The chromosphere and corona of Procyon
(α CMi, F5 IV–V)

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Summary. Observations of Procyon (α CMi), an F5 IV–V star, have been obtained with the International Ultraviolet Explorer (IUE) satellite at low and high spectral resolution. Emission line fluxes are used to find the emission measure distribution. The range of likely electron pressure is discussed, based in part on consideration of line opacities. Models are made for several boundary values of the electron pressure. The chosen model has $P_e = 1.2 \times 10^{14} \text{cm}^{-3} \text{K}$ at $2 \times 10^5 \text{K}$ and $P_e = 3.6 \times 10^{14} \text{cm}^{-3} \text{K}$ at $10^4 \text{K}$, the latter being a factor of 3 greater than in the chromospheric model by Ayres, Linsky & Shine. The maximum temperature for a model with uniform emission is deduced to be $3 \times 10^5 \text{K}$, much lower than in the solar corona. Higher temperature and pressure models are allowed only if the EUV emission originates from limited areas of the stellar disc.

The radiation losses are shown to exceed the net conductive flux, and the energy input required can then be found without knowing the electron pressure. The non-thermal energy content is calculated from the widths of optically thin lines. The energy input is compared with the energy flux carried by acoustic, shock and Alfvén waves. Pure acoustic waves cannot account for the required energy flux. A combination of shock dissipation and Alfvén waves, of the type proposed by Osterbrock cannot be excluded.

1 Introduction

Two F5 stars, α CMi (F5 IV–V) and α Per (Fr Ib) have been observed using the IUE satellite as part of a programme to study the chromospheres and coronae in late-type stars. These are the earliest type of star in the programme, and represent bright examples of near main sequence and supergiant stars in the region where the onset of a sub-photospheric convection zone is expected. The present paper is concerned mainly with α CMi. This is a well studied star at optical wavelengths (Griffin & Griffin 1979) and models of the low chromosphere have been made from Ca II H and K profiles and fluxes (Ayres, Linsky & Shine 1974). Furthermore, prior to observations from IUE α CMi was known to show emission from the EUV resonance lines of H Lyα (1216 Å), Si III (1206 Å) and O VI (1032 Å), through work with
the Copernicus satellite (Evans, Jordan & Wilson 1975) and from C IV (1550 Å) and He II (1640 Å), from observations with the TD-I satellite (Jamar, Macon-Hercot & Praderie 1976). The Mg II resonance lines at ~ 2800 Å have been studied using balloon-borne spectrographs by Kondo et al. (1972) and with the Copernicus satellite by Morton et al. (1977).

These early observations were used by Evans et al. and by Jamar et al. to establish the distribution of the emission measure \( \int R N_e^2 dR \) with temperature over the range \( T_e \sim 2 \times 10^4-2 \times 10^5 \) K. Models based on two limiting values of the electron pressure were made by Evans et al., who concluded that the coronal temperature of α CMI must lie in the range 3 \( \times 10^5 \)-1.3 \( \times 10^6 \) K.

Extensive observations of α CMI have now been obtained with the IUE satellite. The present paper sets out the methods which may be used to analyse the fluxes and linewidths of α CMI and other near main sequence stars. These are applied to α CMI to give the emission measure distribution, the electron pressure and hence a model of the density and temperature as a function of height. The terms of the energy balance equation are then calculated. The profiles of several lines are analysed to give the non-thermal motions. A comparison is then made between the empirical energy input function and that expected from the dissipation of acoustic and Alfvén waves. A brief account of this work was given at the Second European IUE meeting (Brown & Jordan 1980). Other observations of F dwarf stars made with IUE (Saxner 1980) show similar emission lines, but no EUV emission lines were detected from α Car (F0 Ib) by Evans et al. (1975), or from α Per (F5 Ib) by Brown, Jordan & Wilson (1979).

2 Observations and data reduction

Observations of α CMI have been made at both high and low resolution in the long- and short-wavelength regions. Table 1 summarizes the exposures. The large aperture was used in all cases.

A wide range of exposure times is required at low resolution in the short wavelength region because of the rapid variation with wavelength of the continuum emission, and to obtain good exposures of both strong and weak lines. The high resolution spectra were obtained primarily to study the line profiles. In the short wavelength region the useful exposure time is limited by the level of scattered light from longer wavelengths.

In the long wavelength region the exposures were aimed at observing the Mg II emission components and most other parts of the spectra are too heavily exposed for quantitative work on the continuum or absorption lines.

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<th>Exposure time (min)</th>
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<td>1979 Sept. 28</td>
<td>LWR 5713</td>
<td>High</td>
<td>1.5</td>
</tr>
</tbody>
</table>

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The data have been obtained from the geometrically- and photometrically-corrected (GPHOT) \textit{IUE} images from which instrumental non-linearities should be absent. These images have a two-dimensional raster format of 768 \times 768 data values (pixels). Low dispersion images take the form of a single echelle order while high dispersion images consist of a number of orders. Reduction of low dispersion spectra is relatively simple and involves the extraction of a strip of the image parallel to the dispersion and centred on the spectrum. The spectrum and adjacent background strips are then scanned numerically into bins which are usually 6 pixels high and 1\(\frac{1}{2}\) pixels wide. The exact positioning of the spectrum for this process is checked by examining cross-sections of the image strip at prominent lines. Line fluxes are found by removing the background determined by averaging points each side of the line at about \(\pm 6\) \AA{} from line centre.

In the high resolution spectra each order or individual line has been handled as a separate spectrum in the manner used for low resolution spectra. Whereas the bin width at low resolution corresponds to roughly 2.5 \AA{}, at high resolution this width is 0.047 \AA{} being slightly less than 1\(\frac{1}{2}\) pixels wide. The corresponding instrumental resolutions are approximately 6 and 0.1 \AA{}.

Those spectra affected by the SWP intensity transfer function error (Holm 1979) have been corrected by reprocessing the GPHOT image using a routine supplied by Gondhalekar (1979, private communication).

The background is removed after addition of the two background strips and smoothing by a single pass running mean or double pass for the high resolution spectra. Prominent noise spots are removed before smoothing. A linear wavelength calibration is established from the positions of two known emission lines. Absolute flux calibration is achieved using the sensitivity curves of Bohlin & Snijders (1978) and Cassatella & Selvelli (1979) for low and high resolutions, respectively. At short wavelengths, roughly less than 1500 \AA{}, the calibration of Cassatella & Selvelli rises steeply to account for the increasing contamination of the inter-order background expected in sources with a strong continuum. However, for cool stars the background is weak and only a few emission lines are observed. We have therefore also used a linear extrapolation of the \(C_{\lambda}\) curve, taking the Cassatella & Selvelli value at 1600 \AA{} and \(C_{\lambda} = 140\) at 1200 \AA{}.

Low resolution spectra of the same star are added to give coverage at reasonable exposure levels for all wavelength regions. In this process all overexposed sections of spectrum are excluded. The addition can give any desired weight to each spectrum although in general each has been given a weight of unity except for particularly noisy spectra which have weights of a half. The combined spectrum is shown in Fig. 1. The resonance lines of H Ly\(\alpha\), C\(\text{II}\), C\(\text{IV}\), N\(\text{V}\), O\(\text{I}\), Si\(\text{III}\) and Si\(\text{IV}\) are present. The He\(\text{II}\) lines at 1640 \AA{} are weaker than suggested by Jamar \textit{et al.} (1976), possibly because absorption in C\(\text{I}\) at 1656 \AA{} influenced their choice of continuum level. C\(\text{I}\) absorption is also present at \(\sim 1560\) \AA{}.

At high resolution the same emission lines are observed and then Si\(\text{III}\) at 1206 \AA{} is clearly resolved from H Ly\(\alpha\).

No analysis has yet been made of the absorption spectrum. The blackbody curve for 4750 K, the minimum temperature in the model by Ayres \textit{et al.} (1974) is shown, but the continuum below 1650 \AA{} is contaminated by scattered light.

