Mon. Not. R. astr. Soc. (1984) **207**, 339–354

On the nature of the strong cosmic H_2O masers

V. S. Strelnitskij *Astronomical Council of the USSR Academy of Sciences, 48 Pjatnitskaja Str, Moscow 109 017, USSR*

Received 1983 July 25; in original form 1983 March 21

Summary. It is shown that the minimum pump power necessary to explain the observed energy flux in a maser line is proportional to l^{-3}, where l is the amplification length. The probable upper limits for l are estimated from the observed velocity gradients in several H_2O clusters associated with the regions of star formation: they vary from ~0.5 to ~50 AU. Using the 'equivalent pumping transition' approximation, it is argued that with such small l all the previous models fail to explain the observed maser energy fluxes. It is shown that the degree of ionization in the H_2O condensations could be much higher than commonly believed – up to $\sim 10^{-6}$–10^{-5}, and that the electrons could be there appreciably (~10 per cent) cooler than the H_2 molecules. Two possible astrophysical models, based on the collision–collisional pump working in such a two-temperature gas, are proposed. The models develop the concept of the H_2O masers as sites of interaction between the stellar wind of a young star and the gas condensations in its envelope, and they differ mainly by the mass of the star. It is suggested that the origin of the maser condensations can be related with the formation of giant protoplanets.

1 Introduction

This paper is devoted to the problem of the strong H_2O masers observed in the regions of active star formation. It presents the detailed version of a short communication published earlier (Strelnitskij 1980a).

While several of the models of H_2O masers hitherto proposed were satisfactory in explaining the subsources of middle and low power, all of them suffer difficulties in explanation of the extremely strong subsources (ESS), such as observed in W49, W51 and in some other objects. It is shown in Section 2.1 that the specific pump power (cm^{-3} s^{-1}), required for explanation of the observed energy flux in a maser line, is determined by the product of the squared distance to the source, D^2, and the cubed reciprocal amplification length, l^{-3}. On these grounds a short analysis of the previous models is given (Sections 2.2–2.4) and it is argued that the main reason to expose all of them to doubt is the recent VLBI mapping of the ESSs which strongly restricts possible values of l.

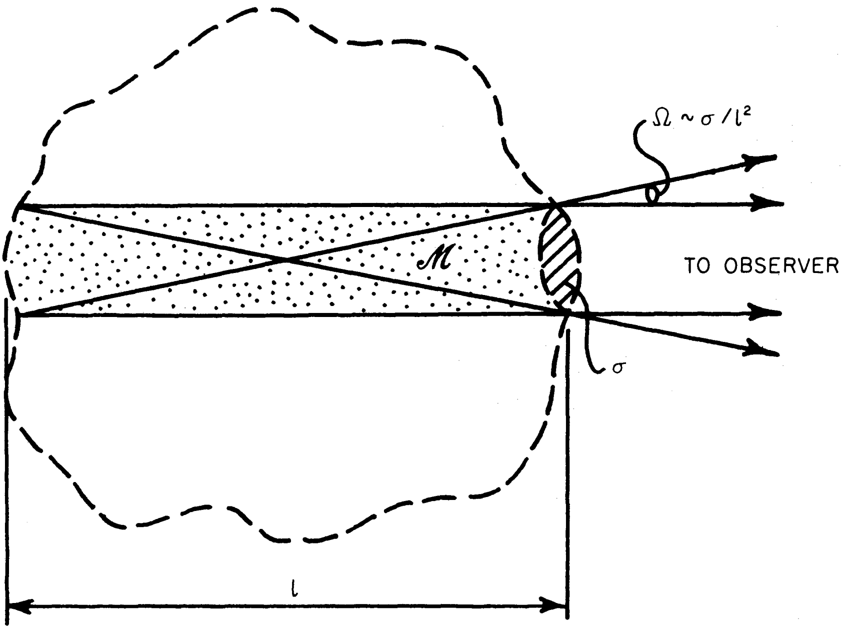

Figure 1. The derivation of equation (3). The apparent emitting surface σ is dashed, the region $\mathcal{M}$ contributing to the observed emission is dotted. The region $\mathcal{M}$ has a volume $V \sim \sigma l$; the emergent maser emission is confined by a solid angle $\Omega \sim \sigma/l^2$.

A new pump mechanism, based on the collisional processes only (Strelnitskij 1980a, b) is reconsidered in Section 3.1 and in Sections 3.2–3.5 the astrophysical models are proposed, which enable this mechanism to work and to provide the observed maser fluxes from the ESSs. The models give concrete expression to the general concept of the H_2O sources as sites of interaction between the stellar wind of a star and the gas condensations in its envelope (Strelnitskij & Syunyaev 1972; Pikelner & Strelnitskij 1976; Norman & Silk 1979; Genzel *et al.* 1978; Elmegreen & Morris 1979).

2 Previous models

2.1 REQUIREMENTS TO THE PUMP POWER

Let us first derive the minimum pump power necessary to explain the observed energy flux in a maser spectral feature.

It is easy to show that the least requirements occur with the elongation along the line-of-sight geometry of the region emitting the observed photons (region $\mathcal{M}$ in Fig. 1). Whatever the cause of the elongation – elongated geometry of the maser condensation itself, or non-uniformity of the saturation across the quasi-spherical maser – the volume of the region $\mathcal{M}$ will be of the order of σl, and the solid angle Ω of the maser beam at the maser's going out, of the order of σ/l^2, where σ is the area of the apparent emitting surface and l is the length of the region $\mathcal{M}$ along the line-of-sight. Thus, the upper limit for the emergent intensity at the line centre, given by the condition of total use of the pump power (full saturation) may be written as

$$I \lesssim \frac{n_1 \Delta P \sigma l h\nu_{12}}{2\Omega\sigma\Delta\nu_{12}} \approx \frac{n_1 \Delta P l^3 h\nu_{12}}{2\sigma\Delta\nu_{12}}, \tag{1}$$

Table 1. Characteristics of some H_2O clusters.

Source	D (kpc)	F (Jy)	dv/dx (km s^{-1})/AU	l_{max} (AU)	$(n_1\Delta P)_{min}$ (cm^{-3} s^{-1})	Features used (km s^{-1})	References
(1)	(2)	(3)	(4)	(5)	(6)	(7)	
W49	14	10^5	>0.02	<50	$\gtrsim 10^3$	−4.7, 8.2, 10.9	(1)†
W51, Main	8	10^3	0.3	3	$\sim 10^4$	62.2, 68.6, 71.1	(2)
W51 N	8	20	0.6	2	$\sim 10^3$	Cluster A	(3)
		5×10^2	0.1	10	$\sim 10^3$	Cluster C	(3)
OMC 1	0.5	2×10^6★	>0.05	<20	$\gtrsim 10^3$	8.3, 18.3, 24.8	(1)
L 1450 (H-H 7−11)	0.5	10^2	2	0.5	$\sim 10^3$	Source A (−7.3, −15.7, −17.3)	(4)
W3 OH	3	10^3	0.1	10	$\sim 10^2 - 10^3$	Clusters 3 and 4	(7)
ON 1	3	10^2	0.3	3	$\sim 10^2 - 10^3$	8.9, 14.3; 12.6, 14.3, 15.3	(5)
ON 2	5.5	2×10^2	0.2	5	$\sim 10^3$	−0.6, −3.2, −21.3	(5)
Cep A	0.7	10^3	0.3	3	$\sim 10^2$	Cluster BCH	(6)
		10^2	1.5	0.7	$\sim 10^3$	Pair IJ	(6)
43.8−0.1	3	30	>0.2	<5	$\gtrsim 40$	45.2, 46.7, 49.1, 51.2	(5)

References
(1) Genzel *et al.* (1978), (2) Genzel *et al.* (1979), (3) Schneps *et al.* (1981), 4) Haschick *et al.* (1980), (5) Downes *et al.* (1979), (6) Lada *et al.* (1981), (7) Giufrida *et al.* (1981).
Notes
★ Maximum F during the flare of the 8.3 km s^{-1} feature in 1979 (Matveenko 1981).
† One deduces the same $(n_1\Delta P)_{min}$ from the analysis of many H_2O clusters in the more detailed VLBI map of W49 by Walker, Matsakis & Garcia-Barreto (1982).

where n_1 is the number density of molecules at the lower signal level, $\Delta P = P_{12} - P_{21}$ is the net rate of pumping from the lower to the upper signal level, ν_{12} is the frequency of the signal transition, $\Delta\nu_{12}$ the observed linewidth (FWHM) and h the Planck constant.

