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SS433: on the uniqueness of cool relativistic jets

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Abstract. The relativistic jets of SS433 are outstanding for their optical thermal radiation, which is produced by small clouds (10^8 cm) with lifetimes about 10^3 times larger than the gas-dynamical crushing time. We show that the clouds reside in thermal and dynamical balance as long as they collide with the wind of the supercritical accretion disk. This interaction is caused by the precessional movement of the jets and takes place only in the sweep-out zone, beyond which the interaction ceases and optical jets merely terminate. The cloud's magnetic field, amplified in course of movement through a medium, could play a role in containing a cloud. This interaction provides the clue to the uniqueness of the optical jets of SS433.

Key words: stars: individual: SS 433 – ISM: jets and outflows

1. Introduction

Since the discovery of jets in SS433 in 1978, there has been discussion about their uniqueness. By now the choice of jets has enlarged into a veritable zoo as hundreds of relativistic jets of AGNs, low velocity jets of young stars and a dozen jets of compact stars were detected. Jets of SS433 stand out among the others, for they harbour cool clouds within the relativistic flow. It is not clear, however, how the clouds could survive in the relativistic jets and be observed. At the same time, the thermal fraction is likely to be dominant in the jets of SS433, whereas, the lack of observational evidence of thermal gas in AGN jets strongly restricts its volume-filling factor ($< 10^{-8}$) and implies that it cannot be important in the energy budget of their jets (Celotti et al. 1998). All these peculiarities make the jets of SS433 unique.

SS433 jets are highly collimated ($\theta_{\rm j} \leq 1^\circ.4$), mildly relativistic ($v_{\rm j} = 7.8\cdot 10^9\,{\rm cm\,s^{-1}}$), and powerful, with kinetic luminosity $L_{\rm k} \approx 10^{39}\,{\rm erg\,s^{-1}}$. They appear from funnels of the supercritical accretion disk and follow its precession rotation with a half opening angle of about 20° and a period of $P_{\rm pr} = 162^{\rm d}.5$ (see Fig. 1). At the distance of a few 10^{12} cm, the jets cool rapidly and radiate in the X-ray domain. The X-ray emission does not show any structure to the jets, so that the X-ray jets are thought to be continuous (Stewart et al. 1987). In contrast, the optical jets are observed to be fragmented (Vermeulen et al. 1993b). Their clumpiness may be roughly deduced from energetic considerations, as follows. The gas must be dense enough to radiate the bulk of the emission:

$$L_{\rm H_\alpha} = \epsilon_{\rm H_\alpha} n_{\rm e}^2 f\,V, \tag{1}$$

and mass loading of the jet by dense gas must approximately satisfy:

$$L_{\rm k} = \dot{M}_{\rm j} v_{\rm j}^2, \tag{2}$$

where $V \approx \theta_{\rm j}^2 R^3/4$ is the volume of the jet, R its length, $n_{\rm e}$ the electron density, and $\epsilon_{\rm H_\alpha}$ the $\rm H_\alpha$ emission rate. Thus, one finds both $n_{\rm e} \approx 5\cdot 10^{11}\,{\rm cm^{-3}}$ and a volume-filling factor of $f \approx 10^{-5}$ for appropriate parameters: luminosity in $\rm H_\alpha$ line $L_{\rm H_\alpha} = 10^{36}\,{\rm erg\,s^{-1}}$; total jet kinetic luminosity $L_{\rm k} = 10^{39}\,{\rm erg\,s^{-1}}$; $\epsilon_{\rm H_\alpha} = 10^{-24}\,{\rm erg\,cm^3\,s^{-1}}$, $R = R_{\rm e} \approx 8.4\cdot 10^{14}$ cm – length of *e*-folding $\rm H_\alpha$ line brightness. It seems that the jets become lumpy in the gap between the X-ray part and the optical jet beginning at about $R_{\rm in} = 2.3\cdot 10^{14}$ cm (Borisov & Fabrika 1987). The gas could condense there due to thermal instability as Brinkmann et al. (1988) have shown in a simulation of the thermal evolution of the X-ray jets. The instability takes place only in the interior jet, up to a distance of 10^{13} cm. The parameters of the clouds found in their work are close to those obtained by Panferov & Fabrika (1997) from the Balmer decrements of the jets: hydrogen density $n \geq 10^{13}\,{\rm cm^{-3}}$ and size of the clouds $l \leq 10^8$ cm. It should be noted that these parameters, as well as the filling factor $f \approx 4\cdot 10^{-6}$ found there, are more appropriate for the distance of maximum optical brightness of the jets $R_{\rm m} \approx 4\cdot 10^{14}$ cm. These parameters will be representative ones in what follows. The clouds are grouped in clusters about 10^{12} cm in size (Panferov & Fabrika 1997). Even larger structures are quite possible in the jets as may be deduced from the structure of the jet emission line.