The line fluxes are given in Table 3. For C\(\text{II}\) and C\(\text{IV}\) fluxes from both low and high resolution spectra are available; the low resolution fluxes are slightly higher. However, the \textit{IUE} fluxes for Si\(\text{III}\), C\(\text{II}\) and H Ly\(\alpha\) are a factor of 2 to 3 higher than those found from the \textit{Copernicus} spectra (Evans \textit{et al.} 1975). On the other hand the C\(\text{IV}\) flux is a factor of \(\sim 2-4\) lower than that found by Jamar \textit{et al.} (1976) from the TD-1 satellite. It was difficult to extract the line flux from background in their spectra, because of the nearby C\(\text{I}\) absorption.
feature. Linsky & Marstal (1981) have discussed IUE spectra of α CMi which they obtained at low resolution using the small aperture. They have used the fluxes given in a previous short account of our work (Brown & Jordan 1980) to obtain the absolute flux calibration. The agreement between relative fluxes is satisfactory for the stronger lines, but less so for weaker features.

3 Analysis of line fluxes

3.1 Lines formed above 2 × 10⁴ K

Methods for analysing solar EUV line fluxes (Pottasch 1964; Jordan & Wilson 1971) can also be applied to stellar EUV spectra, and are summarized briefly here.

The flux $F_\lambda$ (in erg cm⁻² s⁻¹) at the star can be obtained from the flux at the Earth, $F_\odot$, using the measured angular diameter of 5.5 × 10⁻³ arcsec (Hanbury-Brown et al. 1967, for a limb-darkened photosphere; Blackwell & Shullis 1977, give 5.73 × 10⁻³ arcsec).

Assuming a spherically symmetric atmosphere, effectively thin emission lines, and a collisionally excited upper level leads to

$$F_\lambda = \frac{6.8 \times 10^{-22}}{\lambda} \frac{\Omega_{12}}{\omega_1} \frac{N_E}{N_H} \int \Delta h g(T) N_e^2 \, dh,$$

where

$$g(T) = T_e^{1/2} \frac{N_{\text{ion}}}{N_E} \exp \left(- \frac{W_{12}}{k T_e} \right)$$

and where $\Omega_{12}$ is the averaged collision strength (e.g. Seaton 1962), $\omega_1$ is the statistical weight of the lower level and $W_{12}$ is the excitation energy.

The abundances are taken as constant throughout the atmosphere. $N_H$ has been set equal to 0.8 $N_e$ for $T_e > 10^4$ K, and for purposes of illustration only $N_{1}/N_{\text{ion}}$ has been taken as
1.0. In practice the excitation model is examined individually for each line used. \( \Omega_{12} \) is taken as an average quantity appropriate to the expected temperature range.

Following Jordan & Wilson (1971), a logarithmic temperature with \( \Delta \log T = \pm 0.15 \text{ dex} \) is chosen as typical of the temperature range over which an emission line is formed. The percentage of total emission formed in this temperature range is calculated assuming \( \int_{T_1}^{T_2} N_\lambda^2 d\log T \) is constant over this region. Thus an average value of \( g(T) \) is found over \( \Delta \log T \). This procedure is carried out for all lines and allows for the different shapes of the \( g(T) \) curves. Then

\[
\int_{\log T_1}^{\log T_2} g(T) d\log T = G \cdot g(T_m)
\]

is computed to find \( G \) the normalization constant, where \( T_1 \) and \( T_2 \) are well outside \( 1.41 T_m \) and \( 0.71 T_m \), where \( T_m \) is the temperature at the peak of the \( g(T) \) function.

Initially one replaces \( g(T) \) by \( G \cdot g(T_m) \) and takes it outside the integral in equation (1), and hence \( F_\odot \) yields \( \int_{T_1}^{T_2} N_\lambda^2 d\log T \), the usual emission measure. Once the distribution of the emission measure with temperature is established using all available lines the quantity \( \int_{T_1}^{T_2} N_\lambda^2 g(T) d\log T \) can be recomputed, and a new value of \( T_m \) determined.

The ion populations may be derived from tabulations in the literature. Jordan (1969) made two sets of calculations — density independent and density dependent based on a solar model. Comparison with later results by Summers (1972) shows that the density effects were overestimated, and are appropriate to \( P_p \sim 10^{16} \text{ cm}^{-3} \text{ K} \), rather than the quiet sun pressure of \( \sim 6 \times 10^{14} \text{ cm}^{-3} \text{ K} \). Summers’ results have been used for C and O, and weighted means of Jordan’s calculations have been used for nitrogen and silicon. For the C ions Nussbaumer & Storey (1975) have examined the effects of photo-ionization and collisions from metastable levels; neither are important where C II is predominantly formed. Jacobs et al. (1977) have recalculated the ionization balance for silicon ions and the results for Si III and Si IV are essentially the same as those used here.

Table 2. Atomic data used.

<table>
<thead>
<tr>
<th>Ion</th>
<th>( \lambda ) (Å)</th>
<th>( \Omega )</th>
<th>( \log G \cdot g(T_m) )</th>
<th>( T_m ) (K)</th>
<th>( N_\lambda/N_{\lambda H} )</th>
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<td>1334.5</td>
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<td>2802.7</td>
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<td></td>
<td>1402.8</td>
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* For excitation only by collisions from ground term – see text for further discussion.
A. Brown and C. Jordan

Table 3. Observed fluxes and deduced emission measures.

<table>
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<th>Ion</th>
<th>$\lambda$ (Å)</th>
<th>$F_\Phi$ (erg cm$^{-2}$ s$^{-1}$) (low resolution)</th>
<th>$F_\Phi$ (erg cm$^{-2}$ s$^{-1}$) (high resolution)</th>
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<td>$4 \times 10^4$</td>
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<td>1335.7</td>
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<td>$3.3 \times 10^{-12}$</td>
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<td>C IV</td>
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<td>$1.2 \times 10^{-11}$</td>
<td>$5.3 \times 10^{-12}$</td>
<td>$1.0 \times 10^4$</td>
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<td>($\approx 7.0 \times 10^{-13}$)</td>
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<td>($2.8 \times 10^{-12}$)</td>
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<td>1402.8</td>
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<td>($1.3 \times 10^{-12}$)</td>
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* From low resolution fluxes except where only limits are available; in these cases from high resolution fluxes.
† From Copernicus data.

Analyses of photospheric spectra (Griffin 1971) show that the abundances in α CMi are essentially solar. For C II, C IV, O VI and Si IV, the impact parameter method of Burgess (1964) has been used to obtain values of $\Omega$. For Si III values calculated by Nicolas (1977) were adopted. At $N_e \sim 10^{10}$, Si III has over 90 per cent of its population in the ground state. The atomic data and resulting emission measures are summarized in Table 2 and Table 3. The emission measures are from the low resolution fluxes except where only high resolution fluxes are available. The emission measure distribution is shown in Fig. 2, including points based on all measured fluxes. The emission measure derived from the Copernicus observation of O VI is also included. The uncertainty in the average emission measure at a given temperature due to uncertainties in fluxes, abundances, atomic data and $g(T)$ will be about ±50 per cent.

3.2 Lines formed below 2 x 10$^4$ K

Below $T_e \sim 2 \times 10^4$ K a different procedure is used. The chromospheric model of Procyon made by Ayres et al. (1974), and based in part on the Mg II fluxes available at that time, is adopted up to 8000 K. The quantity $\int \Delta h N_e N_H dh$ is calculated from the model, where $\Delta h$ still refers to a region where $\Delta \log T = \pm 0.15$. $N_e N_H$ rather than $N_e^2$ is relevant since hydrogen is not fully ionized below $2 \times 10^4$ K. The resulting values of $\int \Delta h N_e N_H dh$ are shown in Fig. 2.

A simple two-level model for Mg II is used to derive $\int N_e N_H dh$ assuming all observed photons are created by collisions. This provides a cross-check on the model, which was based on earlier Mg II observations. The impact parameter method was used to find $\Omega$. $N(Mg II)/N(Mg)$ was taken as 1.0, from the ionization equilibrium calculations of Jordan (1969). Since the emission measure is sensitive to $T_e$ calculations were made at 6300 and 8000 K.
The values of $f N_e N_H \, dh$ derived are also shown in Fig. 2 and are consistent with the model if the line formation (i.e. photon creation) takes place predominantly at 7000 K.

Of the strong lines observed only H Ly$\alpha$ is formed between $10^4$ and $4 \times 10^4$ K. Because of interstellar absorption the observed H Ly$\alpha$ flux gives only a lower limit to the stellar emission. Assuming collisional excitation of H Ly$\alpha$, with the excitation rate used by Hearn (1967), leads to a corresponding lower limit to $f N_e N_H \, dh$, which is shown in Fig. 2. It is argued below that the intrinsic stellar H Ly$\alpha$ flux should be about a factor of three higher than that observed and the larger value of the emission measure is also indicated.

