Astron. Astrophys. 319, 589–592 (1997)

ASTRONOMY AND ASTROPHYSICS

The origin of strong magnetic fields in circumstellar SiO masers

T.W. Hartquist[1] **and J.E. Dyson**[2,3]

[1] Max-Planck-Institut für extraterrestrische Physik, D-85740 Garching, Germany
[2] Department of Physics and Astronomy, The University of Leeds, Leeds LS2 9JT, UK
[3] Department of Physics and Astronomy, The University of Manchester, Manchester M13 9PL, UK

Received 27 November 1995 / Accepted 20 August 1996

Abstract. We suggest that the Parker instability, acting in the regions where shocks form in response to stellar pulsations, gives rise to the fragments in which SiO masers exist in the atmospheres of evolved stars. The production of maser fragments in this manner is compatible with the Alfvén speeds in them greatly exceeding their kinematic and thermal speeds. The magnetic structure in an evolved star's envelope probably results in momentum transfer to gas from grains subject to radiation pressure being widely distributed even to regions where grains have not yet formed; the consequences for mass loss of the magnetic redistribution of momentum transfer require exploration.

Key words: stars: AGB; post AGB – stars: magnetic fields – stars: mass loss

1. Introduction

Barvainis et al. (1987) have measured the circular polarization of the SiO maser emission from several evolved stars of various types and inferred magnetic field strengths in the masers to be typically in the range of 10 to 100 Gauss. Models of the SiO maser pumps yield estimates for n_H, the number density of hydrogen nuclei, and T, the temperature, of approximately $10^{10} - 10^{11}$ cm^{-3} and 1500 K respectively (e.g. Doel et al. 1995). Barvainis et al. (1987) noted that the magnetic pressure in an SiO maser exceeds the thermal pressure. Of perhaps even greater significance, the Alfvén speed in an SiO maser, given by

$$V_A \approx 1.0 \times 10^3 \text{ km s}^{-1} \left(\frac{B}{50G}\right) \left(\frac{n_H}{10^{10}\text{cm}^{-3}}\right)^{-1/2}, \quad (1)$$

greatly exceeds the speeds (of the order of 10 km s^{-1}) of SiO masers relative to their central stars. We suggest in this research note that the large ratio of maser Alfvén speed to kinematic speed implies that the maser outflow is triggered by the Parker instability.

Send offprint requests to: J.E. Dyson

In Sect. 2 we identify the site of the onset of the Parker instability in an evolved late type star with the layer in which disturbances driven by stellar pulsation steepen into shocks. Sect. 3 concerns the ionization structure in an envelope of such a star and the size scales of the modes that are most unstable to the Parker instability; it also contains comments relevant to maser heating by ambipolar diffusion and the viability of maser pumping arising due to ambipolar diffusion. Sect. 4 consists of speculations about consequences of the Parker instability picture of SiO maser formation for the dynamics of the outer regions of the stellar envelope.

2. The Parker instability in the shock region

A region with an Alfvén speed greatly exceeding the thermal speed and the kinematic speed can be formed by the passage of a shock propagating in a direction antiparallel to a gravitational field, sufficient radiative cooling behind the shock to allow the magnetic field to dominate the postshock pressure, and the onset of the Parker instability leading to buoyancy taking field lines above the shock layer as gas flows along the bowed field lines back down to the shocked layer. The flow of the gas down the field lines creates regions with high Alfvén speed.

The pulsation of an evolved late-type star drives a series of outward moving pressure waves through material in its envelope (e.g. Bowen 1988). The pressure waves develop into strong outward moving shocks behind which gas is heated and compressed; further compression of postshock gas occurs as radiative cooling takes place. During compression, the component of the magnetic field perpendicular to the shock's propagation direction (which is radially outward) increases in strength in proportion to the density; thus, behind each shock, the component of the magnetic field perpendicular to the star's radial gravitational field will be enhanced while the component of the magnetic field parallel to the gravitational field will remain constant, and great compression behind a shock will result in the postshock magnetic field being almost perpendicular to the gravitational field. In this work we assume that radiative cooling and compression result in the postshock magnetic pressure greatly exceeding the postshock thermal pressure. (A condition for the validity of this assumption is given directly before Eq. (4) below.)

The existence of a magnetic field that is nearly perpendicular to the gravitational field results in the medium behind a shock being subject to the Parker instability. The nonlinear development of the Parker instability will give rise to a low density halo having a magnetic field strength comparable to that in the shock layer and extending over a distance roughly equal to the stellar radius.