Brown et al. (1995) consider cloudy structure of the jets as a result of a temporal switch on and off of conditions for thermal instability. They found that the parameters of SS 433 are just suitable for the instability to occur. However, the instability can work with a wide range of the parameters and, in SS 433 jets, the conditions for cloud generation exist rather continuously in time. Indeed, the data on optical emission of the jet (e.g. Vermeulen et al. 1993b) demonstrate that the generation of the clouds in the jets is continuous and the disappearance of the optical jets happens very rarely. The intermittent structure of the jets is natural if the radiating clouds form in such *in situ* variable process as the thermal instability. We see how separated emission peaks grow and decay at an unchanged wavelength, reflecting an evolution of the "bullet" along the jets (Borisov & Fabrika 1987; Vermeulen et al. 1993b). Thus, the clouds form in the interior jets and continue to exist a long time, providing a signature of the optical jets. They terminate at a distance of about $R_\mathrm{j} = 3 \cdot 10^{15}$ cm (Borisov & Fabrika 1987) and then are observed to be quasi-continuous in radio (Vermeulen et al. 1993a).

The actual problem is the determination of how clouds evolve and survive in the relativistic jets? As long as the clouds are involved in a motion through a medium, as in the dense wind of the accretion disk, the applied ram pressure leads to gas-dynamical instabilities that try to disrupt the clouds on a crushing time scale (e.g., Jones et al. 1996):

$$t_\mathrm{cr} = \frac{l}{v_\mathrm{j}}\sqrt{\frac{\rho}{\rho_\mathrm{w}}}, \tag{3}$$

where ρ is the cloud density, and ρ_w is the wind density. For the clouds of the jets of SS 433 this time is in the order of 10^2 s. Besides this, the clouds undergo an evaporation in the hotter intercloud medium. The evaporation's characteristic time may be rather short, too (see Eq. (16)). Namely the capability of the clouds to survive gives us an opportunity to observe the jets of SS 433 in the optical. In this paper we show that the collisional interaction that seems to be the destructive one is a most necessary condition for the clouds' existence.

In Sect. 2 we show that the collisional interaction may be comparably important or even more essential than radiation for the thermal state of the clouds. Although this has been concluded by other authors (e.g., Brown et al. 1995), we use here the latest data on the jets' parameters (Panferov & Fabrika 1997) and on the parameters of the UV radiation of the system – a possible source of heat for the clouds (Dolan et al. 1997). In Sect. 3 we discuss the evolution and confinement of the clouds.

2. Thermal balance of clouds

There are two main ways to heat the clouds: by interaction of jets with the gaseous wind of the supercritical accretion disk or by radiation of the system (Davidson & McCray 1980). Here we consider the competition between these two mechanisms. As long as the cloud sound-crossing time is shorter than the dynamical time-scale $t_\mathrm{d} = r/v_\mathrm{j}$ of the jet, the clouds are in pressure equilibrium with their surroundings. The influence of the ram pressure of the wind on the equilibrium is most essential (see Sect. 3). Density of the wind scales as $\propto r^{-2}$, where r is the radial distance from the source of the jets. Under the plausible condition of approximate constancies of temperature and mass of a cloud, this implies the following dynamical dependencies of the cloud's parameters:

$$\begin{aligned} n &= 1.6 \cdot 10^{14}\, r_{14}^{-2}\ \mathrm{cm}^{-3}, \\ l &= 4 \cdot 10^{7}\, r_{14}^{2/3}\ \mathrm{cm}, \\ f &= 1.6 \cdot 10^{-5}\, r_{14}^{-1}, \end{aligned} \tag{4}$$

where $r_{14} = r/(10^{14}\ \mathrm{cm})$. We have adopted the values of the cloud parameters at R_m, given in Sect. 1. In the following, we assume Eqs. (4) to be valid everywhere along the optical jets.