The values of $f N_e N_H \, dh$ which would be required to produce the observed O I emission through different excitation mechanisms can also be calculated and compared with the model by Ayres et al. Collisional excitation (with collision strengths from Rountree 1977) fails to produce enough emission by several orders of magnitude. Radiative recombination would be sufficient at 6300 K, provided the ionization equilibrium is dominated by charge exchange, as in the chromosphere of $\alpha$ Boo (Haisch et al. 1977). Excitation via photo-excitation by H Ly$\beta$ to the $^3D$ level in O I, followed by cascade through the resonance multiplet (Shine et al. 1976; Haisch et al. 1977) can be examined only in an approximate way since the H Ly$\beta$ profile and flux as a function of height in the chromosphere is not known. The H Ly$\beta$/H Ly$\alpha$ flux ratio is taken as $10^{-2}$, as in the Sun, and the observed lower limit to the H Ly$\alpha$ flux is used. The revised O I transition probabilities discussed by Christensen (1979) have been adopted for the decay routes. These approximations lead to values of $f N_e N_H \, dh$ lower than in the model, which is understandable since the actual excitation will tend to be less efficient than that assumed.
4 Linewidths and profiles

The high resolution spectra are illustrated in Figs 3–6. The causes of line broadening which should be considered are (i) thermal Doppler broadening, (ii) non-thermal motions which arise through the deposition of energy or passage of waves through the atmosphere, (iii) superposition of flows, (iv) scattering into the wings from the line centre due to a high opacity and (v) the instrumental resolution.

The instrumental resolution should have negligible effect on the Mg II lines and H Lyα. For other lines it has been allowed for widely by subtracting a Gaussian FWHM of 0.1 Å from the observed profile treated as a Gaussian. The noise level is a more significant source of error.

4.1 Lines formed at $T_e > 2 \times 10^4$ K, and which may be optically thin

Initially the widths are examined allowing for contributions (i) and (ii) but not (iii). Then one may put

$$\frac{\Delta \lambda}{\lambda} = 7.1 \times 10^{-7} \left( \frac{T_e}{m_i} + \frac{\xi_0^2 m_p}{2k} \right)^{1/2}$$

or

$$\xi_0^2 = 3.3 \times 10^{20} \left( \frac{\Delta \lambda}{\lambda} \right)^2 - 1.66 \times 10^8 \frac{T_e}{m_i},$$

where $\Delta \lambda$ is the FWHM, $m_i$ is the ion mass relative to the proton mass, $m_p$, and $\xi_0$ is the derived ‘non-thermal’ most probable speed. The rms value of $\xi_0$ is given by $\langle V_T^2 \rangle = 3/2 \xi_0^2$ provided the motions have a Maxwellian distribution.

The speeds can be compared with the sound speed given by

$$C_S = (\gamma P/\rho)^{1/2}.$$  

For the solar molecular weight at $T_e \geq 2 \times 10^4$ K and with $\gamma = 5/3$

$$C_S = 1.46 \times 10^4 T_e^{1/2} \text{ cm s}^{-1}.$$  

In Table 4 we give the observed values of $\Delta \lambda$, from the 90 min exposure, the adopted value of $T_e$ and the calculated values of $\xi_0$, $C_S$ and $\langle V_T^2 \rangle$.

<table>
<thead>
<tr>
<th>Ion</th>
<th>$\lambda$ (Å)</th>
<th>$\Delta \lambda$ (Å)</th>
<th>$T_e$ (K)</th>
<th>$\xi_0$ (km s$^{-1}$)</th>
<th>$\langle V_T^2 \rangle^{1/2}$ (km s$^{-1}$)</th>
<th>$C_S$ (km s$^{-1}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mg II</td>
<td>2795.5</td>
<td>0.18*</td>
<td>&gt;6.5(3)</td>
<td>11</td>
<td>14</td>
<td>8.4</td>
</tr>
<tr>
<td>C II</td>
<td>1335.71</td>
<td>0.39</td>
<td>4.0(4)</td>
<td>53</td>
<td>65</td>
<td>26</td>
</tr>
<tr>
<td></td>
<td>1335.66</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Si III</td>
<td>1206.5</td>
<td>0.43</td>
<td>5.0(4)</td>
<td>64</td>
<td>79</td>
<td>29</td>
</tr>
<tr>
<td>Si IV</td>
<td>1393.8</td>
<td>0.18</td>
<td>7.4(4)</td>
<td>22</td>
<td>28</td>
<td>37</td>
</tr>
<tr>
<td>C IV</td>
<td>1548.2</td>
<td>0.39</td>
<td>1.0(5)</td>
<td>44</td>
<td>54</td>
<td>46</td>
</tr>
<tr>
<td></td>
<td>1550.8</td>
<td>0.21</td>
<td>(22)</td>
<td>(26)</td>
<td>(26)</td>
<td></td>
</tr>
<tr>
<td>O VI</td>
<td>1031.9</td>
<td>&lt;0.31</td>
<td>2.8(5)</td>
<td>&lt;52</td>
<td>&lt;63</td>
<td>77</td>
</tr>
</tbody>
</table>

* See text.
† Upper limit from *Copernicus* data.

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The estimates of opacity later in this section show that the lines formed at highest temperature are least likely to be optically thick, so we begin the discussion with those lines.

Although the N V 1238 Å line is observed at high resolution it is not sufficiently strong to yield a reliable value of the FWHM.

The two lines of the C IV resonance multiplet have a peak intensity ratio of ~ 2:1 as expected under optically thin conditions. The linewidths are much greater than expected with the local electron temperature. The 1548 Å line is broader than the 1551 Å line, but given the noise level in the profiles this is not interpreted as significant. Both values are given in Table 4. The value of \( \langle V_2^2 \rangle \) derived from the 1548 Å line is close to the local sound speed and is higher than the value of 24 km s\(^{-1}\) found in the solar transition region (e.g. Brueckner & Moe 1973; Moe & Nicolas 1976; Boland et al. 1975, and later authors).

Only the stronger of the two components of the Si IV resonance multiplet can be used to derive a width; this is smaller than the C IV widths. The other component of the Si IV doublet is substantially weaker, consistent with the peak flux ratio of 2:1 expected under optically thin conditions.

The Si III profile is noisy and the FWHM gives a value of \( \xi_0 \) which is larger than that found from Si IV or C IV. From their *Copernicus* spectra, Evans et al. found that the Si III line had three adjacent points significantly above background giving an approximate width of 50 km s\(^{-1}\) (± 15 km s\(^{-1}\)), consistent with the present value of 64 km s\(^{-1}\).

The C II profiles can be obtained from the 30 and 90 min exposures, and these are illustrated in Fig. 3. Inspection of the original photo-write shows that the strong peak in the 1334.5 Å line in the 90 min exposure (dotted line) is caused by a small noise spot. Apart from this one would expect the 90 min exposure to give more accurate profiles but the noise level is still quite high. However, the three other profiles on the photo-write show clear evidence for self-reversal. The line at ~ 1335.7 Å is a close blend of the \( ^2P_{3/2} - ^2D_{3/2} \) and \( ^2P_{3/2} - ^2D_{5/2} \) components at 1335.66 Å and 1335.71 Å respectively. The ratio of the 1334.5: 1335.7 peak intensities should be 1:2 whether the lines are optically thin or thick. The shape of the lines makes it difficult to define a FWHM. Averaging across the peaks of the 1335.7 blend leads to a FWHM of 0.39 Å. The apparent self-reversal of the 1335.7 line indicates a high opacity and then the measured width will overestimate the non-thermal motions.