On the other hand, I is related to the observed flux density F by

$$F \approx I\omega \approx I(\sigma/D^2), \tag{2}$$

where ω is the solid angle subtended by the surface σ, as viewed at the source distance D from the Earth.

From the equations (1) and (2) one derives the sought for lower limit for the specific pump power $n_1\Delta P$:

$$\left(\frac{n_1}{\text{cm}^{-3}}\right)\left(\frac{\Delta P}{\text{s}^{-1}}\right) \gtrsim 10 \left(\frac{l}{\text{AU}}\right)^{-3} \left(\frac{D}{\text{kpc}}\right)^2 \left(\frac{F}{\text{Jy}}\right) \left(\frac{\Delta\nu_{12}/\nu_{12}}{10^{-6}}\right). \tag{3}$$

The independence of the saturated maser luminosity of σ and its cubic (but not linear) dependence on l, which at the first glance may appear paradoxical, are explained by the directivity of the stimulated emission. The l^3-dependence of the brightness temperature of a saturated cylindrical maser was first pointed out by Goldreich & Keeley (1972). Passing from elongated to the spherical geometry of the region $\mathscr{M}$ raises the numerical coefficient at the right-hand side of equation (3) only to $\sim 10^2$. So, practically this equation is valid independently of the source geometry.

The upper limit for l may be obtained from the analysis of the structure obtained by interferometry of the region of maser activity. Recently it was established by VLBI that maser condensations with different radial velocities are often concentrated in compact 'centres of activity' (e.g. Genzel *et al.* 1978), or even in 'multiple velocity knots'; very compact (not resolved by the VLBI) regions containing several velocity components (e.g. Genzel *et al.* 1979; Downes *et al.* 1979). Generally the distribution of the velocity knots within the cluster appears quite random. This, and also the facts of correlation of the widths of velocity features with the flux densities in their centres (Lada *et al.* 1981) and with the velocity dispersion within the cluster (Schneps *et al*, 1981) make it very probable that the velocity dispersion is chiefly of a turbulent nature. Then with the natural assumption that the line-of-sight and the perpendicular velocity gradients, dv/dx [(km s^{-1})/AU] are of the same order and, taking into account that for the formation of the observed line-widths, $\Delta v \lesssim 1$ km s^{-1}, the corresponding velocity coherence along the amplification path is needed, one obtains the probable upper limit for l(AU) as $l_{\max} \sim 1/(dv/dx)$.

Table 1 presents the results for several bright H_2O sources. The probable values of dv/dx, estimated from the observed projected distances (or their upper limits) between the condensations with the given radial velocities, are indicated in column 4, and the corresponding $l_{\max}$ are listed in column 5. For all the features we assumed $\Delta\nu_{12}/\nu_{12} = 2\times10^{-6}$ (which is always correct within a factor 2–3) and calculated by means of equation (3) the minimum required pump power, given in column 6. It is obvious that for many sources the required pump power exceeds $\sim 10^2$–10^3 cm^{-3} s^{-1} and may even attain $\sim 10^4$ cm^{-3} s^{-1}.

We show in the next subsections that the previous models fail to provide such high pump powers. We briefly analyse the four most elaborated models, to be short, brief three-letter designations (Strelnitskij 1980a, b), the first letter giving the type of the energy source (C, collisional; R, radiative), the second the type of the energy sink and the third the type of the transitions constituting the pump cycles (r, rotational; v, vibrational).

2.2 LITVAK MODEL (RRv)

Litvak (1969) has suggested that the pumping occurs behind a shock front propagating in a collapsing protostar. The vibration–rotational photons radiated by the H_2O molecules near the shock excite the cooler molecules further downstream; subsequent radiative de-excitations tend, according to Litvak, to overpopulate the upper signal level 6_{16} relative to the lower level 5_{23}.

The pump power $n_1\Delta P$ would be of the order of net rate of population transfer in an 'equivalent' three-step pump cycle (Fig. 2), which, due to the smallness of the signal frequency ν_{12} as compared with the pump frequencies ν_{13} and ν_{23}, may be approximated by the net rate of population transfer in an 'equivalent pumping transition' 1–3 (or 2–3) (Strelnitskij 1980b; 1983).

So, in the case of the RRv pump one has: $n_1\Delta P_{\rm RRv} \sim n_3A_{32}Z_{32}$, where $A_{32}\equiv A_{\rm v}$ is the Einstein coefficient for spontaneous emission and Z_{32} is the net radiative bracket (Mihalas 1978). Throughout the greatest part of a layer of optical thickness $\tau_0 \gg 1$ one has: $Z(\tau) \sim L(\tau_0/2)$ (Borodina 1975; Schmeld, Strelnitskij & Muzylev 1976), where $L(t)$ [$\simeq (2t\sqrt{\pi \ln t})^{-1}$ with the Doppler line broadening] is the transmission function. Hence, one can admit for the whole maser condensation:

$$n_1\Delta P_{\rm RRv} \sim \frac{n_1 \exp(-E_{\rm v}/kT_{\rm v})\, A_{\rm v}}{\tau_0^{\rm v}\sqrt{\pi \ln(\tau_0^{\rm v}/2)}} \lesssim \frac{8\pi c(\Delta\nu_{\rm v}/\nu_{\rm v})}{\lambda_{\rm v}^3 d[\exp(E_{\rm v}/kT_{\rm v})-1]\sqrt{\pi\ln(\tau_0^{\rm v}/2)}}, \quad (4)$$

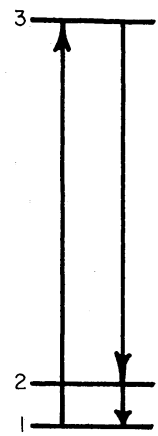

Figure 2. The three-step pump cycle, equivalent to the real pumping process. The net population transfer is indicated by the arrows. 1–3, the excitation step; 3–2, the sink step; 1–2, the signal step closing the cycle. If $\nu_{12} \ll \nu_{13}$ and ν_{23}, the pump rate can be approximated by the net rate of population transfer in a two-level system with the frequency $\nu_{13}(\simeq\nu_{23})$ and with the same excitation and sink processes.

where λ_v is the wavelength, ν_v the frequency and E_v the energy of the typical pump photon, $\Delta\nu_v$ is the width of the pump line, T_v the vibrational temperature, d the smallest size of the maser condensation, τ_0^v its optical thickness at the centre of the pump line and c the velocity of light. It may be shown (Bolgova, Strelnitskij & Umanskij 1982) that RRv pumping of the ESSs requires $\tau_0^v \gtrsim 10^2$. Taking for d the observed sizes of the ESSs ($\sim$1 AU) and $T_v \lesssim 2000$ K (see Section 3.2), one obtains from (4) for ν_2 and ν_3 bands: $n_1 \Delta P_{\mathrm{RRv}} \lesssim 10\,[(\Delta\nu_v/\nu_v)/10^{-6}]$. Taking also into account that maser lines could at best be 5–6 times narrower than the pump lines (e.g. Strelnitskij 1974) one has $n_1 \Delta P_{\mathrm{RRv}} \lesssim 10^2$ cm^{-3} s^{-1}. Hence, RRv pump power is insufficient for most ESSs listed in Table 1.

2.3 GOLDREICH–KWAN (RCv) AND NORMAN–SILK (CRv) MODELS

In the model by Goldreich & Kwan (1974) the infrared radiation of the circumstellar dust heated by the radiation of a variable star excites the circumstellar H_2O molecules to the (010) vibrational state, while the energy sink of the pump is provided by de-excitation of this state in collisions with the cooler (than dust) H_2 molecules. The requirement that the thermal relaxation between dust and gas be sufficiently slow puts the upper limit to H_2 number density: $n_H \lesssim 10^9$ cm^{-3}.