Let us compare the rates of cloud heating by the impinging gas and by radiation. The jets sweep the gas out of the wind thanks to the change of the jet's direction in the course of the precession and to nodding motions and small amplitude jitters of unclear nature (Margon & Anderson 1989). The energy deposition into a cloud by crossing protons with energy $\epsilon_\mathrm{p} = m_\mathrm{p} v_\mathrm{j}^2/2$ depends on the collisional thickness, measured by the ratio $\sigma_\mathrm{c} = N/N_\mathrm{s}$, where $N = n\,l = 6.4 \cdot 10^{21} r_{14}^{-4/3}$ cm^{-2} is the column density of the cloud, and N_s is the column density required to stop a proton. The rate of energy loss of a high energy proton in an ionized gas is approximately (Ginzburg & Syrovatskii 1964):

$$\Gamma_\mathrm{p} = \frac{2.2 \cdot 10^{-8}\, K_\mathrm{c} n}{v_\mathrm{j}} \left(1 + 1.21 \frac{m_\mathrm{p}}{m_\mathrm{e}} \frac{k\,T}{\epsilon_\mathrm{p}}\right)^{-3/2} \mathrm{erg\,s}^{-1}, \tag{5}$$

for temperatures $k\,T \ll m_\mathrm{e} c^2$, $k\,T \ll \epsilon_\mathrm{p}$. $K_\mathrm{c} = 30\,\Lambda/n_\mathrm{e} n \sim 1$, where Λ is the Coulomb logarithm. In an ionized plasma almost all of the lost energy goes into heating the electron gas. The stopping column density of a proton with energy $\epsilon_\mathrm{p} = 32$ MeV is:

$$N_\mathrm{s} = n\, v_\mathrm{j} \frac{\epsilon_\mathrm{p}}{\Gamma_\mathrm{p}} \approx 1.4 \cdot 10^{23}\ \mathrm{cm}^{-3}. \tag{6}$$

When $N < N_\mathrm{s}$, which is satisfied for the entire jet (see Eq. (9)), the cloud heating rate by the fast protons is:

$$H_\mathrm{c} \approx \sigma_\mathrm{c}\, \epsilon_\mathrm{p} n_\mathrm{w} v_\mathrm{j} l^2 = \epsilon_\mathrm{p} n_\mathrm{w} v_\mathrm{j} n\, l^3/N_\mathrm{s}. \tag{7}$$

This is the lowest estimate. The magnetic field may be amplified to an equipartition value at the surface of a cloud moving through a medium with a weak magnetic field (Jones et al. 1996). In the magnetic field of Eq. (18) the gyro-radius of a 32 MeV proton is $3 \cdot 10^5\, r_{14}$ cm. This is smaller than the cloud's radius, therefore the trajectories of protons may be tangled and the heat input boosted.

The clouds could also be heated by the collimated radiation from funnels of the accretion disk (Arav & Begelman 1993; Panferov & Fabrika 1993) and by the disk UV radiation. The heating is realized by δ-electrons – high energy electrons broken away from an atom due to ionization. Already for ionization degree ≥ 0.5, almost all the energy of δ-electrons is utilized in the heating of the electron gas of a cloud. We take into account radiation only with frequency beyond the Lyman edge. A cloud intercepts radiation flux with density $j_\nu = \pi(R_\star/r)^2 B_\nu$, where $R_\star$ is the radius of the radiation source, B_ν is the Planck function, and relativistic factor $\gamma^{-3} = (1 - v_\mathrm{j}^2/c^2)^{3/2} \sim 1$ is omitted. The part of the radiation which is absorbed is $1 - e^{-\tau}$, where the cloud's optical thickness is $\tau = \sigma\, n_\mathrm{H}\, l = \sigma\,(1-\zeta)\, n\, l$, and the ionized fraction is $\zeta = n_\mathrm{p}/(n_\mathrm{p} + n_\mathrm{H})$. Then the radiative heating rate of one cloud is:

$$H_\mathrm{r} = l^2 \int_{\nu_1} \mathrm{d}\nu\, j_\nu (1 - e^{-\tau_\nu}), \tag{8}$$

where ν_1 is the frequency of the Lyman edge. We adopt here $\sigma_\nu \approx 1.73 \cdot 10^{-22} \epsilon_{\rm keV}^{-8/3}\,{\rm cm}^2$, for a photon with energy $\epsilon = 1{\rm keV} \cdot \epsilon_{\rm keV}$ in the range 13.6 eV – 12.4 keV, for cosmic abundances (Cruddace et al. 1974).