![Figure 3. Emission profiles of C II, Si III, Si IV and C IV resonance lines. All profiles are from a 90 min exposure except the solid line profile of C II which is from a 30 min exposure.](https://academic.oup.com/mnras/article-abstract/196/4/757/983083)
4.2 Lines which are likely to have high opacity

The O I lines are shown in Fig. 4. The spectra are noisy and the definite contributions from O I are indicated by shading. These emission components have intensity ratios of 9:5.5:2.5, close to the optically thin ratio of 9:5.4:3, and definitely not in the ratio of 1:1:1 expected with high opacity. However, each complete O I line consists of the obvious blue wing emission plus weaker emission at longer wavelengths beyond an absorption component. The O I lines are formed in essentially the same region of the atmosphere as the Mg II lines and as seen from Fig. 5 these are also highly asymmetric, with a strong blue wing enhancement. If, as indicated above, the C II lines are optically thick then the O I lines should also be optically thick, and the model by Ayres et al. would indeed lead to \( \tau \sim 500 \) in the strongest O I line. Given the complexity of the O I line formation, it is clear that a full radiative transfer calculation, including velocity fields is required to interpret both the O I and Mg II profiles. The strong blue asymmetry does not apparently continue to the height of C II. This behaviour is consistent with an atmosphere which is decelerating rather than accelerating, only the latter indicating mass-loss. This is obvious from the following considerations. Photons in the red wing of a line created at some depth in the chromosphere, when reaching greater heights are reabsorbed closer to the local line centre, whereas photons in the original blue wing are reabsorbed further from the line centre. Thus the blue wing becomes enhanced.

Figure 4. High resolution spectrum of O I resonance triplet. The definite O I emission is indicated by shading.

Figure 5. Emission components of Mg II doublet. The spectra are taken directly from traces provided by ESA VILSPA reduction procedures. The dashed and solid lines represent data from different orders.
relative to the red wing. The degree of enhancement depends on the variation of the source function with depth. The photons in the blue wing reflect conditions in the region of last scattering and it is reasonable that the ratios of the blue components are close to their optically thin values. A decelerating downflow, as suggested for the solar supergranulation regions, cannot be excluded as an alternative cause of the asymmetric profiles.

### 4.2.1 The Mg II profiles

The Mg II profiles are shown in Fig. 5. These have been used to derive the flux in the Mg II emission lines, the width of the lines at their base, and the separation of the emission components. The profiles are strongly asymmetric with a blue wing enhancement. A small interstellar absorption component could be present at the rest wavelengths.

The total flux in the lines above an interpolated background is $1.7 \times 10^6 \text{erg cm}^{-2} \text{s}^{-1}$; the $k$ line flux is $1.1 \times 10^6 \text{erg cm}^{-2} \text{s}^{-1}$. These can be compared with previous values in the literature; e.g. Linsky & Ayres (1978) give a total of $3.5 \times 10^6 \text{erg cm}^{-2} \text{s}^{-1}$, based on observations with the BUSS experiment, assuming zero central intensity for the broad absorption component. Making the same assumption our spectra result in a total flux of $2.8 \times 10^6 \text{erg cm}^{-2} \text{s}^{-1}$. Weiler & Oegerle (1979) give a lower limit for the $k$ line as $> 4.2 \times 10^5 \text{erg cm}^{-2} \text{s}^{-1}$. In an earlier paper Ayres et al. (1974) found a $k$-line flux integrated over 1 Å, of $8.5 \times 10^5 \text{erg cm}^{-2} \text{s}^{-1}$, based on observations made by Kondo et al. (1972). The present value lies between earlier estimates.

The width of the ‘base’ of the $k$ line is 1.41 Å, leading to $W = 151 \text{ km s}^{-1}$ in good agreement with the previous measurement of 1.5 Å by Kondo et al. (1972). This width is in excellent agreement with that expected from the empirical ‘Wilson–Bappu’ relations given by Kondo, Morgan & Modisette (1976a, b). For $M_{\nu} = +2.65$ these predict $W = +2.10$ and $+2.17$ respectively, compared with the present value of 2.18.

The separation of the emission peaks, $(2W_2)$, is 0.62 Å. If it is assumed that the emission peaks are formed near the optical depth, $\tau \sim 1$, then $W_2$ can be combined with the measured flux $F(\text{Mg II})$ to give $\Delta \lambda_{T}$ (the non-thermal 1/2–1/e width) and $\tau_0$ the opacity at line centre.

It is assumed that $W_2$ lies in the Doppler core of the line. Then

$$W_2/\Delta \lambda_{T} = (1n \tau_0/2)^{1/2}$$

or

$$W_2/\Delta \lambda = (1n \tau_0/2)^{1/2}/1.66,$$

where $\Delta \lambda$ is the FWHM, as used in equation (8). From the absorption coefficient for a Doppler broadened line,

$$\tau_0 = 1.2 \times 10^{-14} \lambda(\text{Å}) f_{12} m_i^{1/2} \frac{N_E}{N_H} \int_{\Delta h} \frac{N_i}{N_E} \cdot \frac{N_i}{N_H} \cdot \frac{N_H}{T_{1/2}^{1/2}} dh$$

(7)

where

$$T_{1/2}^{1/2} = \Delta \lambda m_i^{1/2}/\lambda 7.1 \times 10^{-7},$$

$f_{12}$ is the oscillator strength and $m_i$ is the atomic weight. But the total flux $F(\text{Mg II})$ is given by equation (4), which contains similar quantities. Replacing $\Omega$ in equation (4) by $\Omega f$ (Gaunt factor) one finds

$$\tau_0/F(\text{Mg II}) = 6.1 \times 10^{-6} \lambda^2(\text{Å}) T_e^{3/2} \exp (W/k T_e)/\Omega f P_e \Delta \lambda.$$

(8)

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$W_2$ and $F$ (Mg II) can be measured. For Procyon $P_e$ is known as a function of $T_e$ from the model by Ayres et al. (1974). Thus $\tau_0$ and $\Delta \lambda$ can be found from equations (9) and (11). The value of $\Delta \lambda$ is obviously insensitive to $\tau_0$. The resulting values vary between $2.6 \times 10^3$ and $7.6 \times 10^3$ depending on $T_e$ and $P_e$ and agree with that found by integrating the opacity in the Ayres et al. model down to $T_e \sim 6200$ K, where $\tau_0 \sim 2.4 \times 10^3$. The value of $\Delta \lambda$ is then 0.18 Å, leading to $\xi_0 = 11.5$ km s$^{-1}$. An integration of the Ayres et al. model down to $\sim 6200$ K is also required to match the observed flux in the $k$ line, and most of the emission would be formed near 6500 K. The present $k$ line flux observed is about twice that used in the derivation of the model by Ayres et al., which may therefore need some modification. The value of $\xi_0 = 11.5$ km s$^{-1}$ is larger than the microturbulent velocity of 5 km s$^{-1}$ assumed in the model by Ayres et al. at $T_e \sim 6500$ K. It is larger than the local sound speed ($c_s \sim 8.5$ km s$^{-1}$). From $\Delta \lambda$ one finds $W_2/\Delta \lambda_T \sim 2.8$ consistent with the emission peaks being just within the Doppler core. The same method applied to the observations discussed by Lemaire & Skumanich (1973) (quiet Sun average) leads to $\xi_0 = 5.6$ km s$^{-1}$, compared with a value of $\xi_0 = 6.6$ km s$^{-1}$ deduced by Athay (1976). The profiles of the solar Mg II lines also show blue wing enhancements but these are less pronounced than in Procyon.

4.2.2 The hydrogen Lyman $\alpha$ line

The H Ly$\alpha$ line of Procyon has been observed previously by Evans et al. (1975) using the Copernicus satellite at low resolution, and by Anderson et al. (1978) using the Copernicus satellite at high resolution. The present IUE observations have lower resolution but are an improvement in that they extend over a much wider wavelength range. Fig. 6(a) and (b) show the profile obtained from the 90 min exposure and the high resolution data from the Copernicus satellite, from Anderson et al. Both sets of data show absorption near the wavelength of D Ly$\alpha$ at $\sim 1215.34$ Å. The observed profiles of H Ly$\alpha$ are in broad agreement, but the geocoronal contribution is far less important in the IUE spectra.

Neither the intrinsic stellar line profile nor the absolute wavelength scale are known, and since the stellar Ly$\alpha$ may, like the Mg II lines, be strongly self-reversed and asymmetric it is difficult to estimate the contribution from interstellar absorption.

If the wavelength of the geocoronal component is adopted as the rest wavelength then the wavelength of the centre of the absorption and of the total emission and the wavelength of D Ly$\alpha$ can be measured with respect to this. The results are given in Table 5.