The magnetic pressure in a postshock region in which the thermal pressure is negligible is comparable to the ram pressure of the shock. In Bowen's (1988) models of Mira variable shocks, the shock speed, V_s, is 25–30 km s^{-1} and the shock forms in gas with $n_H \simeq 1 \times 10^{13}$ cm^{-3} (see also Cherchneff et al. 1992). Thus, the magnetic field produced behind the shock in a Mira variable will have a strength of

$$B \approx 70\,G \left(\frac{V_s}{30\,\mathrm{km\,s^{-1}}}\right)\left(\frac{n_H}{1\times 10^{13}\,\mathrm{cm^{-3}}}\right)^{1/2} \tag{2}$$

in good agreement with the field strengths measured for the SiO masers.

Though the Parker instability causes the field to rise high in a halo, one might think that very nearly all of the mass on a field line will run down to the shock layer as the field bows; this would be the case if only a single sized perturbation were to grow. However, Parker (1967) has shown that perturbations of a rather broad size distribution grow at nearly the maximum rate during the linear development of this instability. We suggest that the masers are associated with perturbations on scales somewhat shorter than the largest that grow rapidly. As a larger scale perturbation grew, much of the material along the field line did flow to the shock layer, but as a somewhat smaller scale perturbation, superposed on the larger one, developed, some of the material on the field line flowed into the region associated with that smaller scale perturbation. The continued growth of the larger scale perturbation caused the smaller scale local density maximum on it to move outwardly with respect to the shock layer. Of course, the action of the instability on smaller scales causes this density maximum to be turbulent.

Only a few numerical simulations of the 3-D development of the Parker instability have been preformed (Matsumoto & Shibata 1992; Matsumoto et al. 1993) and the picture described in the previous paragraph of the nonlinear development in a region where magnetic pressure dominates over thermal pressure must remain somewhat speculative, because those simulations are for media in which the thermal pressure is initially equal to the magnetic pressure, leading to the retardation of the linear growth of large wavenumber perturbations. There are certainly signs in Fig. 5 of Matsumoto et al. (1993) showing the magnetic field structure arising in the nonlinear regime for one run that localized troughs in the magnetic field exist at some places in the halo, as suggested above.

3. The ionization structure and the maser masses, temperatures, and pumping

The fractional ionization, $X_i \equiv n(e)/n_H$ where $n(e)$ is the electron number density, is at least about 1×10^{-6} in the post-shock layer because the temperature is high enough for nearly all sodium to be ionized. As described below in more detail, an evolved star with SiO masers may have a chromosphere emitting sufficiently in the ultraviolet to produce other ions including Si^+ by photoionization, and X_i may be 1×10^{-5} or larger in the post-shock layer.

We follow Parker (1967) in estimating the effective Reynolds number, R_m, due to ambipolar diffusion in the shocked gas and obtain

$$R_m \approx 1.6 \times 10^4 \left(\frac{M_*}{1.2 M\odot}\right)^{1/2}\left(\frac{R_s}{300\mathrm{R}_\odot}\right)^{-1}$$

$$\left(\frac{\Lambda}{1\times 10^{12}\mathrm{cm}}\right)^{3/2}\left(\frac{X_{i,P}}{10^{-5}}\right)\left(\frac{n_{H,P}}{1\times 10^{14}\,\mathrm{cm^{-3}}}\right)\left(\frac{V_{A,P}}{10\,\mathrm{km\,s}}\right)^{-2} \tag{3}$$

where M_* and R_s are the stellar mass and shock radius, the subscript P signifies post-shock values, and Λ is the density scale-height in the post-shock gas. Here $n_{H,P}$ is a very uncertain quantity since the pre-shock magnetic field strength is unknown and we cannot determine to what extent the post-shock gas must be compressed before the magnetic pressure in it is comparable to the shock ram pressure. (For the validity of the two assumptions that the presence of the upstream magnetic field does not affect the formation and propagation of shocks and that the post-shock pressure is dominated by the magnetic pressure, the component $B_\perp$ of the upstream field perpendicular to the shock propagation must lie in the range $B/\sqrt{2} > B_\perp \gg B(C_c/V_s)^2$, where B is given by Eq. (2) and C_c is the isothermal sound speed in the downstream cooled gas.) Because $n_{H,P}$ is uncertain, $V_{A,P}$ is also uncertain, as is Λ where

$$\Lambda \approx 1.4 \times 10^{12}\,\mathrm{cm}$$

$$\left(\frac{M_*}{1.2\,\mathrm{M}_\odot}\right)^{-1}\left(\frac{R_s}{300\,\mathrm{R}_\odot}\right)^{2}\left(\frac{V_{A,P}}{10\,\mathrm{km\,s^{-1}}}\right)^{2}. \tag{4}$$