Parameters of the collimated radiation are not well known since its orientation to the Earth is unfavourable. Here we adopt rather rough estimates for its luminosity $L_{\rm X} \sim 10^{39}\,{\rm erg\,s^{-1}}$ and temperature $T_{\rm X} \sim 1\,{\rm keV}$ (Arav & Begelman 1993; Panferov & Fabrika 1993). Then the radius of a source of collimated radiation in blackbody approach is $R_{\rm X} \approx 2 \cdot 10^7\,{\rm cm}$. The supercritical accretion disk is a source of near isotropic UV radiation, its parameters in blackbody approach are: $T_{\rm UV} \approx 72000\,{\rm K}$ and $R_{\rm d} \approx 1.5 \cdot 10^{12}\,{\rm cm}$ (Dolan et al. 1997). For given parameters and $\tau < 1$, Eq. (8) gives $H_{\rm r}^{\rm UV}/H_{\rm r}^{\rm X} \approx 3 \cdot 10^5$, thus, heating by UV radiation of the disk could dominate heating by the X-ray collimated radiation. This ratio might be bigger because the clouds have $\tau > 1$ for the UV radiation, whereas $\tau < 1$ for the X-ray radiation. However, the situation is further complicated by the screening of the clouds in a jet.

The cross-over path length of a particle moving in the direction of the jet velocity across the jet is determined by the curvature of the jet. The motion of the jet caused by nodding is more rapid than that caused by precession and its rate is $\dot\phi \approx 5.7 \cdot 10^{-7}\,{\rm rad\,s^{-1}}$ (Borisov & Fabrika 1987). Using this rate we obtain the cross-over length for a jet with fixed pattern $l_{\rm f} = v_{\rm j}\theta_{\rm j}/\dot\phi = 3.4 \cdot 10^{14}\,{\rm cm}$, with an accuracy of factor 2. Column density along this path is:

$$N_{\rm f} = \int_r^{r+l_{\rm f}} {\rm d}r\, f\, n \approx \frac{1.3 \cdot 10^{23}}{r_{14}^2} \times \left(1 - \frac{r_{14}^2}{(r_{14} + l_{\rm f\,14})^2}\right)\ {\rm cm}^{-2}. \qquad (9)$$

Consequently, we have $N_{\rm f} < N_{\rm s}$ everywhere in the optical jets and the jets are transparent to the rapid protons. So, the screening effect is not important for the collisional heating. Meanwhile the screening factor of clouds in the jet $N_{\rm f}/N$ may be essential for the UV radiation. We now determine $N_1 = 1/\sigma_\nu(1-\zeta)$ as the column density of a cloud layer where $\tau = 1$ for the UV radiation. Then the screening factor for the UV radiation is $N_{\rm f}/N_1 \sim 10^6(1-\zeta)$. This essentially depends on the ionization degree and may be much larger than 1. Therefore we will consider the heating $H_{\rm r}^{UV}$ of Eq. (8), of the cloud by the UV radiation, as an upper limit of the radiative heating. Really, this heating is $N_{\rm f}/N_1$ times stronger than the averaged over the whole jet heating rate of one cloud; the opacity of the jet to UV radiation causes a dominance of the collimated radiation in the heating of screened clouds.

The ratio of the heating rate from fast protons (Eq. (7)) to the radiation heating rate (Eq. (8)) is:

$$\frac{H_{\rm c}}{H_{\rm r}^{\rm UV}} = \frac{\epsilon_{\rm p} v_{\rm j} n\, l\, n_{\rm w} r^2}{\pi\, N_{\rm s} R_\star^2 \int_{\nu_1} {\rm d}\nu B_\nu^{\rm UV}} \approx 10\, r_{14}^{-4/3}, \qquad (10)$$

where we have used the reduced form of Eq. (8) for $\tau > 1$ (substituting $1 - e^{-\tau_\nu}$ for 1) and the density of wind of the accretion disk:

$$n_{\rm w} = \frac{\dot M_{\rm w}}{4\,\pi\, m_{\rm p}\, r^2 v_{\rm w}} \approx 1.5 \cdot 10^8\, r_{14}^{-2}\ {\rm cm}^{-3}. \qquad (11)$$

The mass loss rate and wind velocity are taken to be $\dot M_{\rm w} = 10^{-4}\, M_\odot/{\rm y}$ (van den Heuvel 1981) and $v_{\rm w} = 2000\,{\rm km/s}$ (from the width of the emission lines) respectively. Thus, the heating of the jet by fast protons is at least comparable to the radiation heating. Their rates become equal at a distance of $5.6 \cdot 10^{14}\,{\rm cm}$, up to where as much as 64% of the heat input into the whole jet given by Eq. (13) take place. The collisional heating may be much larger than the radiative one, because the value of $H_{\rm r}^{\rm UV}$ is overestimated approximately $10^6(1-\zeta)$ times due to the screening effect.