The pump rate can be estimated by the 'equivalent pumping transition' approximation (see an intermediate step in equation 7):

$$\Delta P_{\mathrm{RCv}} \sim n_H q_H^v [\exp(-E_v/kT_v) - \exp(-E_v/kT_H)]$$
$$\lesssim n_H q_H^v [\exp(-E_v/kT_d) - \exp(-E_v/kT_H)], \tag{5}$$

where T_d is the dust temperature, T_H the H_2 (kinetic) temperature, q_H^v the rate of collisional de-excitation of an H_2O vibration–rotational transition by the H_2 molecules and we have taken into account that $T_d > T_v > T_H$.

Goldreich & Kwan (1974) assumed the gas-kinetic value $q_H^v \sim 10^{-9}$ cm^3 s^{-1}, which seems to be a great overestimation (Strelnitskij 1980a; Deguchi 1981). Available experimental data on the de-excitation rates of H_2O vibrational levels by N_2, He and Ar and comparison with the data on the de-excitation of NH_2 vibrational levels by H_2 and Ar lead to the conclusion that for H_2O $q_H^v \lesssim 10^{-13}\,cm^3\,s^{-1}$ at $T_H \sim 10^3$ K (Bolgova *et al.* 1982). With this upper limit and with $T_d = 1000$ K, $T_H = 800$ K assumed by Goldreich & Kwan, one obtains from (5): $\Delta P_{RCv} \lesssim 10^{-14} n_H$. Assuming $n_H = 10^9\,cm^{-3}$, the cosmic oxygen abundance (hence, $n_{H_2O} \lesssim 10^{-3}\,n_H$) and, taking into account that $\eta_r \equiv n_1/n_{H_2O} \sim 10^{-2}$ at rotational temperatures $T_r \sim 10^2$–10^3 K, one has $n_1 \Delta P_{CRv} \lesssim 10^{-1}\,cm^{-3}\,s^{-1}$. Hence, this model is very inadequate for pumping the ESSs.

Moreover, with such a low pump rate the possibility of inversion becomes doubtful. It may be shown that the inversion is possible only if

$$\Delta P > C_{rel}(E_m/kT_H), \tag{6}$$

where E_m is the energy of a maser quantum; $C_{rel}(\gtrsim n_H q_H^r)$ is the total rate of relaxation collisional transitions from a signal level into other rotational levels; $q_H^r \sim 10^{-11}\,cm^3\,s^{-1}$ (Bolgova *et al.* 1982) is the average rate of collisional de-excitation of a typical rotational transition by H_2 molecules. Putting into (6) the above $\Delta P_{CRv} \lesssim 10^{-14}\,n_H$, $E_m/k = 1.07$ K, $T_H = 800$ K, and considering the fact that C_{rel} is at least several times greater than $n_H q_H^r$, one sees that condition (6) is scarcely fulfilled and hence the inversion is hardly possible.

In the model by Norman & Silk (1979) the pumping cycle consists of a collisional excitation of H_2O to the (010) state by hot H_2 molecules with subsequent spontaneous de-excitation into the (000) state, the signal transition closing the cycle. Energy is supplied by the stellar wind blowing around the maser condensation and heating the H_2 molecules therein to ~1500 K. An obvious upper limit for the CRv pump rate is given by the rate of collisional excitation of the (010) state: $\Delta P_{CRv} < n_H q_H^v \cdot \exp(-E_v/kT_H)$. With $n_H \sim 10^9\,cm^{-3}$ assumed by Norman & Silk (1979) ΔP_{CRv} is practically as small as ΔP_{CRv} given by equation (5), hence, the Norman–Silk model is equally inadequate for pumping of the ESSs. Note, however, that if n_H could be raised up to $\sim 10^{12}\,cm^{-3}$, the CRv process would provide the necessary pump rate (see Section 4), but the possibility of inversion (condition 6) remains questionable for this process at *any* n_H.

2.4 CRr MODELS

In these models [Hills (see Turner 1970); Strelnitskij 1971, 1973; de Jong 1973; Shmeld *et al.* 1976] pump cycles are constructed on rotational transitions and include collisional (or chemical) excitations and subsequent spontaneous de-excitations. Difficulties with the energy sink indicated by Goldreich & Kwan (1974) are avoidable in the models by de Jong, or by Shmeld *et al.* (see Strelnitskij 1981, and also a recent model by Deguchi 1981). But in all these models l was assumed to be $\sim 10^3$ AU, which by far overpasses the upper limits for l, deduced from the recent interferometry (Table 1). Our calculations give as a firm upper limit $n_1 \Delta P_{CRr} < 10^3\,cm^{-3}\,s^{-1}$, but the real $n_1 \Delta P_{CRr}$ seems to be at least one order of magnitude smaller. The CRr process, like the preceding pump processes, seems to be insufficient to pump ESSs either.

3 New models

3.1 CCr PUMP

To avoid these difficulties a new pump mechanism based on collisonal processes between the rotational levels (CCr) was proposed (Strelnitskij 1980a, b; Bolgova 1981). CCr pumping is possible if there are two kinds of particles, e.g. H_2 molecules and electrons, with different kinetic temperatures T_H and T_e. Collisions with the hot particles provide the excitation of molecules in the pump cycles, whereas collisions with the cold particles provide the energy sink. Solutions to the balance equations for the rotational level populations (Bolgova 1981) show, that superheating and inversion of the 6_{16}–5_{23} transition are possible only with $T_H > T_e$, and only such temperature relation will be considered below.

In this case the pump rate, approximated by the net rate of population transfer in an 'equivalent' transition 1–3 (Fig. 2), is given by

$$\Delta P_{CCr} \sim [n_1 n_H q_H^r(13) - n_3 n_H q_H^r(31)]/n_1 = n_H q_H^r[\exp(-E_{13}/kT_H) - \exp(-E_{13}/kT_{13})], \tag{7}$$

where T_{13} is the excitation temperature of the 1–3 transition. Under temperatures $\gtrsim 10^3$ K one can consider that $E_{13}(\equiv E_r)$ is much smaller than kT_H, kT_e, kT_{13} and then equation (7) is simplified:

$$\Delta P_{CCr} \sim n_H q_H^r (E_r/k)(T_{13}^{-1} - T_H^{-1}). \tag{7a}$$

Solution of the statistical equilibrium equations for a two-level system (e.g. Spitzer 1978, p. 83) gives:

$$T_{13}^{-1} = (n_H q_H^r T_H^{-1} + n_e q_e^r T_e^{-1})/(n_H q_H^r + n_e q_e^r),$$

where n_e is the electron number density and q_e^r is the rate coefficient for the de-excitation of a typical rotational transition by electron impact. Substituting this into equation (7a) leads to:

$$\Delta P_{CCr} \sim \frac{E_r}{k}(T_e^{-1} - T_H^{-1})\left(\frac{1}{n_e q_e^r} + \frac{1}{n_H q_H^r}\right)^{-1}. \tag{8}$$

It is easy to show that ΔP_{CCr} reaches its maximum at $n_e q_e^r = n_H q_H^r$. Assuming $q_e^r \sim 10^{-6}$ $cm^3\,s^{-1}$ (Itikava 1972) and $q_H^r \sim 10^{-11}$ $cm^3\,s^{-1}$ (see above), one finds that the maximum of ΔP_{CCr} is attained at $n_e/n_H \sim 10^{-5}$. At smaller n_e/n_H (when $n_e q_e^r \ll n_H q_H^r$) the equation (8) gives:

$$\Delta P_{CCr} \sim n_e q_e^r \frac{E_r}{kT_e}\left(\frac{T_H - T_e}{T_H}\right). \tag{8a}$$

There is practically no upper limit for density at which the CCr pump can work at provided the necessary n_e/n_H and $T_H > T_e$ (one can realize the CCr pump at high densities, typical for laboratory experiment).