This procedure results in an average shift of $+0.04 \pm 0.03$ Å for the interstellar H Ly$\alpha$ and D Ly$\alpha$ absorption. The centre of the stellar emission is then blue-shifted by $-0.06 \pm 0.02$ Å. Anderson et al. also adopted the geocoronal H Ly$\alpha$ as the rest wavelength. They argued that the predicted heliocentric bulk velocity of interstellar hydrogen in the direction of Procyon is $18$ km s$^{-1}$, which with a stellar radial velocity of $-3$ km s$^{-1}$ should lead to $21$ km s$^{-1}$ for the apparent shift of the interstellar absorption components. The D Ly$\alpha$

<table>
<thead>
<tr>
<th>Table 5. The H Ly$\alpha$ profile.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Parameter</td>
</tr>
<tr>
<td>$\lambda$ of geocoronal H Ly$\alpha$ — adopted as $1215.67$ Å</td>
</tr>
<tr>
<td>$\lambda$ of centre of H Ly$\alpha$ absorption $1215.72 \pm 0.02$ Å</td>
</tr>
<tr>
<td>$\lambda$ of centre of D Ly$\alpha$ absorption $1215.37 \pm 0.03$ Å</td>
</tr>
<tr>
<td>$\Delta \lambda$ of H Ly$\alpha$ absorption $+0.05 \pm 0.02$ Å</td>
</tr>
<tr>
<td>$\Delta \lambda$ of D Ly$\alpha$ absorption $+0.03 \pm 0.03$ Å</td>
</tr>
<tr>
<td>$\lambda$ of centre of stellar emission $1215.61 \pm 0.02$ Å</td>
</tr>
<tr>
<td>$\Delta \lambda$ of centre of stellar emission $-0.06 \pm 0.02$ Å</td>
</tr>
</tbody>
</table>

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The chromosphere and corona of Procyon

Figure 6. (a) H Ly\(\alpha\) emission profile from 90 min \textit{IUE} exposure, (b) Central region of Ly\(\alpha\) profile as observed by Anderson \textit{et al.} (1978) using the Copernicus satellite.

should appear 1215.42 Å, and H Ly\(\alpha\) at 1215.75 Å. Our measured values are smaller, 1215.37 ± 0.03 Å and 1215.72 ± 0.02 Å, corresponding to a Doppler velocity of only 10 ± 6 km s\(^{-1}\). The blueshift of the stellar emission as a whole corresponds to a velocity of 15 ± 5 km s\(^{-1}\), suggesting that the intrinsic stellar line has a blue-enhanced wing similar to that seen in the Mg \(\text{II}\) lines.

The equivalent width of the interstellar absorption is difficult to determine and Anderson \textit{et al.} find this by calculating a line profile as a function of column density allowing for Doppler and Lorentzian broadening. This is not attempted here and their criticism that the previous determination of the interstellar absorption by Evans \textit{et al.} (1975) is likely to be an underestimate is accepted. The geocoronal emission fills in the absorption core and in order to account for the great width of the absorption that is observed, a Lorentzian component and a zero intensity core would be required, otherwise a very large velocity dispersion would be required in the interstellar region. Anderson \textit{et al.} find \(n_H \sim 0.11\) cm\(^{-3}\), a factor of \(~3.7\) larger than the upper limit given by Evans \textit{et al.} This results in an equivalent width, \(W\),
of 0.79 Å (through \( N_H = 1.9 \times 10^{18} \text{ cm}^{-2} \), where \( N_H \) is the column density). Approximate reconstruction of the stellar profile suggests that the intrinsic emission is about a factor 2.5 stronger than observed. However, in view of the highly asymmetric Mg\,\textsc{ii} profiles we consider that until a new model of the chromosphere can be constructed, linking the Mg\,\textsc{ii} forming region to that above \( 2 \times 10^4 \text{ K} \), the interstellar H Ly\( \alpha \) absorption remains poorly determined.

### 4.3 Estimates of Opacity in Emission Lines

It was shown above that the opacity at line centre, \( \tau_0 \), can be related to the line flux through the emission measure and electron pressure. Equations (1) and (7) are combined together as in equation (8). The Doppler linewidth must also be known. Then the product \( \tau_0 P_e \) can be found for each line, and some results are given in Table 6. This approach was developed for the analysis of solar limb to disc line ratios (Burton et al. 1971).

The C\,\textsc{iv} lines show a peak flux ratio close to that expected under optically thin conditions. Thus to give \( \tau_0 < 1 \), we require \( P_e > 2.5 \times 10^{14} \text{ cm}^{-3} \text{ K} \), using the broader linewidth. The Si\,\textsc{iv} lines should then also be optically thin, as suggested by the line ratio.

Using the observed width for Si\,\textsc{iii} leads to \( P_e < 1.6 \times 10^{14} \text{ cm}^{-3} \text{ K} \), if the line is optically thick. But then some of the observed width will be due to the high opacity. To obtain a more realistic upper limit the non-thermal velocity is set equal to the sound velocity and then \( P_e < 4.3 \times 10^{14} \text{ cm}^{-3} \text{ K} \).

Similarly, the apparent width of the C\,\textsc{ii} 1335.7 Å line leads to \( P_e < 1.5 \times 10^{14} \text{ cm}^{-3} \text{ K} \) if the line is optically thick. But setting \( V_T = C_v \) would give \( P_e < 3.8 \times 10^{14} \text{ cm}^{-3} \text{ K} \).

To summarize it can be seen that requiring C\,\textsc{iv} to be optically thin and at the same time requiring C\,\textsc{ii} to be optically thick puts stringent limits on the electron pressure. It must be \( \geq 2.5 \times 10^{14} \text{ cm}^{-3} \text{ K} \) at \( 10^5 \text{ K} \) but \( \leq 3.8 \times 10^{14} \text{ cm}^{-3} \text{ K} \) at \( 4 \times 10^4 \text{ K} \). These values are used to guide the modelling below.

### 5 Models of the Structure

Once the emission measure distribution has been established the temperature and density structure can be derived using methods developed in earlier solar work (e.g. Jordan & Wilson 1971). The generalization to stellar atmospheres has been discussed by Brown et al. (1979) and Jordan (1980a).

The emission measure, derived from each line is

\[
E_m = \int_{\Delta h} N^2 d h ,
\]

which can be rewritten as

\[
dT/dh = P_e^2 / 1.4 E_m T_e ,
\]

(9)
assuming initially that $P_e$ and $dT/dh$ are constant over the region of line formation but not over the whole atmosphere. If $P_e$ is allowed to vary according to hydrostatic equilibrium then

$$\frac{dP_e}{dh} = -7.14 \times 10^{-9} P_e g_* / T_e,$$

(10)

where $g_*$ is the stellar gravity. Thus if $P_e$ is known at some $T_e$, and $F_m$ is known as a function of temperature the structure of the atmosphere is determined by equations (9) and (10). Ideally $P_e$ could be determined from density sensitive line ratios, but in Procyon no suitable lines are observed. Thus models must be made for several chosen values of $P_e$, based on the limits derived above from arguments concerning the opacity, and the limits discussed below.

A lower limit to $P_e$ can be found by assuming that the hottest line observed is formed in an isothermal corona over one density scale height. Then

$$P_{\text{min}} = 1.29 \times 10^{-4} (E_m g_* T_{\text{max}})^{1/2}.$$  

(11)

(A factor $1/(0.43)^{1/2}$ was inadvertently omitted from the similar expression given by Jordan 1980a.)

For $\alpha$ CMi, $g_*$, the surface gravity is $1.2 \times 10^4$ cm$^{-2}$ s$^{-1}$ (Allen 1973).

The hottest line observed in the spectrum of $\alpha$ CMi is that due to O VII, observed with the Copernicus satellite. With $T_{\text{max}} = 2.8 \times 10^5$ K, and $E_m = 4.8 \times 10^{-5}$, $P_{\text{min}} = 5.2 \times 10^{13}$ cm$^{-3}$ K.

If the N V line is used instead, then

$$P_{\text{min}} = 2.5 \times 10^{14} \text{ cm}^{-3} \text{ K},$$

the value found from requiring the C IV lines to be optically thin.

From equations (9) and (10) the pressure at any $T_e$ can be expressed explicitly in terms of the pressure $P_0$, at $T_0 = 2.5 \times 10^5$ K, and the emission measure.

Writing $dP_e/dh$ as $dP_e/dT_e$ and substituting for $dh/dT$ leads to

$$P_e^2 = P_0^2 + 2.4 \times 10^{-4} \int_{T_e}^{T_0} E_m T_e d \ln T_e.$$  

(12)

Either a functional form for $E_m$ or point by point values can be used.