This idea concerning the collisional heating of the jet clouds is verified by observations: the H_α emission of the jets is anisotropic, the maximum of its directional pattern is directed to the side of the jet movement (Panferov et al. 1997). This implies that some screening effect exists for the impinging protons. This is possible if a magnetic field tangles the trajectories of the protons in a cloud resulting in the smaller stopping column density $N_{\rm s}$ than that given by Eq. (6). However, the directional pattern can not be explained straightforwardly as a result of a head-on collision of the clouds with the wind, because the direction of maximal radiation is inclined to the jet axis. The inclination of the maximum is about 40° from the jet axis in the direction of the precessional motion. This tends to show that a maximum outcome of the emission line radiation is from jet sides moving in the wind.

As a result of collision with the wind, a cloud decelerates at a rate:

$$\frac{\Delta v}{v} \approx \frac{\Delta M}{M} = \int {\rm d}t \frac{\sigma_{\rm c} n_{\rm w}\, v_{\rm j}}{n\, l} = \int_{R_{\rm m}}^{R_{\rm j}} {\rm d}r \frac{n_{\rm w}}{N_{\rm s}} \approx 2 \cdot 10^{-2}. \qquad (12)$$

Here we used $R_{\rm m}$ as the lower limit of the integration because evolution of the jet at distances $< R_{\rm m}$ is not well established. The deceleration (12) agrees well enough with the observed value of the deceleration $\leq 10^{-2}$ (Kopylov et al. 1987). Jets sweep wind gas out with a pattern speed of $\dot\phi\, r$. Then the rate of the heating of the whole of the optical jet provided by fast protons is:

$$H = \int_{R_{\rm in}}^{R_{\rm j}} {\rm d}r\, \frac{N_{\rm f}}{N_{\rm s}} \epsilon_{\rm p} n_{\rm w} (\dot\phi\, r)(\theta_{\rm j} r) \approx 2.2 \cdot 10^{37}\ {\rm erg\,s^{-1}}. \qquad (13)$$

This integral is calculated numerically and the dependence of $N_{\rm f}$ on a geometry of the precessing conical jet is taken into account. The calculated value of the whole heat H agrees with the observable radiation output of the jets in H_α line $L_{{\rm H}_\alpha} \approx 10^{36}\,{\rm erg\,s^{-1}}$ (Panferov et al. 1997): a part < 0.1 of absorbed energy is emitted in H_α.

The small clouds in question will overheat and their H_α radiation will cease unless the clouds are dense enough to establish a balance between radiative losses and the heating by the fast protons. In the temperature region near 10^4 K, which is the threshold ionization temperature of hydrogen, the efficiency of the radiative cooling $\lambda(T)$ strongly depends on temperature and changes by 3 orders of magnitude: $\lambda = 10^{-25}$ – $10^{-22}\,\mathrm{erg\,cm^3\,s^{-1}}$ (Kaplan & Pikelner 1979). From this and Eq. (4) one finds limits of possible values of the radiative energy loss of the clouds, corresponding to the end and the beginning of the jet:

$$3\cdot 10^{-3}\,\mathrm{erg\,cm^{-3}\,s^{-1}} \le \lambda n^2 \le 9\cdot 10^4\,\mathrm{erg\,cm^{-3}\,s^{-1}}. \quad (14)$$

Since limits on the heating rate H_c from Eq. (7) correspond to:

$$9\cdot 10^{-2}\,\mathrm{erg\,cm^{-3}\,s^{-1}} \le H_\mathrm{c}/l^3 \le 3\cdot 10^3\,\mathrm{erg\,cm^{-3}\,s^{-1}}, \quad (15)$$

for the jet end and beginning respectively, the heating rate never can exceed the cooling capability, given by inequality (14), and the thermal balance of the clouds can be maintained near a temperature of 10^4 K throughout the whole of the optical jets.