Models of the strong H_2O sources, in which these conditions would be fulfilled, are proposed below, in Sections 3.2 and 3.3. It is now commonly assumed that the strong H_2O masers are associated only with massive, luminous young stars (e.g. Genzel & Downes 1977). An argument against this point of view was found in infrared observations: there are no sufficiently strong infrared sources towards many strong H_2O sources (Forster, Welch & Wright 1977; Forster *et al.* 1978). Although this could in principle be explained by the absorption of infrared radiation in the surrounding gas–dust cloud (Genzel *et al.* 1978) or in the circumstellar disc (Elmegreen & Morris 1979) it is, however, interesting to consider the possibility of the ESSs pumping in the surroundings of a moderately luminous star as well (Norman & Silk 1979; Strelnitskij 1980a) because for such a star the absence of strong

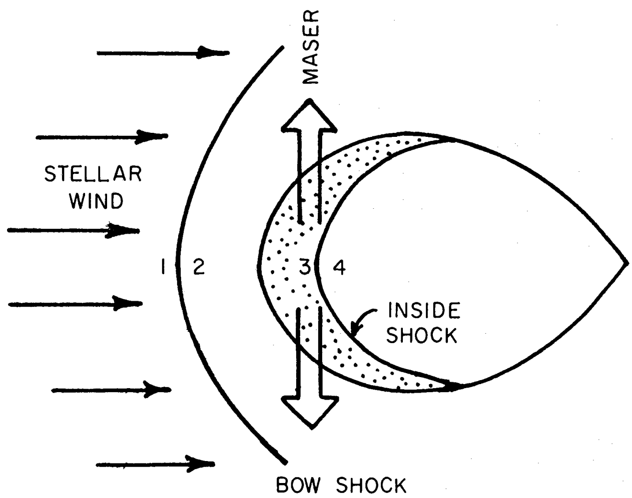

Figure 3. Simplified structure of an ESS. 1, the free stellar wind; 2, the compressed post-shock stellar wind; 4, the unperturbed part of the gas condensation; 3, the part of the condensation compressed by the shock 3–4 (the pump region).

infrared source is explained without special hypotheses. Both possibilities are considered below.

3.2 MODEL 1: YOUNG STAR OF GREAT MASS

If the mass of a young star is great, say, $M_* \gtrsim 5–10\,M_\odot$, its luminosity may be as high as $\sim 10^5\,L_\odot$. Infrared observations of the Orion Molecular Cloud OMC-1, where one of the brightest H_2O sources is observed, show that at least one such star is present here (presumably IRc2) and that it gives off an extremely strong stellar wind with a mass-loss rate $\dot{M} \sim 10^{-3}\,M_\odot\,\mathrm{yr}^{-1}$ and velocities $v_\mathrm{w} \sim 20–200\,\mathrm{km\,s}^{-1}$ (Downes *et al.* 1981). At later stages of evolution massive OB-stars lose mass usually at a rate $\dot{M} \sim 10^{-5}\,M_\odot\,\mathrm{yr}^{-1}$ with the typical velocities $v_\mathrm{w} \sim 1500\,\mathrm{km\,s}^{-1}$ (de Jager 1980). Next we show that both these sets of stellar wind parameters (we shall call them model 1a and model 1b respectively) are convenient to provide the CCr pumping of the ESSs.

Following Genzel *et al.* (1978) and Elmegreen & Morris (1979) we assume that each centre of H_2O maser activity is associated with a circumstellar gas–dust disc around a young star. Suppose that at the moment the strong stellar wind is 'switched on' there is a density condensation within the disc at $r \sim 5 \times 10^{14}–10^{15}$ cm from the star, with the typical size $\sim 3 \times 10^{14}$ cm. The origin of the condensation (see, for example, Norman & Silk 1979; Strelnitskij 1980a, and below) is not so important: even the weak initial density inhomogeneities would be strengthened by the stellar wind. Important is the involved gas dynamics, which may be highly complicated. The wind hindered by the condensation produces a bow shock 1–2 (Fig. 3). The pressure of the compressed wind will produce a shock wave 3–4 within the condensation itself. Due to inhomogeneous density distribution within the condensation this shock may have a complicated form. Some parts of it may intersect with each other and thereby strengthen. There, where the shock reaches the surface of the condensation, the gas will expand and an expansion wave will move in the opposite direction. The picture as a whole may be complicated and non-stationary, due, in particular, to the absence of stationarity of the stellar wind itself. However, the net effects would be acceleration of the condensation, as a whole, away from the star (Strelnitskij & Syunyaev 1972) and the energy dissipation within the condensation. We consider the latter process as the pump source for the H_2O masers and we will make now some crude quantitative estimates to show the principal possibility of the sufficiently powerful CCr pumping in such conditions.

The height of the circumstellar disc at the distance r from the star is (Ward 1976):

$$h(r) \sim 2c_s(r^3/GM_*)^{1/2}, \tag{9}$$

where $c_s(\sim 10^5 \text{ cm s}^{-1})$ is the velocity of sound, G the gravitational constant. At the assumed distance from the star is found from (10) to be $\sim 10^{10}$ cm^{-3}. To obtain order-of-magnitude that the surface density of the disc, δ, is proportional to r^{-1} (Elmegreen 1978), one finds the average n_H at the distance r from the star:

$$n_H(r) \sim (\delta/\text{hm}) \sim (G^{1/2}M_D M_*^{1/2})/(8\pi c_s r_D r^{5/2}), \tag{10}$$

where m is the mass of H_2 molecule, M_D is the mass of the disc and r_D its radius. With the possible values $M_D \sim 1\ M_\odot$, $r_D \sim 10^{16}$–10^{17} cm the typical value of n_H at the assumed distance from the star is found from (10) to be $\sim 10^{10}$ cm^{-3}. To obtain order-of-magnitude estimates of gas-dynamical parameters we approximate the condensation in its initial state by a homogeneous cylinder of this density, with the axis orientated along the radius-vector from the star.

Gas temperature behind the front 1–2 is determined by (Zel'dovich & Raizer 1966):

$$T \simeq \left(\frac{\gamma-1}{\gamma+1}\right)^2 \left(\frac{3\mu v_w^2}{Rg}\right), \tag{11}$$

where $\gamma = 5/3$ is the ratio of specific heats of the mono-atomic gas, $\mu \simeq 1/2$ is the mean molecular weight of the gas, R_g the gas constant. From equation (11) one finds for models 1a and 1b respectively: $T_2 \sim 5 \times 10^5$ K and 10^7 K.

The velocity of the shock 3–4 relative to the condensation can be evaluated from the condition of conservation of momentum:

$$v_{34} \simeq v_w(\rho_1/\rho_4)^{1/2} \simeq (\dot{M}v_w/4\pi r^2\rho_4)^{1/2},$$

where ρ is the mass density. For both models v_{34} turns out to be ~ 10 km s^{-1}. The gas temperature T_3 just behind the front 3–4 is evaluated from equation (11); with $v_{34} \sim 10$ km s^{-1}, $\mu \simeq 2$ and $\gamma = 7/5$ (H_2 with the excited rotations) T_3 is of the order of several thousand K for both models. Further downstream T is increasing to ~ 1500 K due to the excitation of the H_2 vibrations, and then it continues to drop more slowly due to the vibrational excitation of the polar molecules (CO, H_2O) with subsequent spontaneous de-excitation. The velocity of the gas relative to the front 3–4 just behind the front is $v_3 \simeq v_{34}(\gamma-1)/(\gamma+1) \sim 1$ km s^{-1}, and the gas number density is $n_3 \simeq n_4(\gamma+1)/(\gamma-1) \sim 10^{11}$ cm^{-3}. Further downstream while T is decreasing, n_H is increasing and v decreasing, but within the interval of interest for the CCr pumping ($T_H \sim 1000$–2000 K) v and n_H change at maximum by several times.

Let us estimate the width of this region, roughly the region which is cooled by radiative losses on the polar molecules. Since H_2O and CO are almost equally efficient in cooling, one can assume, for simplicity, that all the oxygen is tied in H_2O, i.e. $H_2O/H_2 \sim 10^{-3}$. Then the width of the region will be of the order $d \sim (3/2)n_H k T_3 v_3/\mathscr{L}$, with the cooling rate $\mathscr{L} \sim n_H n_{H_2O} q_H^v \mathscr{X}_H^v \exp(-E_v/kT_H)$, where $\mathscr{X}_H^v \sim kT_H/E_r \sim 5$ is the typical number of transitions induced by collisions with H_2 transitions from a given rotational level of the (000) state to the rotational levels of the (010) state. Substituting $n_{\mathrm{H}}=10^{11}\,\mathrm{cm}^{-3}$, $n_{\mathrm{H_2O}}=10^8\,\mathrm{cm}^{-3}$, $q_{\mathrm{H}}^{\mathrm{v}} \sim 10^{-14}\,\mathrm{cm}^3\,\mathrm{s}^{-1}$, $E_{\mathrm{v}}/k \sim 2000$ K, $T_3 \sim 1500$ K, $v_3 \sim 1\,\mathrm{km\,s}^{-1}$, one finds: $d \sim 10^{11}$ cm.