Four values of $P_0$ were investigated, $P_0 = 2.5 \times 10^{14}$, $1.0 \times 10^{14}$, $5.0 \times 10^{13}$ cm$^{-3}$ K and zero. A value of $P_0 > 2.5 \times 10^{14}$ cm$^{-3}$ K would give pressures at $\sim 4 \times 10^4$ K which are larger than those imposed by the opacity constraints discussed above. With $P_0 = 0$, the emission measure imposes a minimum value of $P_e$, and $P_e$ is not sensitive to $P_0$ for $P_0 \leq 10^{14}$ cm$^{-3}$ K. The models of temperature, a function of height, found with $P_0 = 0$, $10^{14}$ cm$^{-3}$ K and $2.5 \times 10^{14}$ cm$^{-3}$ K are shown in Fig. 7. The pressures are given in Table 7. The model with $P_0 = 0$ is referred to as $P_{\text{min}}$.

Using the emission measure down to $10^4$ K the minimum electron pressure at that temperature is $3.5 \times 10^{14}$ cm$^{-3}$ K, a factor of three larger than that given in the model by Ayres et al. at 8000 K. The formulation above assumes that hydrogen is sufficiently ionized for the total gas pressure to be the sum of the proton and electron pressures, which are approximately equal. Below $10^4$ K this assumption is inappropriate, and at 8000 K $P_e = 1/3 P_g$ in the model of Ayres et al. $P_{\text{min}}$ is already $2.8 \times 10^{14}$ cm$^{-3}$ K by $4 \times 10^4$ K, above which the emission measure curve is well determined. Postulating concentration of the emission to small areas of the stellar disc would raise the emission measure and $P_{\text{min}}$. The flux calibration is the only likely source of the emission measure being overestimated. However, the difference of a factor of three in the pressure is not large given the difficulty of making chromospheric models.

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Figure 7. Variation of temperature with height for three values of the limiting pressure, $P_e$.

Because the pressure limits from the consideration of the opacities agree with the value of $P_{\text{min}}$ from treating the O vi emission as 'coronal', and because the present pressures are already greater than those in the chromospheric model, there is no evidence for material above $T_e \sim 2-3 \times 10^8$ K.

Observations in the X-ray region provide other limits on coronal emission measures and temperatures. Mewe et al. (1975) found for α CMi an upper limit to the luminosity in the 0.2–0.28 keV range of $< 10^{28}$ erg s$^{-1}$. Crudace et al. (1975) found $< 4 \times 10^{28}$ erg s$^{-1}$ at $10^7$ K and $< 2 \times 10^{28}$ erg s$^{-1}$ at $10^7$ K, so the upper limit from Mewe et al. is more stringent. Converting the $10^{28}$ erg s$^{-1}$ to an emission measure, using the radiation parameters adopted by Crudace et al., leads to an upper limit of $\int N_e^2 dh < 10^{17}$ at $\sim 1.2 \times 10^6$ K. From Fig. 2 it can be seen that this is comparable with the value at $\sim 2 \times 10^8$ K, and there is no evidence of an emission measure rising to give a hot corona. The pressure limits imposed by this emission measure are $\geq 4.7 \times 10^{14}$ cm$^{-3}$ K at $1.2 \times 10^6$ K and $\geq 6 \times 10^{14}$ cm$^{-3}$ K at $2.5 \times 10^5$ K. These are not consistent with the observed line profiles unless the EUV emission originates from limited regions of the stellar surface (see below).

Table 7. Models of the pressure.

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<th>Model (b) $P_e = 5.0 \times 10^{13}$ (10$^{14}$ cm$^{-3}$ K)</th>
<th>Model (c) $P_e = 1.0 \times 10^{14}$ (10$^{14}$ cm$^{-3}$ K)</th>
<th>Model (c) $P_e = 2.5 \times 10^{14}$ (10$^{14}$ cm$^{-3}$ K)</th>
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<td>3.1</td>
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<td>3.0</td>
<td>3.2</td>
<td>3.9</td>
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<td>3.6</td>
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<td>6.3</td>
<td>6.4</td>
<td>6.4</td>
<td>6.8</td>
</tr>
</tbody>
</table>
From Fig. 7 it can be seen that the chromosphere of α CMi extends to ~ 4500 km, compared with ~ 2000 km in the Sun. The rise to \( T_e \approx 10^5 \) K is less steep than in the Sun, but this is not apparent on the scale of the figure. For the low pressure model the temperature flattens off to a ‘coronal’ value of ~ \( 2.8 \times 10^5 \) K. Higher pressures would imply \( T_e > 3 \times 10^5 \) K. The steep temperature gradient arises solely from the emission measure and pressure used — thermal conduction is unimportant in the pressure range adopted.

6 The energy balance

In general the energy input, \( \Delta F_m \), must be balanced by radiation losses, \( \Delta F_R \), the net conductive flux, \( \Delta F_C \), and losses in a stellar wind \( \Delta F_W \). No useful limits can be placed on \( \Delta F_W \) and in the absence of evidence for a strong wind the term is neglected. Then

\[
\Delta F_m = \Delta F_R + \Delta F_C .
\]

(13)

The radiation losses are given by

\[
\Delta F_R = \int_{\Delta h} 0.80 N^2 \, P_{\text{rad}} \, d h ,
\]

(14)

where \( P_{\text{rad}} \) is the power loss for a chosen composition of plasma. The calculations by McWhirter, Thonemann & Wilson (1975) have been used.

The quantities \( \Delta \) refer to regions where the logarithmic temperature interval is \( \log T_e \pm 0.15 \) dex, corresponding to the interval over which line emission predominantly occurs. Then we can write

\[
\Delta F_R \approx 0.80 \, E_m \, P_{\text{rad}} .
\]

(15)

Since \( P_{\text{rad}} \) is determined only by the temperature it can be seen that between different stars, \( \Delta F_R \) will scale directly as \( E_m \). This allows one to say immediately that the radiation losses in Procyon are about a factor of 10 larger than from the Sun in a given temperature range.

\( \Delta F_R \) does not depend on knowledge of the electron pressure. The calculated values of \( \Delta F_R \) (for convenience over intervals of 0.10 dex) are given in Table 8. (Because the observed

<table>
<thead>
<tr>
<th>( \log T_e )</th>
<th>( \Delta F_R ) (( \text{erg cm}^{-2} \text{s}^{-1} ) over 0.10 dex)</th>
<th>( F_m = F_R ) (( \text{erg cm}^{-2} \text{s}^{-1} ))</th>
</tr>
</thead>
<tbody>
<tr>
<td>3.9</td>
<td>1.7(5)</td>
<td>2.4(6)</td>
</tr>
<tr>
<td>4.0</td>
<td>2.2(5)</td>
<td>2.2(6)</td>
</tr>
<tr>
<td>4.1</td>
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<td>2.0(6)</td>
</tr>
<tr>
<td>4.2</td>
<td>3.4(5)</td>
<td>1.8(6)</td>
</tr>
<tr>
<td>4.3</td>
<td>1.8(5)</td>
<td>1.4(6)</td>
</tr>
<tr>
<td>4.4</td>
<td>1.1(5)</td>
<td>1.2(6)</td>
</tr>
<tr>
<td>4.5</td>
<td>6.2(4)</td>
<td>1.1(6)</td>
</tr>
<tr>
<td>4.6</td>
<td>5.8(4)</td>
<td>1.1(6)</td>
</tr>
<tr>
<td>4.7</td>
<td>8.8(4)</td>
<td>1.0(6)</td>
</tr>
<tr>
<td>4.8</td>
<td>1.6(5)</td>
<td>9.2(5)</td>
</tr>
<tr>
<td>4.9</td>
<td>2.5(5)</td>
<td>7.6(5)</td>
</tr>
<tr>
<td>5.0</td>
<td>2.1(5)</td>
<td>5.1(5)</td>
</tr>
<tr>
<td>5.1</td>
<td>1.5(5)</td>
<td>3.0(5)</td>
</tr>
<tr>
<td>5.2</td>
<td>1.2(5)</td>
<td>1.5(5)</td>
</tr>
<tr>
<td>5.3</td>
<td>3.1(4)</td>
<td>3.1(4)</td>
</tr>
</tbody>
</table>
The energy transported by thermal conduction is given by

$$F_C = -\kappa T_e^{\frac{5}{2}} \frac{dT}{dh} \text{ erg cm}^{-2} \text{s}^{-1},$$

where over the range of interest $\kappa \sim 10^{-6}$ (Spitzer 1962).

