravitational settling to play some role in differentiating between species) and also from a regime which exhibits a strong FIP effect even if apparently in hydrostatic equilibrium. In fact it seems likely that this

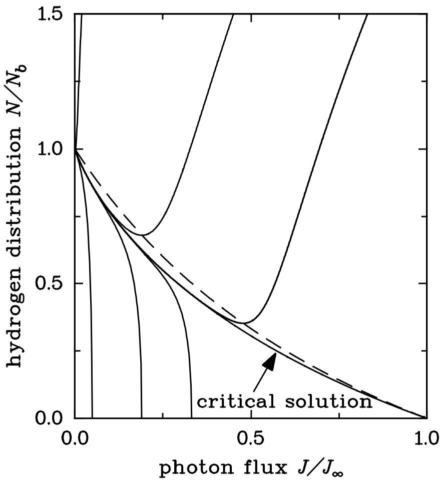

Fig. 6. The phase trajectories of the structure equation in the (N, J) plane. The broken curve is the locus of the minima of the integral curves. There is a unique critical solution which links the bottom of the layer ($N = N_b$, $J/J_\infty = 0$) to the top ($N/N_b = 1$, $J = 0$).

results primarily from magnetic separation of low FIP ions in partially ionized regions by upwards plasma drift as observed in the ionosphere of the Earth, especially in equatorial regions where the magnetic field is horizontal. Such a process preferentially injects low FIP ions but it must be inherently non-steady or spatially non-uniform since upwards motions are always associated with downwards motions if the plasma cannot escape directly into space.

Appendix A: the structure equation in the (N, J) plane

It is instructive to analyze the structure of the photoionization layer in the N (hydrogen), J (photon flux) plane. This is accomplished by dividing Eq. (7) by Eq. (1) and using (4), (5) and (8) to eliminate n, v, and V in terms of J and N so as to obtain.

$$\frac{dN}{dJ} = c_H \frac{[NJ - (J_\infty - J)(N_b - N)/v_1]}{NJ(N + (N_b - N)/v_1)}, \tag{A1}$$

where c_H is given by (12b) and v_1 is the speed ratio (14b). The point $N/N_b = 1$, $J = 0$ is a critical point of the differential equation and the phase trajectories in this neighbourhood can be approximated by the bundle

$$\frac{N}{N_b} - 1 = \frac{\overline{c}_H}{(\overline{c}_H/v_1 - 1)} \frac{J}{J_\infty} + c \left(\frac{J}{J_\infty}\right)^{\overline{c}_H/v_1}, \tag{A2}$$

where c, the constant of integration, generates the family of curves emanating from ($N/N_b = 1$, $J = 0$). The phase trajectories can exhibit a minimum ($dN/dJ = 0$) on the locus corresponding to the zero of the numerator of (A1). Phase trajectories lying above this locus have $dN/dJ > 0$ whereas below it $dN/dJ < 0$. The point $N = 0$, $J = J_\infty$ is also a critical point of the differential equation for which there is one critical solution which permits N to smoothly go to zero according to

$$\frac{N}{N_b} = m\left(1 - J/J_\infty\right), \tag{A4a}$$

$$m = \frac{v_1 \overline{c}_H}{2}\left(\sqrt{1 + 4/\overline{c}_H v_1^2} - 1\right), \tag{A4b}$$

The phase trajectories are shown in Fig. (6). The structure of the layer is described by the critical solution which is the only phase trajectory which links the bottom of the layer ($N = N_b$, $J = 0$) to the top of the layer ($N = 0$, $J = J_\infty$). In Appendix (B) we show that the locus $dN/dJ = 0$ (given by the broken curve on Fig. (6)) provides a useful asymptotic form for the critical solution.

Appendix B: asymptotic solution for layer structure for large c_H

If c_H, or more precisely its normalized version given by (16), is very much greater than unity the critical solution in the (N, J) plane is well approximated (to O($\overline{c}_H^{-1}$)) by the zero of the numerator of (A1), i.e.

$$\frac{N}{N_b} = \frac{1 - J/J_\infty}{1 + (v_1 - 1)J/J_\infty}. \tag{B1}$$

Physically this corresponds to the situation in which there are many hydrogen–proton collisions in a characteristic photoionization length so that the speeds are almost equal, i.e. $v \approx V$, and given by

$$v = V = V_b(1 + (v_1 - 1)J/J_\infty). \tag{B2}$$

Hence the proton density is given by

$$n = J/V_b(1 + (v_1 - 1)J/J_\infty). \tag{B3}$$

Using (B1) to eliminate N in favour of J in the photon flux decay Eq. (1) facilitates its exact integration

$$\frac{J/J_\infty}{(1 - J/J_\infty)^{v_1}} = \exp z/L. \tag{B4}$$

where $L^{-1} \equiv \alpha N_b$. Eqs. (B1) to (B4) provide a neat analytic form for the structure of the ionization layer. Eq. (B4) shows that at the lower boundary ($z = -\infty$) the photon flux decays exponentially to zero with the photonionization length scale L, whereas it approaches its asymptotic value J_∞ at the upper boundary ($z = +\infty$) with that length scale augmented by the speed ratio v_1. This asymptotic solution is shown in Figs. (3).

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