In this estimate we have completely ignored the return of the energy to the gas by quenching collisions. This is justified when the H_2O vibrational temperature T_{v} is much lower than T_{H}, i.e. when the source function S_{v} of the vibration–rotational transitions is much smaller than the Planck function $B(T_{\mathrm{H}})$. This may be due to two processes: (1) escape of the photons from the medium and (2) their absorption by the cold dust. If the first process dominates then, for the central part of a plane-parallel slab, $S_{\mathrm{v}} \simeq \epsilon B(T_{\mathrm{H}})/\{L(\tau_0^{\mathrm{v}}/2) + \epsilon[1 - L(\tau_0^{\mathrm{v}}/2)]\}$ (e.g. Mihalas 1978), where $\epsilon \simeq n_{\mathrm{H}} q_{\mathrm{H}}^{\mathrm{v}}/(n_{\mathrm{H}} q_{\mathrm{H}}^{\mathrm{v}} + A_{\mathrm{v}})$. If the second process dominates then $S_{\mathrm{v}} \simeq \epsilon B(T_{\mathrm{H}})/[\beta\delta + \epsilon(1-\beta\delta)]$, where $\beta = \sigma_{\mathrm{v}}^{\mathrm{c}}/\sigma_{\mathrm{v}}^{\mathrm{m}}$ is the ratio of the absorption coefficients for the vibration–rotational photons due to the dust and to the molecules respectively, and $\delta(\simeq\sqrt{\ln(1/\beta)}$ for $\beta \ll 1$ and Doppler line broadening) accounts for the increase of the probability for a photon to die on a dust grain due to the brokenness of the photon's path (Ivanov 1969). With the above values of n_{H} and $q_{\mathrm{H}}^{\mathrm{v}}$ and assuming $A_{\mathrm{v}} \sim 10\,\mathrm{s}^{-1}$, $d \sim 10^{12}$ cm (see below), $\tau_0^{\mathrm{v}} \sim n_{\mathrm{H_2O}} \eta_{\mathrm{r}} k_{\mathrm{v}} d \sim 10^2$ ($k_{\mathrm{v}} \sim 10^{-16}\,\mathrm{cm}^2$ is a typical value of the absorption coefficient in a ν_2 line), one finds: $\epsilon \sim 10^{-4}$, $L(\tau_0^{\mathrm{v}}/2) \sim 3 \times 10^{-3}$, and hence $S_{\mathrm{v}} \sim 3 \times 10^{-2} B(T_{\mathrm{H}})$.

If the photons are removed from the medium primarily by the cold dust, S_{v} can be even smaller. The maximum absorption coefficient of the dirty dielectric grains may attain $\sigma^{\mathrm{c}} \sim 3 \times 10^{-23} n_{\mathrm{H}} (M_{\mathrm{d}}/M_{\mathrm{H}}) \lambda^{-1}\,(\mathrm{cm}^{-1})$ (Strelnitskij 1981), which for $M_{\mathrm{d}}/M_{\mathrm{H}} \sim 10^{-2}$, $n_{\mathrm{H}} \sim 10^{11}\,\mathrm{cm}^{-3}$, and $\lambda \simeq 6 \times 10^{-4}$ cm (ν_2 band) gives $\sigma_{\mathrm{v}}^{\mathrm{c}} \sim 5 \times 10^{-11}\,\mathrm{cm}^{-1}$. Then β (and $\beta\delta$) would be of order several tenths and hence $S_{\mathrm{v}} \sim (10^{-4}$–$10^{-3})\,B(T_{\mathrm{H}})$. However, neither the quantity nor the properties of the dust in the maser condensations are known. The real value of $\sigma_{\mathrm{v}}^{\mathrm{c}}$ can be significantly (1–2 orders) smaller, than the assumed maximum value and the reasonable conclusion would be that the dust can in principle be as efficient as the escape in removing the vibrational photons to provide $S_{\mathrm{v}} \ll B(T_{\mathrm{H}})$, and so $T_{\mathrm{v}} \ll T_{\mathrm{H}}$.

The residual thermalizing trapping of the vibrational photons (hence, somewhat greater $\mathscr{L}$), as well as somewhat greater values of $\dot{M} v_{\mathrm{w}}$ (hence, greater v_3) than assumed above, could maybe increase d up to one order but it would still remain smaller than the observed transverse dimensions of the maser spots ($\sim 10^{13}$ cm). This discrepancy could be avoided if one supposes that the observed size is an effective one, determined by a more or less complicated shock wave system within the condensation. The existence of such structures, with dimensions $<10^{13}$ cm, is indeed suggested by the VLBI (Matveenko 1981).

It is easy to see that with the typical values for the rotational transitions of H_2O, $k_{\mathrm{r}} \sim 10^{-14}\,\mathrm{cm}^2$, $q_{\mathrm{H}}^{\mathrm{r}} \sim 10^{-11}\,\mathrm{cm}^3\,\mathrm{s}^{-1}$, $A_{\mathrm{r}} \sim 1\,\mathrm{s}^{-1}$, S_{r} must be very close to $B(T_{\mathrm{H}})$ and hence contrary to T_{v}, $T_{\mathrm{r}} \simeq T_{\mathrm{H}}$ (of course, this is true only for a typical rotational transition and does not exclude significant differences between T_{r} and T_{H} for some special transitions, like the 6_{16}–5_{23}).

Next we shall demonstrate that in the radiatively cooled region T_{e} can be appreciably smaller than T_{H}. The electrons are cooled primarily by the collisional excitation of the H_2O vibrations (because, as stated above, $T_{\mathrm{v}} \ll T_{\mathrm{H}}$), with the rate per electron

$$\Lambda_{\mathrm{e}} \sim n_{\mathrm{e}} n_{\mathrm{H_2O}} q_{\mathrm{e}}^{\mathrm{v}} \mathscr{X}_{\mathrm{e}}^{\mathrm{v}} E_{\mathrm{v}} \exp(-E_{\mathrm{v}}/kT_{\mathrm{e}}) \qquad (\mathrm{erg\,s}^{-1}) \tag{12}$$

and they are heated collisionally by the thermalized to T_{H} degrees of freedom, primarily by: (1) H_2 translations, (2) H_2 vibrations, (3) H_2 rotations and (4) H_2O rotations (Fig. 4). For the sake of simplicity we shall compare the rates of these heating processes only at one typical pair of temperatures: $T_{\mathrm{H}} = 1500$ K, $T_{\mathrm{e}} = 1200$ K, which will be shown below.

(1) H_2 translations:

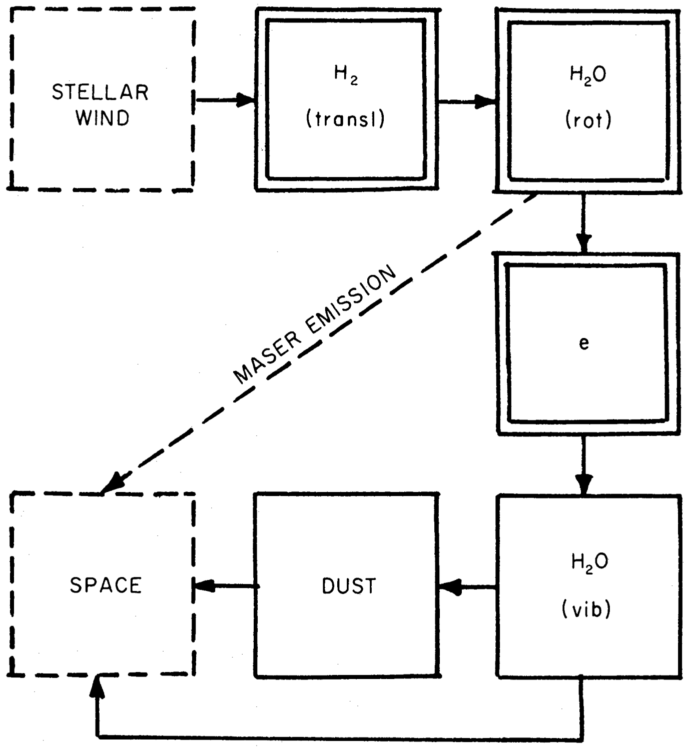

Figure 4. The main energy drain causing the CCr pumping in the radiatively cooled region behind the shock 3–4. The double squares mark the energy reservoires which directly participate in the CCr pump cycles.

$$\Gamma_e(H_2 - t) \sim n_H q_e^{H(t)} \frac{m_e}{m_H} k(T_H - T_e) \sim 4 \times 10^{-14} \qquad (\mathrm{erg\,s^{-1}}),$$

where $q_e^{H(t)} \sim 3 \times 10^{-8}\ \mathrm{cm^3\,s^{-1}}$ (Drukarev 1978) is the rate of elastic collisions between electrons and H_2 molecules for the momentum exchange and the mass ratio m_e/m_H provides the transition to the rate of energy exchange.