The value of $dT/dh$ is given by equation (9) and depends on the pressure as well as the emission measure. Using the highest pressure model (c) with $P_0 = 2.5 \times 10^{14} \text{ cm}^{-3} \text{s}^{-1}$, the value of $\Delta F_C$ at $2 \times 10^5$ is $\sim 6 \times 10^4 \text{ erg cm}^{-2} \text{s}^{-1}$, less than the value of $\Delta F_R = 3 \times 10^4 \text{ erg cm}^{-2} \text{s}^{-1}$. At lower temperatures and lower pressures the net conductive flux is far less than the radiation losses. Thus the energy input has only to balance the radiation losses and the value of $\phi_m$ (erg cm$^{-2}$s$^{-1}$) above each $T_e$ can be found by summing $\Delta F_R$ and is also given in Table 8.

The power loss function below 8000 K was not calculated by McWhirter et al. An estimate of the energy input required between $\sim 6300$ and 8000 K can be made by summing the observed line fluxes for Ca II (1.9 $\times$ 10$^6$ – Ayres et al. 1974), and Mg II (1.7 $\times$ 10$^6$ – this work). This lower limit is $\sim 3.6 \times 10^6 \text{ erg cm}^{-2} \text{s}^{-1}$. Adding this to the value of $2.4 \times 10^6 \text{ erg cm}^{-2} \text{s}^{-1}$ radiated above 8000 K leads to the conclusion that $\sim 6.0 \times 10^6 \text{ erg cm}^{-2} \text{s}^{-1}$ must pass through the atmosphere at $\sim 6000$ K to account for all the radiation losses above. A larger Mg II flux would of course increase this estimate.

In Section 5 the observed linewidths were used to derive non-thermal velocities, and it was found that these are comparable with the local sound velocity. A comparison is now made between the values of $\phi_m$ deduced from the radiation losses and the energy input expected from the dissipation of either acoustic or Alfvén waves.

The formulation used previously for analysing solar non-thermal motions is adopted (e.g. Boland et al. 1975). The non-thermal energy density, $E$, is given

$$E = \frac{1}{2} \rho \langle V_T^2 \rangle \text{ erg cm}^{-3},$$

where $\rho$ is the gas density and $V_T$ is the non-thermal velocity. The energy flux is then

$$\phi_m = 2EC \text{ erg cm}^{-2} \text{s}^{-1},$$

where $C$ is the appropriate propagation velocity for a plane wave. The factor 2 arises from the equipartition of energy between the potential and kinetic energy in the wave. This simple model gives an upper limit to the energy carried by acoustic waves, since departures from plane waves will reduce the energy flux. A plane wave would also be appropriate in the presence of a strong magnetic field with propagation along flux tubes (Lighthill 1978). The sound speeds calculated from equation (5) are given in Table 4.

The non-thermal energy densities and acoustic flux derived using the measured linewidths and gas pressures from model (b) are given in Table 9, and are shown in Fig. 8 where they can be compared with the mechanical energy flux required to account for the observed radiation losses. The pressures in the different models are not very different for $T_e \leq 10^5 \text{K}$.

From Fig. 8 it appears that there is a significant difference between the energy flux deduced from the radiation losses and that implied by the width of the Si IV line. It is unlikely that the $F_m$ ($= F_R$) values are overestimated, unless the IUE calibration is incorrect. The calculations of $P_{\text{rad}}$ by McWhirter et al. (1975) are in fact lower than earlier ones such as those by Cox & Tucker (1969). The calculated fluxes depend on $P_e$, but if the pressure were increased then the discrepancy between the present values at $\sim 10^4 \text{K}$ and those in the model by Ayres et al. would be even larger. One is forced to the conclusion that the sound
Table 9. Non-thermal energy densities and acoustic flux.

<table>
<thead>
<tr>
<th>Ion</th>
<th>λ (Å)</th>
<th>E (erg cm⁻³)</th>
<th>φ (erg cm⁻³ s⁻¹)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mg II*</td>
<td>2795</td>
<td>&lt;1.5</td>
<td>2.6(6)</td>
</tr>
<tr>
<td>C II</td>
<td>1335.7</td>
<td>&lt;0.21</td>
<td>&lt;1.1(6)</td>
</tr>
<tr>
<td></td>
<td>1335.6</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Si III</td>
<td>1206.5</td>
<td>&lt;0.24</td>
<td>&lt;1.4(6)</td>
</tr>
<tr>
<td>Si IV</td>
<td>1393.8</td>
<td>0.018</td>
<td>1.3(5)</td>
</tr>
<tr>
<td>C IV</td>
<td>1548.2</td>
<td>0.044</td>
<td>4.0(5)</td>
</tr>
<tr>
<td></td>
<td>1550.8</td>
<td></td>
<td></td>
</tr>
<tr>
<td>O VI†</td>
<td>1031.9</td>
<td>&lt;0.009</td>
<td>&lt;1.4(5)</td>
</tr>
</tbody>
</table>

* ρ from model by Ayres et al., other values from model (b).
† From Copernicus data.

velocity is not appropriate for converting the observed energy density to a propagating energy flux.

It is also useful to compare \( F_R \) with the maximum energy carried by a pure acoustic wave with \( V_T = C_s \). Then

\[
\phi_m = \rho C_s^3 \text{ erg cm}^{-2} \text{ s}^{-1}.
\]

(19)

This is shown by the dashed line in Fig. 8. It can be seen that the energy flux at \( 10^4 \text{K} \) is too small to account for the radiation losses above. In a previous brief account of this work (Brown & Jordan 1980) comparisons were made only around \( T_e \approx 10^5 \text{K} \), where the discrepancies are not so large, and the agreement between the acoustic flux and radiation losses appeared satisfactory.

The energy carried in the region where the Mg II lines are formed can also be examined. Using the values of \( P \) and \( \rho \) from the model of Ayres et al. \( \phi_m \sim 2.6 \times 10^6 \text{erg cm}^{-2} \text{ s}^{-1} \). Although this could account for the Mg II emission the value of \( \phi_m \) drops so rapidly above 6300 K, due to the rapid decrease of \( \rho \) that there is insufficient energy left to account for the total of \( \sim 2 \times 10^6 \text{erg cm}^{-2} \text{ s}^{-1} \) above \( 10^4 \text{K} \). Although the emission measure distribution is uncertain between 8000 K and \( 4 \times 10^4 \text{K} \) it can be seen from Fig. 8 that even with \( \text{no} \) emission in this region, an energy flux of \( \sim 10^6 \text{erg cm}^{-2} \text{ s}^{-1} \) would be required to account for the losses above \( 4 \times 10^4 \text{K} \).

![Figure 8. Comparison of the mechanical flux input and the maximum energy available from pure acoustic waves.](https://academic.oup.com/mnras/article-abstract/196/4/757/983083/983083)
It can be concluded that there are severe problems in attributing the mechanical energy deposition to the damping of acoustic waves. Neither the total radiation losses nor the linewidths converted to fluxes are consistent with this mechanism. Similar problems have been encountered in trying to attribute heating of the solar atmosphere to the dissipation of acoustic waves. (Athey & White 1978; Bruner 1978; Jordan 1976, 1980b).