Parker (1967) has given results (his Figure 1 and the discussion around it) for the linear growth rates of perturbations for cases when $R_m = 1 \times 10^4$; they show that for such a large magnetic Reynolds number, the development of the instability at wavenumbers comparable to the inverse of the scaleheight proceeds mush as it does in the ideal MHD case. Growth at nearly the maximum rate occurs for a very broad range of perturbations including those with inverse wavenumbers comparable to Λ. That maximum rate is $(g/\Lambda)^{1/2}$, where g is the gravitational field strength; note that typically this maximum rate is greater than the stellar pulsation rate. We would expect the nonlinear development of the instability to favour the survival of the largest scale rapidly growing perturbations as smaller fragments coalesce. The most massive fragments, and hence the most massive possible SiO masers, would have masses of roughly

$$M_{\max} \simeq \rho_P \Lambda^3 \simeq 3 \times 10^{-7} \mathrm{M}_\odot \left(\frac{M_*}{1.2\mathrm{M}_\odot}\right)^{-3} \left(\frac{R_s}{300\mathrm{R}_\odot}\right)^6 \left(\frac{B_P}{50G}\right)^2 \left(\frac{V_{A,P}}{10\mathrm{km\ s}^{-1}}\right)^4 . \quad (5)$$

VLBI mapping of SiO masers in R Cas suggest that the maser spots have diameters of 2 to 6×10^{13} cm (McIntosh et al. 1989) which, if spherical geometry is assumed, correspond to masses of about 4–100 $\times 10^{-8}$ $\mathrm{M}_\odot$ $(n_H/1 \times 10^{10}\ \mathrm{cm}^{-3})$ comparable to $M_{\max}$ given by (5).

The above discussion is based on the assumption that the magnetic field lines remain reasonably well coupled to the neutral flow on lengthscales comparable to those of the perturbations that become SiO masers. Barvainis et al. (1987) estimated the ambipolar diffusion timescales in the masers to be of the order of only 10^3s, much shorter than the maser lifetimes; in fact, the ambipolar diffusion speed associated with such a short timescale is comparable to the speed of light, a clearly unrealistic result. The reasons that Barvainis et al. (1987) obtained such a small estimate for the diffusion time is that they assumed $n_H \approx 1 \times 10^9$ cm^{-3} and a fractional ionization of only 10^{-8} which would be appropriate for that value of n_H if the only source of ionization were low energy cosmic rays. In fact, ion-neutral collisions will induce further ionization if the ambipolar diffusion speed exceeds about 40 km s^{-1} (Draine et al. 1983). The ambipolar diffusion speed is approximately

$$V_D \approx 20\,\mathrm{km\,s}^{-1} \left(\frac{B}{50G}\right)^2 \left(\frac{L}{1 \times 10^{14}\mathrm{cm}}\right)^{-1} \left(\frac{X_i}{10^{-6}}\right)^{-1} \left(\frac{n_H}{1 \times 10^{10}\mathrm{cm}^{-3}}\right)^{-2} \quad (6)$$

where L is the lengthscale over which the field varies. Hence, the fractional ionization in SiO masers would be maintained by the ambipolar diffusion itself at a value of order 10^{-6} if $n_H \approx 10^{10}$ cm^{-3} and of order 10^{-8} if $n_H \approx 10^{11}$ cm^{-3}.

Since most SiO masers move away from the central stars at speeds less than 40 km s^{-1} and the field strengths can be maintained only if the ambipolar diffusion speed is at most comparable to the maser speed, ion-neutral collisions cannot be the dominant source of ionization in the SiO masers. Ultraviolet, infrared, and VLA observations show that late-type evolved oxygen-rich stars have chromospheres which emit in the ultraviolet and that the ultraviolet emission is weaker in later-type stars (Stencel et al. 1986; Judge 1989). The ultraviolet emission of α Ori is sufficient to maintain a considerable fraction of silicon as Si^+ in the stellar envelope (Clegg et al. 1983; Haas & Glassgold 1993) and the ultraviolet emission of the chromospheres of other late-type stars may also affect the ionization in their SiO maser regions. The fact that SiO maser emission is not seen from Mira variables with spectral types earlier than M6 (Patel et al. 1992) may be a consequence of the ultraviolet emission from earlier-type Miras preventing the chemistry that in late-type Miras results in SiO formation in regions where conditions for maser pumping obtain.