(2) H_2 vibrations:

$$\Gamma_e(H_2 - v) \sim n_H q_e^{H(v)} E_v(H_2)\{\exp[-E_v(H_2)/kT_H] - \exp[-E_v(H_2)/kT_e]\}$$
$$\sim 3 \times 10^{-14} \qquad (\mathrm{erg\,s^{-1}}),$$

where $q_e^{H(v)} \sim 3 \times 10^{-11}\ \mathrm{cm^3\,s^{-1}}$ is the rate of de-excitation of the $v = 1-0$ transition by electron impact (Drukarev 1978).

(3) H_2 rotations:

$$\Gamma_e(H_2 - r) \sim n_H q_e^{H(r)} E_r(H_2)\{\exp[-E_r(H_2)/kT_H] - \exp[-E_r(H_2)/kT_e]\} \sim 1 \times 10^{-13} \quad (\mathrm{erg\,s^{-1}}),$$

where $q_e^{H(r)} \sim 3 \times 10^{-10}\ \mathrm{cm^3\,s^{-1}}$ (Drukarev 1978).

(4) H_2O rotations:

This rate is obtained by means of the equation (8a):

$$\Gamma_e(H_2O - r) \sim n_{H_2O} q_e^r \mathscr{K}_e^r \frac{E_r^2}{kT_e} \left(\frac{T_H - T_e}{T_H}\right) \sim 3 \times 10^{-13} \qquad (\mathrm{erg\,s^{-1}}), \tag{13}$$

where $\mathscr{K}_e^r$ is the typical number of permitted transitions induced by electrons from a rotational level to other rotational levels ($\mathscr{K}_e^r \sim 3$) and $E_r \sim 3 \times 10^{-14}$ erg.

Now we can find the relative temperature difference $\Delta T/T \equiv (T_H - T_e)/T_H$ by equating the dominant heating rate of the electrons, $\Gamma_e(H_2O - r)$, with the cooling rate (12):

$$\frac{\Delta T}{T} \sim \frac{q_e^v \mathscr{X}_e^v}{q_r^r \mathscr{X}_e^r} \frac{E_v}{E_r} \frac{kT_e}{E_r} \exp\left(-\frac{E_v}{kT_e}\right). \tag{14}$$

Taking as the typical values: $T_e \sim 1500$ K; $E_v \sim 3 \times 10^{-13}$ erg; $q_e^v \mathscr{X}_e^v \sim 3 \times 10^{-8}$ cm^3 s^{-1} [deduced by Bolgova *et al.* 1982 from the experimental results by Seng & Linder 1976]; $E_v/E_r \sim 10$–15, $kT_e/E_r \sim 10$, one obtains: $\Delta T/T \sim 0.2$.

How reliable is this estimate? E_v/E_r and kT_e/E_r are known to a factor ~ 1; $q_e^v \mathscr{X}_e^v$ and $q_e^r \mathscr{X}_e^r$ would be erroneous by no more than several times. Hence, one can hope that the found $\Delta T/T$ is reliable within an order of magnitude.

Substituting this estimate into equation (8a) one finds $\Delta P_{CCr} \sim 2 \times 10^{-8}\, n_e$. Then it is readily shown that for establishing the inversion (condition 6) $n_e/n_H \gtrsim 2 \times 10^{-6}$ is sufficient. To fix the parameters let us assume, for example, $n_e/n_H = 3 \times 10^{-6}$. Then, with $n_H \sim 10^{11}$ cm^{-3}, $n_{H_2O} \sim 10^{-3}\, n_H$, $\eta_r \sim 10^{-2}$, one has $n_1 \Delta P_{CCr} \sim 6 \times 10^3$ cm^{-3} s^{-1}, which is roughly sufficient for all the ESSs listed in Table 1.

Next we consider some mechanisms which would provide such n_e/n_H.

One possible mechanism is connected with the well-known high efficiency of particle acceleration by the plasma turbulence behind shocks (Kaplan & Tsytovich 1972), in our case behind the shock 1–2. A high energy particle, with energy ϵ(MeV) loses an essential part of ϵ for ionization, having passed a column density $N_H \sim 10^{23}\, \epsilon^2$ (Kaplan & Pikelner 1963). Hence, to produce the ionization in the whole pumping layer ($N_H \sim 10^{23-24}$ cm^{-2}) particles with $\epsilon \sim 1$–10 MeV are needed. In stationary conditions ionization rate equals recombination rate, αn_e^2, where $\alpha \sim 10^{-12}$ cm^3 s^{-1} is the recombination coefficient. Since each particle produces $\epsilon/\epsilon_i \sim 10^5$ ionizations, where $\epsilon_i \sim 30$ eV is the average energy lost at each collision, the total energy lost for ionization within the pumping region of volume $V \sim 10^{41-42}$ cm^3 will be $E_i \sim \epsilon_i \alpha n_e^2 V \sim 10^{29-30}$ erg s^{-1}, which makes up a quite real quantity – only $\sim 10^{-5}$ of the stellar wind energy transmitted to the shock 1–2.

Another source of ionization can be connected with the phenomena on the stellar surface, similar to the solar chromospheric flares. It is known (e.g. Tverskoj 1967) that the chromospheric flares are accompanied by acceleration of protons up to tens of MeV, with the total power of the flow exceeding $\sim 10^{-5}\, L_\odot$. If a similar part of the luminosity of a young massive star is transformed into the energy of MeV protons in flare-like phenomena, the required n_e/n_H would be ensured. This ionization mechanism can provide observed sporadic variability of the maser emission: the rate of the CCr pump is proportional to n_e and, if maser is not fully saturated, it must be very sensitive to n_e/n_H.

One more ionization mechanism possible in condensations with large dust grains will be mentioned in the next subsection.

4.3 MODEL 2: YOUNG STAR OF SMALL MASS

This model is based on the current ideas on the formation of planetary systems around young stars of approximately solar mass (Cameron 1978; Ward 1976) and on the possibility for such stars to produce an intense stellar wind (Norman & Silk 1979; Cohen, Bieging & Schwartz 1982; see also below).

Planetary systems are formed from the circumstellar gas–dust discs. Initial condensations necessary to form the H_2O masers could be connected with the formation of giant protoplanets. According to one of the models (Perry & Cameron 1973) the giant protoplanet is produced by accretion of the gas from the disc on a solid nucleus, generated itself by the gravitational instability in the thin equatorial layer of solid bodies (Goldreich & Ward 1973). A more recent model (Cameron 1978) connects the formation of the giant protoplanets with the gravitational instabilities in the gaseous disc itself. It is suitable that in both models a protoplanet of several Jovian masses situated at $10^{14}-10^{15}$ cm from the star has the diameter of $10^{13}-10^{14}$ cm and the average gas number density $n_H \sim 10^{10}-10^{11}$ cm^{-3} (though rising rather steeply to the centre of the protoplanet).