The above arguments do not, however, exclude heating by shock waves, and most predictions of stellar coronal structure have been made on the basis of shock dissipation of acoustic waves generated in the convective zone (e.g. Kuperus 1965; Ulmschneider 1967; de Loore 1970). These estimates suggest that ample energy is available to account for the energy losses of the low and middle chromosphere (e.g. Ca II, Mg II, H-). de Loore predicts that for a star like α CMi, with $T_{\text{eff}} \sim 6510$ K, (Code et al. 1976) an energy flux of $7 \times 10^7$ erg cm$^{-2}$ s$^{-1}$ is available from the convective zone, compared with $1.4 \times 10^7$ erg cm$^{-2}$ s$^{-1}$ for the Sun. Schmitz & Ulmschneider (1980) have recently calculated the energy losses expected for α CMi up to $T_e \sim 6000$ K, from the model by Ayres et al., and have compared these with energy dissipated through the damping of shock waves. At the height where shocks are formed (below the temperature minimum) they find $\sim 10^8$ erg cm$^{-2}$ s$^{-1}$, which is reduced to $\sim 2 \times 10^7$ erg cm$^{-2}$ s$^{-1}$ by the height of the temperature minimum. Although this is slightly lower than their estimate of the H$^-$ losses ($2.5 \times 10^7$ erg cm$^{-2}$ s$^{-1}$) it is sufficient to account for the Mg II and Ca II losses ($3.6 \times 10^6$ erg cm$^{-2}$ s$^{-1}$) and the radiation from the atmosphere above 8000 K ($2 \times 10^6$ erg cm$^{-2}$ s$^{-1}$). Schmitz & Ulmschneider did not model the higher layers or predict coronal conditions. According to de Loore, shock heating should lead to coronal pressures and temperatures which are higher than in the Sun, but this does not appear to be the case. Landini & Monsignori-Fossi (1973) used de Loore’s models to predict coronal X-ray fluxes. For α CMi, as Cruddace et al. point out, the value is a factor of 70 above their upper limit, and also far exceeds the lower upper limit found by Mewe et al. Similarly Ulmschneider (1967) proposed that the coronal temperature should be proportional to the acoustic flux.

Although shock heating of the low atmosphere cannot be excluded the predicted rate of dissipation does not apparently fit the observed radiation losses. This can also be deduced from Fig. 8. Following Kuperus (1965), the flux carried by a shock wave of Mach number $M$ is

$$
\phi_m = \frac{4}{3} \rho C_s^3 \frac{(M^2 - 1)^2}{(\gamma + 1)^2 M}.
$$

At $10^4$ K, setting $\phi_m = F_R$ leads to $M \sim 4.7$; at $4 \times 10^4$ K, $M \sim 4$ and at $10^5$ K, $M \sim 3.5$. Although these values cannot be ruled out below $T_e \sim 5 \times 10^4$ K the linewidths at $10^5$ K indicate $M \leq 1$, with little energy passing through to higher temperature regions.

Osterbrock (1961) proposed that shock dissipation of fast mode magneto-acoustic waves may heat the solar chromosphere with slow mode Alfvén waves becoming more important in the corona, the transfer of energy between modes occurring through colliding shocks. Since the theory of these processes in the context of even the solar atmosphere is still incomplete only a simple exploration of the likely relevance of these modes can be made for α CMi.

The Alfvén speed (Alfvén 1947) and energy flux can be found using

$$
C_A = B/(4\pi \rho)^{1/2} \text{ cm s}^{-1}
$$

and $\phi_A = 2 E C_A$. The strength of the magnetic field of Procyon is not known and only an upper limit of 180 Gauss is available (Borra & Landstreet 1973).

Using the observed values of $V_A^2$, the values of $V_A$ and $B$ required to match the observed radiation losses can be found. Equations (5) and (21) lead to the condition that $C_A > C_s$.
provided $B \geq 1 \text{G}$, sufficiently low to make MHD modes of interest. To give sufficient flux a maximum field strength of 7 G would be needed at $\sim 7 \times 10^4 \text{K}$ where Si iv is formed. On the basis of solar field strengths this is not unreasonably large. The relevant damping process would be viscous damping and short periods would be implied to give sufficiently small damping lengths as in the solar case (Jordan 1976).

Finally the assumption of a uniformly emitting atmosphere, made throughout the analyses must be examined, since in the Sun it is known that the EUV emission lines are around an order of magnitude stronger in the supergranulation boundaries than in cell centres.

Suppose that the emission comes solely from an area $a = X(T_e)A$, where $A$ is the whole disc area and $X(T_e) < 1$. In deriving the emission measure distribution all values must be increased by a factor $1/X(T_e)$. The radiation losses deduced will also increase by the same factor. For the same pressure, the temperature gradient and hence the conductive flux will be reduced. Thus the conclusion that the net conductive flux is unimportant compared with radiation remains valid. The estimate of the minimum pressure from equation (11) would be increased by $1/X(T_e)^{1/2}$. Using this value as $P_e$ in equation (9) would return $dT/dh$ to the original value. The pressure change between $10^8$ and $10^9 \text{K}$ (equation 12) would be greater, thus some increase in $dT/dh$ at low temperatures would occur.

The pressures estimated through the opacity arguments would increase by $1/X(T_e)$. Thus if $X(T_e)$ were substantially less than 1.0 at $10^9 \text{K}$ there would be a discrepancy of $\sqrt{X(T_e)}$ between these pressures and that deduced as $P_{\text{min}}$.

The agreement between the two methods of estimating the pressure is consistent with a coronal temperature of $3 \times 10^5 \text{K}$ and $X(T_e) = 1$ in this region. It does not rule out concentration to supergranulation boundaries at lower temperatures.

With $X(T_e) < 1$ at $2.5 \times 10^5 \text{K}$ models with larger pressure and temperature become possible. For example if the upper limit of coronal emission measure discussed above was a positive detection, and hence $P_e \sim 6 \times 10^{14} \text{cm}^{-3}\text{K}$ at $2.5 \times 10^5 \text{K}$, an area factor of $\leq 1/2$ would be required to give consistency with the C ii line profiles.

The estimate of energy densities and acoustic flux depends on $P_e \propto 1/X(T_e)^{1/2}$. Hence, since $\Delta F_R \propto 1/X(T_e)$ the problem in accounting for the radiation losses becomes worse if area factors are important, and our conclusions concerning the likely heating processes remain unaltered.

It can be seen that a positive detection of an X-ray flux with simultaneous measurement of $T_e$ would place strong constraints not only on the pressure but also the uniformity of the region emitting the EUV flux.

7 Conclusions

The emission line fluxes from α CMi allow the emission measure distribution to be established. Over the temperature range $\sim 10^4$–$10^5 \text{K}$ the emission measure is an order of magnitude larger than in the Sun. Models made as a function of $P_e$ at $2 \times 10^5 \text{K}$ show that for $P_0 < 5 \times 10^{14} \text{cm}^{-3}\text{K}$, the net conductive flux is less than the radiation losses. The lower limit to the pressure is found to be $5 \times 10^{13} \text{cm}^{-3}\text{K}$ if the O vi lines are formed in the 'corona', leading to $P_e \sim 2 \times 10^{14} \text{cm}^{-3}\text{K}$ at $10^5 \text{K}$. The peak line flux ratio and profiles also suggest $P_e \sim 2 \times 10^{14} \text{cm}^{-3}\text{K}$ at $10^5 \text{K}$. When combined with the emission measure distribution and hydrostatic equilibrium these pressures result in an electron pressure at $10^4 \text{K}$ which is about three times that in the model by Ayres et al.

If $P_e \leq 5 \times 10^{14} \text{cm}^{-3}\text{K}$ then the radiation losses must be balanced directly by dissipation of non-thermal energy. Since the radiation losses do not depend on the pressure neither does the estimate of the mechanical energy input. The EUV observations suggest a model in
which the ‘coronal’ temperature is only $\sim 3 \times 10^5$ K. This is consistent with published upper limits for X-ray fluxes (Mewe et al. 1975).

The linewidths lead to velocities which are comparable with the sound velocity, but a high opacity adds to the broadening of lines formed below $\sim 5 \times 10^4$ K giving apparently supersonic velocities.

Although at $\sim 10^5$ K the energy carried by and dissipated through acoustic waves could account for the line broadening and radiation losses, at lower temperatures a pure acoustic mode clearly fails to give sufficient energy to balance the radiation losses. Apart from Joule heating, which is not possible to check from these observations, two other processes cannot be excluded. Dissipation of shock-waves has been suggested as the cause of heating in the low chromosphere of $\alpha$ CMi (Schmitz & Ulmschneider 1980), and this may continue into the higher chromosphere. However, the coronal parameters predicted by shock-heating do not fit the observations. Viscous damping of short-period Alfvén waves is a possible cause of heating above $\sim 2 \times 10^4$ K. The sequence of processes, changing fast-mode shocks to Alfvén modes proposed by Osterbrock (1961) may well be appropriate.

The main uncertainties in the modelling of the structure and energy balance come from the weakness of the lines above $10^5$ K, and the high level of noise in line profiles. Owing to the scattered light from the strong continuum above 1700 Å it may be difficult to improve the high resolution observations by extending the exposure time beyond the maximum so far of 90 min, but this will be attempted in our future IUE observing time.

Acknowledgments

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