3. Evolution and confinement of clouds

The lifetime of the clouds in the jets is about 4 days. This is much longer than the cloud sound-crossing time. Therefore, the clouds must be in pressure equilibrium with the ambient medium and should not be exposed to destructive processes. The pressure of the ionizing radiation $p_\mathrm{r} = H_\mathrm{r}^\mathrm{UV}/l^2 c \approx 9\,r_{14}^{-2}$ dyn is much smaller than the gas pressure in a cloud $p = 2\,n\,k\,T \approx 10^3\,r_{14}^{-2}$ dyn, where we use a temperature of $T = 20000$ K (Panferov & Fabrika 1997). Therefore, radiation is not important for the dynamical balance of a cloud. The pressure equilibrium of clouds with the surroundings imposes constraints on the density $n_\mathrm{h} = n\,f \sim 10^7\,\mathrm{cm^{-3}}$ and on the temperature $T_\mathrm{h} = T\,n/n_\mathrm{h} \sim 10^{10}$ K of the surrounding gas at a distance R_m. In this hot medium, the clouds will evaporate with a timescale of:

$$t_\mathrm{ev} = \frac{m}{\dot m} = \frac{\rho\, l}{6\,\rho_\mathrm{h}\, c_\mathrm{h}\, F(\sigma_0)} \ll 10^4\ \mathrm{s}, \quad (16)$$

where m is the mass of the cloud, c_h is the sound velocity in the hot medium, $F(\sigma_0)$ is some function (given by Eq. (62) in Cowie & McKee 1977) which is $\gg 1$ in our case. The evaporation of the clouds with such a short timescale is in contradiction with the lifetime of the clouds. Therefore the clouds' evaporation must be suppressed. On the other hand, the ram pressure of the wind is important for the pressure balance of the clouds: $p_\mathrm{w}/p = \rho_\mathrm{w}\, v_\mathrm{j}^2/2\,n\,k\,T \approx 17$. The ram pressure of the wind gas seems to be strongest and must govern the clouds' pressure. But again, gas-dynamical instabilities are capable of disrupting the clouds on a timescale given by Eq. (3). This is the puzzle of SS 433 jets: how can thermal gas survive there?

The optical jets end at a distance of about $3\cdot 10^{15}$ cm. Further, at a distance of $3.7\cdot 10^{15}$ cm, the jets brighten in radio which is a consequence of jet expansion (Vermeulen et al. 1987). It seems that this expansion is connected to the termination of the optical jet. As long as the jets precess and move through the dense wind of the accretion disk, they sweep the wind gas out to the surface of the precession cone. After the passage of the jets, the

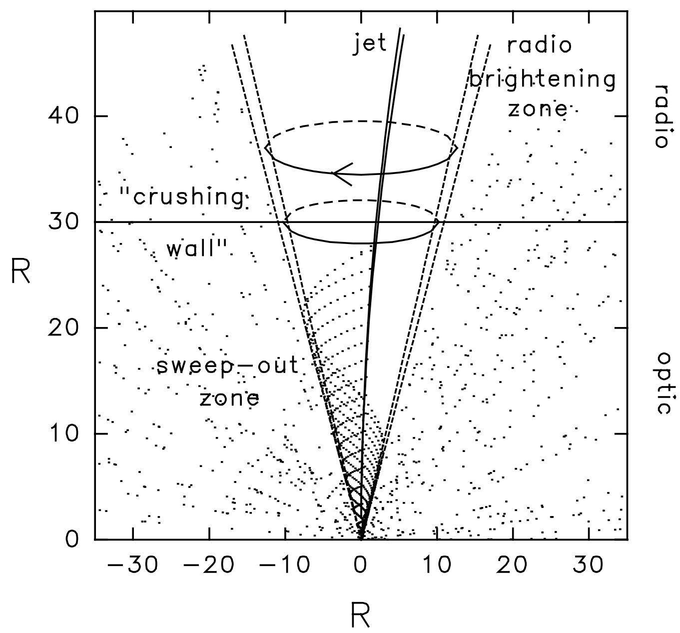

Fig. 1. Sketch of the precessing jet of SS 433. Horizontal axis lies in the accretion disk plane. Both axes are in units of 10^{14} cm. Scales are exact. The wind of the disk fills all space except the jet precession surface after distance $R_\mathrm{sw} \approx 3\cdot 10^{15}$ cm, that is the distance of the "crushing wall". And in bounds of the sweep-out zone the wind fills the jet precession surface only along a helical pattern moving with the precession period. The expansion of the optically radiating clouds occurs beyond the "crushing wall" between two rings.