Following Norman & Silk (1979) we suppose that the young central star loses mass intensively during its pre-T Tauri stage. A protostar of $M_* \sim 1-2\,M_\odot$ may have the luminosity up to $L_* \sim 500\,L_\odot$ (Larson 1969). Even if the mass loss is caused only by the radiative pressure, in the optically thick case the kinetic power of the wind, $\dot{M}v_w^2/2$, can make a considerable part of L_* (e.g. Phillips & Beckman 1980). The recent radio continuum observations by Cohen *et al.* (1982) show that the kinetic power of the stellar wind in T Tauri stars can reach and even surpass their observed bolometric luminosity. Thus, it would not be unrealistic to assume for a pre-T Tauri star: $\dot{M} = 10^{-5}\,M_\odot\,\mathrm{yr}^{-1}$, $v_w = 500\,\mathrm{km\,s^{-1}}$, which, with the assumed L_*, corresponds to $\dot{M}v_w^2/2 \sim 0.3\,L_*$.

It is easily shown that the switched on stellar wind with the assumed parameters will tear off almost all the envelope of the protoplanet. The subsequent evolution of the envelope will be quite similar to the evolution of the gaseous condensation in models 1a and 1b (Fig. 3). Gas temperature behind the front 1–2 is found from equation (11) to be $T_2 \sim 3 \times 10^6$ K. If one approximates very crudely the envelope by a homogeneous condensation of radius $\sim 10^{14}$ cm and $n_H \sim 10^{10}$ cm^{-3}, then the other parameters will be of the same order as in the models 1a and 1b: $v_{34} \sim 10\,\mathrm{km\,s^{-1}}$, $T_3 \sim$ several thousand K, $v_3 \sim 1\,\mathrm{km\,s^{-1}}$, $n_3 \sim 10^{11}$ cm^{-3}, $d \sim 10^{11-12}$ cm, $\Delta T/T \sim 10^{-1}$. Hence, the sufficiently powerful CCr pumping is possible in model 2. The electrons can be provided in this model by the mechanisms described above and also by photoionization of Na (Chugai, private communication; Strelnitskij 1980b). If the dust of the condensation is presented by particles larger than ~ 1 mm in size (a rather probable situation in a protoplanetary disc), this mechanism would give $n_e/n_H \sim 4 \times 10^{-6}$ throughout the maser condensation with the cosmic gas phase Na abundance.

4 Discussion

The aim of this paper was not to propose a universal model of the H_2O maser subsources in the regions of star formation. Most probably such a task is senseless: at least the low and the high velocity H_2O condensations appear to be subjects of rather different external conditions and would have different physical states, geometries and pump mechanisms. Thus, only the particular purpose of interpreting the intrinsically most powerful subsources (ESSs) was pursued here because all the previous models seem to be unable to do it.

It was shown that the required pump rate can be provided by the CCr process in a circumstellar gas condensation pressed by the stellar wind of a young star to $n_H \sim 10^{11}$ cm^{-3} and that not only massive stars (models 1a and 1b), but also stars of solar mass can in principle give off the wind with the necessary parameters (model 2). At somewhat higher densities ($n_H \sim 10^{12}$ cm^{-3}) the CRv pumping can also provide the necessary $n_1\Delta P \sim 10^4$ cm^{-3} s^{-1}. However, the kinetic tendency (heating, or cooling of the signal transition?) is unknown for the CRv process and, due to the smallness of q_H^v, there would be a principal difficulty for all the collisional v-mechanisms in overtaking thermalization and establishing inversion (see Section 2.3). So, we prefer the CCr pump, which was shown to have a strong tendency to invert the $6_{16}-5_{23}$ transition (Bolgova 1981) and requires smaller densities.

Not only the high specific powers of the ESSs, but also many other properties of the

H_2O sources can find explanation within the limits of the proposed models. The recurrent maser activity with the typical time-scale of ~1 yr (Haschick *et al.* 1980) can be explained by the alteration of the shock and the expansion waves in the condensations: with $v_{34} \sim 10\,\mathrm{km\,s^{-1}}$ the condensation of $l \sim 10^{14}$ cm will be crossed by the shock in ~1 yr. Occasional powerful outbursts, like the outburst of the $+8\,\mathrm{km\,s^{-1}}$ in Orion (Matveenko 1981; Abraham *et al.* 1981) could have origins in the fortunate projection of that relatively thin layer where the CCr pump power reaches its maximum, or in the fortunate orientation of the line of intersection of such layers (Strelnitskij 1982).

The high velocity condensations in the centres of activity are also naturally explained as a result of acceleration by the stellar wind. It was shown by Strelnitskij & Syunyaev (1972) that a condensation with a column density $\delta\,(\mathrm{g\,cm^{-2}})$ and the initial distance r_0 from the star is accelerated by the stellar wind to a velocity $v(r) = [(r_0^{-1} - r_0^{-1})\dot{M}v_w/2\pi\delta]^{1/2}$. With the typical values in the above models of $\dot{M}v_w \sim 10^{29}$ dyn and $\delta \sim 1\,\mathrm{g\,cm^{-2}}$ this equation shows that the observed velocities $v \sim 10^7\,\mathrm{cm\,s^{-1}}$ are reached if the acceleration starts at $r_0 \sim 10^{14}$ cm. The acceleration is quick: the final velocity is reached practically within the interval $\sim 10^{14}$–10^{15} cm, after which the condensation is flying practically due to inertia up to $r \sim 10^{16}$–10^{17} cm, and this takes a time $t \sim r/v \sim 10^{9-10}$ s. Since the ram pressure drops quickly with r (more quickly than r^{-2}), the condensation would expand up to $l \sim t \cdot c_s \sim 10^{14}$–$10^{15}$ cm, i.e. by several times, with the consequent drop in the density down to $n_H \sim 10^{8-9}\,\mathrm{cm^{-3}}$. One can suppose that the condensation will taken an elongated form along the direction of the wind (Pikelner 1973; Pikelner & Strelnitskij 1976). As long as the velocity of the wind surpasses that of the condensation, there will be dissipation of the wind energy within the condensation by one or another mechanism (Norman & Silk 1979; Strelnitskij 1980a; and above) heating the gas up to several hundred kelvins. At such densities and temperatures CCr pumping will vanish, for the cooling of the electrons by excitation of the H_2O vibrations drops quickly with T_e (equation 12) and T_e approaches T_H.

More probable for the high velocity condensations is the CRr pump with the possible sink of the rotational photons on the cold dust. Calculations show that with, for example, $n_1 \sim 5 \times 10^{-9}\,n_H$, $n_H \sim 5 \times 10^8\,\mathrm{cm^{-3}}$, $T_H \sim 500$ K, $T_d \sim 100$ K the observed fluxes from the high velocity features are reached with $l \sim 10^{15}$ cm, which is close to the expected value in the above model and is consistent with the VLBI of the high velocity features (Heckman & Sullivan 1976).

The strong stellar wind, which pumps the ESSs according to the proposed models, would sweep the ESSs away from the star and so must limit the 'ESS phase' of a new-born star by approximately $\tau \sim 5 \times 10^{14}\,\mathrm{cm}\;10^{-6}\,(\mathrm{cm\,s^{-1}}) \sim 5 \times 10^8\,\mathrm{s} \sim 20$ yr. According to the current estimates of the star formation rate R (e.g. Massevitch *et al.* 1976), ~0.4 stars of $M_* \sim 1$–$2\,M_\odot$ and ~0.06 stars of $M_* \gtrsim 5M_\odot$ are born in the Galaxy per year. So, model 2 predicts the total number of ESSs in the Galaxy $N \sim R_{1-2}\tau \sim 10$, whereas model 1 gives $N \sim R_{>5}\,\tau \sim 1$. Since we always see at any moment ~10 ESSs the former model seems to be more adequate from this point of view.

It is doubtful, however, that all the H_2O sources observed in the regions of star formation, would originate in such non-stationary conditions. Typical lifetimes of the centre of activity would then be $\tau_{CA} \sim 10^{17}\,\mathrm{cm}\;\;10^{-7}(\mathrm{cm\,s^{-1}}) \sim 10^{10}\,\mathrm{s} \sim 300$ yr, and the predicted total number of such centres, observed at a given time, $N_{CA} \sim \tau_{CA}R_{tot} \sim 10^2$, which is comparable with the number of detected H_2O sources but probably one order of magnitude less than the total number of existing sources (Genzel & Downes 1977). So we do not exclude that a part of the H_2O sources originates in quasi-stationary conditions, like the quasi-stationary discs suggested by Elmegreen & Morris (1979).