A lower bound on X_i in SiO maser regions can be set from the requirement that ambipolar diffusion does not heat the maser region gas too much. The heating timescale due to ambipolar diffusion is roughly

$$t_h \approx \frac{5k_B T}{2\alpha_{in} X_i \rho_H V_D^2} \approx 3 \times 10^3\mathrm{s} \left(\frac{T}{1500\mathrm{K}}\right) \left(\frac{X_i}{10^{-6}}\right) \left(\frac{n_H}{1 \times 10^{10}\mathrm{cm}^{-3}}\right)^3 \left(\frac{B}{50G}\right)^{-4} \left(\frac{L}{1 \times 10^{14}\mathrm{cm}}\right)^2 \quad (7)$$

where α_{in} is the ion-neutral rate coefficient, k_B is Boltzmann's constant and ρ_H is the mass density. Muchmore et al. (1987) give 5×10^4 s as an estimate of the cooling timescale at $n_H \approx 10^{10}$ cm^{-3} in the envelope of a star with an effective temperature of 3000K. If n_H in SiO masers is as low as 1×10^{10} cm^{-3}, the requirement that t_h exceed this cooling time would imply that $X_i \approx 10^{-5}$ for the normalization parameters adopted for Eqs. (6) and (7). Obviously X_i may be much less in maser regions with higher number densities.

The values of X_i associated with $n_H = 1 \times 10^{10}$ cm^{-3} imply that SiO masers at that number density might be pumped as a consequence of charged particle–neutral relative streaming and the differences between electron, ion and neutral temperatures which the streaming establishes (Kylafis & Norman 1987).

4. Consequences for outer envelope dynamics

In the current "standard" model (e.g. Bowen 1988; Cherchneff et al. 1992; MacGregor & Stencel 1992; Habing et al. 1994) of mass loss from highly evolved late-type stars (with sufficiently weak chromospheric ultraviolet emission that it does not prevent dust formation), radiation pressure on dust grains and dust–gas momentum transfer are responsible for the acceleration of gas to its final speed, whereas the shocks formed as a consequence of stellar pulsation lead to the extension of the atmosphere to the dust condensation radius. We are suggesting that buoyancy as well as the kinematic pressure plays a role in establishing the structure and dynamics of gas up to at least the dust condensation radius. Buoyancy by itself is not sufficient to drive mass loss because field lines remain anchored in denser gas further down in the atmosphere and magnetic tension eventually arrests the motion of less dense gas relative to that deeper in the stellar atmosphere. However, if the Parker instability and buoyancy influence the structure and dynamics of gas up to the dust condensation radius they affect the mass loss rate because they help determine the density at that radius.

In addition, the fact that field lines in dusty parts of the shell are anchored in gas at higher densities has at least two consequences for the transfer of momentum from grains to gas. One is that some grains which might otherwise decouple from the gas (MacGregor & Stencel 1992; Habing et al. 1994) remain well coupled to it by their interaction with the magnetic field (Hartquist & Havnes 1994). The other is that the momentum transfer from the grains to the gas is not a local process, but is rather distributed throughout the gas in the dusty regions and the denser gas interior to those regions but connected to them by the magnetic field; this implies that the effects of radiation pressure are important at depths greater than the dust condensation radius. These consequences for momentum transfer from dust to grains do not alter the conclusion of Hartquist et al. (1995) that the evolutionary tracks of some stars in the vicinities of some active galactic nuclei (AGNs) probably differ greatly from those of galactic stars because the radiation field of a sufficiently nearby AGN suppresses dust formation in stellar envelopes.

Dyson et al. (1989) have suggested that clumps existing in the planetary nebula NGC 7293 were formed in the same manner as the clumps associated with SiO masers. On the basis of survival arguments Dyson et al. (1989) concluded that the planetary nebula clumps must have masses of about $10^{-4}\mathrm{M}_\odot$. Inspection of Eq. (5) shows that the formation of such massive clumps by the Parker instability would have required

$$\left(\frac{M_*}{1.2\mathrm{M}_\odot}\right)^{-3}\left(\frac{R_s}{300\mathrm{R}_\odot}\right)^{6}\left(\frac{B_P}{50G}\right)^{2}\left(\frac{V_{A,P}}{10\mathrm{km\,s}^{-1}}\right)^{4}\approx 300 \quad (8)$$

a criterion which, given the high powers at which the quantities on the left hand side of (8) appear, is conceivably met in some evolved late-type stars.

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