wind fills the jet channels, and, during the precessional period, a region of length $R_\mathrm{sw} = P_\mathrm{pr} v_\mathrm{w} \approx 3\cdot 10^{15}$ cm can be refilled. Consequently, at distances larger than the sweep-out distance, the lobe along the precession cone surface is empty (see Fig. 1). Thermal diffusion of wind gas is not capable of filling the lobe: $c_\mathrm{h} P_\mathrm{pr} \ll \theta_\mathrm{j} r$ at $r > R_\mathrm{sw}$. The dependencies (4) of the clouds' parameters on distance and the clouds' heating balance have no peculiarities in point $r = R_\mathrm{sw}$. But the conditions in the clouds' ambient medium change abruptly. Possibly, this forces the confinement of the clouds to switch off and results in the expansion of the jet.

From laws in Eqs. (4) we derive the approximate size of the cloud $l \approx 4\cdot 10^8$ cm and the filling factor $f \approx 5\cdot 10^{-7}$ at the end of the optical jets. With these parameters the clouds start to expand freely and fill the whole volume of the jets in time:

$$t_\mathrm{ex} \sim \frac{l}{c_\mathrm{s} f^{1/3}} \approx 0^\mathrm{d}.5, \quad (17)$$

where c_s is the sound velocity in the cloud. The estimated time of expansion corresponds to the time interval required for the jet to move between the end of the optical part and the radio brightening zone. The latter could develop due to the increase of the rate of generation of relativistic particles, when the clouds expand and turbulence is enhanced.

The *coincidence* between the lengths of the optical jets and the sweep-out zone indicates that the ram pressure of the wind gas causes confinement of the clouds. The ram pressure switch off is equivalent to a "hard wall" which crushes the clouds. However, the ram pressure by itself is not able to prevent clouds from stripping off the gas and thermal evaporation. On the contrary, it seems from Eq. (3) that the higher the ram pressure, the shorter the lifetime of a cloud. In reality, ionized clouds and their surroundings harbour magnetic fields, which should be important for the clouds' stability. The magnetic field inferred from radio observations is $B \approx 0.08$ G at the distance of radio brightening (Vermeulen et al. 1987). Magnetic flux constancy implies that B scales as r^{-1}, and we assume that the magnetic field in the jet is:

$$B = 3\, r_{14}^{-1}\ \mathrm{G}. \tag{18}$$

This field could not confine the clouds to having internal pressure $\sim 10^3 r_{14}^{-2}$ dyn. Jones et al. (1996) showed in their numerical simulations that a "magnetic shield" forms around a cloud moving through a medium. The field lines of the medium are caught by a cloud and stretched and the magnetic field is amplified as a result. The "shield" has a magnetic pressure comparable to the ram pressure and is able to quench gas-dynamical instabilities and prevent evaporation of a cloud. Energy of the "magnetic shield" is converted from the kinetic energy of the cloud. So, the magnetic field needed for the confinement of the clouds in the jets of SS 433 can be generated by the motion of the clouds through the wind of the accretion disk. This mechanism of confinement naturally explains the termination of the optical jet beyond the "crushing wall" at a distance $R_{\rm sw}$, where the collisional interaction of the clouds with the wind ceases.

4. Conclusions

We have considered some consequences of the collisional interaction of clouds in the jets of SS 433 with the wind of the accretion disk. This interaction appears to be a determining factor of the clouds' state and evolution. This is emphasized by the fact that the clouds exist only in boundaries of the sweep-out zone over which the jet sweeps the wind out. We propose that in the sweep-out zone the clouds are prevented from destruction essentially by a magnetic field which is amplified due to a collisional interaction with the wind. The rapid expansion of the clouds after their exit from the sweep-out zone naturally explains the disappearance of the jet hydrogen line emission and the brightening of the jets at radio wavelengths.

The collisional interaction of the jets of SS 433 with their surroundings is possible for two reasons: the powerful wind from the supercritical accretion disk and the precession with the cone opening being much larger than the jet opening angle. These factors are necessary for the existence of the cold clouds radiating the hydrogen lines. Their combination may be the clue to the uniqueness of the optical relativistic jets of SS433.

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