The control experiment of choosing between these two possibilities for the ESSs would be the measurement of their proper motions within the centres of activity: if the non-stationary models are correct, the proper motions of the ESSs must be of radial type and with a secular acceleration. If, on the contrary, emission of the ESSs originates in quasi-stationary discs, the sub-sources would show no secular proper motions.

More thorough statistical investigations of the association between the H_2O masers and the compact infrared and H II sources would help to draw conclusions about the masses of the central stars.

References

Abraham, Z., Cohen, N. L., Opher, R. & Rafaelli, J. C., 1981. *Astr. Astrophys.*, **100**, L10.

Bolgova, G. T., 1981. *Nauchnye Informacii Astrosovijeta AN SSSR*, **47**, 9.

Bolgova, G. T., Strelnitskij, V. S. & Umanskij, S. Ya., 1982. *Nauchnye Informacii Astrosovijeta AN SSSR*, **51**, 22.

Borodina, E. G., 1975. *Solnechnye Dannye (GAO AN SSSR), Bull.*, **6**, 80.

Cameron, A. G. W., 1978. In *Protostars and Planets*, p. 453, ed. Gehrels, T., University of Arizona Press, Tucson, Arizona.

Cohen, M., Bieging, J. H. & Schwartz, P. R., 1982. *Astrophys. J.*, **253**, 707.

Deguchi, S., 1981. *Astrophys. J.*, **249**, 145.

de Jager, C., 1980. *The Brightest Stars*, Reidel, Dordrecht, Holland.

de Jong, T., 1973. *Astr. Astrophys.*, **26**, 297.

Downes, D., Genzel, R., Moran, J. M., Johnston, K. J., Matveyenko, L. I., Kogan, L. R., Kostenko, V. I. & Rönnäng, B., 1979. *Astr. Astrophys.*, **79**, 233.

Downes, D., Genzel, R., Becklin, E. & Wynn-Williams, C., 1981. *Astrophys. J.*, **244**, 869.

Drukarev, G. F., 1978. *Collisions of the Electrons with Atoms and Molecules*, Nauka, Moscow.

Elmegreen, B. G., 1978. *The Moon and the Planets*, **19**, 261.

Elmegreen, B. G. & Morris, M., 1979. *Astrophys. J.*, **229**, 593.

Forster, J. R., Welch, W. J. & Wright, M. C. H., 1977. *Astrophys. J.*, **215**, L121.

Forster, J. R., Welch, W. J., Wright, M. C. H. & Baudry, A., 1978. *Astrophys. J.*, **221**, 137.

Genzel, R. & Downes, D., 1977. *Astr. Astrophys. Suppl.*, **30**, 145.

Genzel, R., Downes, D., Moran, J., Johnston, K. J., Spencer, J. H., Walker, R. C., Haschik, A., Matveyenko, L. I., Kostenko, V. I., Rönnäng, B., Rydbeck, O. E. H. & Moiseev, I. G., 1978. *Astr. Astrophys.*, **66**, 13.

Genzel, R., Downes, D., Moran, J. M., Johnston, K. J., Spencer, J. H., Matveyenko, L. I., Kogan, L. R., Kostenko, V. I., Rönnäng, B., Haschick, A. D., Reid, M. J., Walker, R. C., Giuffrida, T. S., Burke, B. F. & Moiseev, I. G., 1979. *Astr. Astrophys.*, **78**, 239.

Giuffrida, T. S., Greenfield, P. E., Burke, B. F., Haschick, A. D., Moran, J. M., Rydbeck, O. E., Rönnäng, B. O., Bääth, L., Invesson, K. S., Matveyenko, L. I., Kostenko, V. I., Kogan, L. R. & Moiseev, I. G., 1981. *Pis'ma v Astron. Zh.*, **7**, 358.

Goldreich, P. & Keeley, D. A., 1972. *Astrophys. J.*, **174**, 517.

Goldreich, P. & Kwan, J., 1974. *Astrophys. J.*, **191**, 93.

Goldreich, P. & Ward, W. R., 1973. *Astrophys. J.*, **183**, 1051.

Haschick, A. D., Moran, J. M., Rodriguez, L. F., Burke, B. F., Greenfield, P. & Garcia-Barreto, J. A., 1980. *Astrophys. J.*, **237**, 26.

Heckman, T. M. & Sullivan III, W. T., 1976. *Astrophys. Lett.*, **17**, 105.

Itikava, I., 1972. *J. Phys. Soc. Japan*, **32**, 217.

Ivanov, V. V., 1969. *Radiative Transfer and Spectra of Celestial Bodies*, Nauka, Moscow.

Kaplan, S. A. & Pikelner, S. B., 1963. *Mezhsvesdnaja Sreda*, Fismatgis, Moscow. (English translation: 1970. *The Interstellar Medium*, Harvard University Press, Cambridge, Massachusetts.)

Kaplan, S. A. & Tsytovich, V. N., 1972. *Plazmennaja Astrofisika*, Nauka, Moscow.

Lada, C. J., Blitz, L., Reid, M. J., Moran, J. M., 1981. *Astrophys. J.*, **243**, 769.

Larson, R., 1969. *Mon. Not. R. astr. Soc.*, **145**, 271.

Litvak, M. M., 1969. *Science*, **165**, 855.

Massevitch, A. G., Tutukov, A. V., Yungelson, L. R., 1976. *Astrophys. Space Phys.*, **40**, 115.

Matveenko, L. I., 1981. *Pis'ma v Astron. Zh.*, **7**, 100.

Mihalas, D., 1978. *Stellar Atmospheres*, W. H. Freeman & Co., San Francisco.

Norman, C. & Silk, J., 1979. *Astrophys. J.,* **228,** 197.

Perry, F. & Cameron, A. G. W., 1974. *Icarus,* **22,** 416.

Phillips, J. P. & Beckman, J. E., 1980. *Mon. Not. R. astr. Soc.,* **193,** 245.

Pikelner, S. B., 1973. *Astrophys. Lett.,* **15,** 91.

Pikelner, S. B. & Strelnitskij, V. S., 1976. *Astrophys. Space Sci.,* **39,** L19.

Schneps, M. H., Lane, A. P., Downes, D., Moran, J. M., Genzel, R. & Reid, M. J., 1981. *Astrophys. J.,* **249,** 124.

Seng, G. & Linder, F., 1976. *J. Phys. B.,* **9,** 2539.

Shmeld, I. K., Strelnitskij, V. S. & Muzylev, V. V., 1976. *Astr. Zh.,* **53,** 728.

Spitzer, L., Jr., 1978. *Physical Processes in Interstellar Medium,* J. Wiley & Sons, New York.

Strelnitskij, V. S., 1971. *Astr. Tsirk.,* **609,** 1.

Strelnitskij, V. S., 1973. *Astr. Zh.,* **50,** 1133.

Strelnitskij, V. S., 1974. *Usp. Fis. Nauk,* **113,** 463. (*Sov. Phys. Uspechy,* **17** (4), 507, 1975).

Strelnitskij, V. S., 1980a. In *Interstellar Molecules, IAU Symp. 87,* p. 591, ed. Andrew, B. H., Reidel, Boston, London.

Strelnitskij, V. S., 1980b. *Pis'ma v Astron. Zh.,* **6,** 354.

Strelnitskij, V. S., 1981. *Astr. Zh.,* **58,** 660.

Strelnitskij, V. S., 1983. *Nauchnye Informacii Astrosovijeta AN SSSR,* **52,** 75.

Strelnitskij, V. S., 1982. *Pis'ma v Astron. Zh.,* **8,** 165.

Strelnitskij, V. S. & Syunyaev, R. A., 1972. *Astr. Zh.,* **49,** 704.

Turner, B. E., 1970. *J. R. astr. Soc. Canada,* **64,** 221; 282.

Tverskoj, B. A., 1967. *Zh. Exp. Teor. Fis.,* **53,** 1417.

Walker, R. C., Matsakis, D. N. & Garcia-Barreto, J. A., 1982. *Astrophys. J.,* **255,** 128.

Ward, W. R., 1976. In *Frontiers of Astrophysics*, p. 1, ed. Avrett, E. H., Harvard University Press, Cambridge, USA and London.

Zel'dovich, Ya. B. & Raizer, Y. P., 1966–67. *Physics of Shock Waves and High-Temperature Phenomena.* Academic Press, New York.