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Two-flow model for extragalactic radio jets

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Summary. We propose a model for extragalactic radio jets in which two different flows of particles are taken into account. One flow is a beam of relativistic particles, presumably electrons and positrons, assumed to be extracted from the funnel or the innermost part of the accretion disc and accelerated to a bulk Lorentz factor of the order of 10. It is believed to be responsible for the VLBI parsec-scale jet and for the observed superluminal motion, and carries only a negligible fraction of the total energy of the jet. The second flow consists of a classical or mildly relativistic wind, presumably made of electrons and protons coming out from all parts of the accretion disc. It is related to the large jet observed on the scale of several kiloparsecs and represents the main contribution to the energy budget of the extended radio source. The beam and the wind are both assumed pre-existing. We study their mutual interaction and their interaction with any ambient plasma or radiation field. We show that this configuration is not destroyed by the plasma-beam two-stream instability as long as the magnetic field, assumed longitudinal, is strong enough – that is to say, as long as the electron gyrofrequency ω_c is greater than the ambient plasma frequency ω_p. In weak magnetic field zones, the beam loses its energy and its momentum mainly through the development of strong Langmuir turbulence in the wind or in the ambient plasma. It then disappears quietly after some relaxation length where heating and entrainment of the ambient medium occur. This emphasizes one aspect of the important role likely played by the magnetic field in extragalactic jets. Detailed observational data on the radiosource 3C120 suggest that the magnetic field is strong ($\omega_c > \omega_p$) at least in the first 1.4 kpc of the jet. The transition from strong to weak magnetic field can occur at the location of the peculiar 4 arcsec knot described by Walker *et al.* which is then interpreted as the relaxation zone of the relativistic beam. It is natural to consider a two-flow model for extragalactic jets since formation of both beam and wind from an active galactic nucleus is expected. Beams and winds are in fact currently, but separately, studied in the literature. Such a model has two major advantages. First, it greatly reduces the problem of jet formation, since only a very small fraction of electrons and positrons has to reach the bulk Lorentz factor required to explain the superluminal motion. The main part of the jet material fuelling the radio lobes and hotspots needs to attain only the classical or mildly relativistic bulk speed expected at the large kiloparsec-scale. Second, it solves in a simple way the apparent, although still debated, discrepancy between the highly relativistic speed detected at VLBI parsec-scale and the classical or mildly relativistic speed usually expected at large kiloparsec-scale, without introducing any specific break in the continuity of the jet structure from small to large scale. Detection of any observational signature of the wind at VLBI scale, such as very faint non-superluminal radio emission associated with superluminal jets, would be an important clue in favour of such a two-flow model.

1 Introduction

The confrontation of observational data to theories and numerical simulations shows that, on the large kiloparsec-scale, extragalactic jets behave as a fluid well described, at least to first order, by the laws of magnetohydrodynamics. On the other hand, theoretical studies of black hole magnetospheres, and VLBI polarization maps of compact radio sources, both suggest a well-ordered magnetic configuration, whereas the effect of severe synchrotron losses ensures that the particles flow mainly along the magnetic field lines in the region of jet formation (for a review see Asséo & Sol 1987, and references therein). As the fluid picture is not particularly suitable for motions along the magnetic field, it appears that near the central engine the MHD approximation is not fully applicable to the jet. Whatever the mechanism of jet formation is, kinetic theory would be more appropriate. Indeed, any acceleration mechanism is expected to act on all charged particles present in the acceleration zone, namely protons, electrons and positrons when pair creation occurs, therefore leading to beams of particles with different velocities which probably produce some population inversion.

As far as we know, no theoretical attempt has yet been made to account for the transition from the formation zone of the jet to the parsec- and kiloparsec-scales. In this paper, we go just one step further in complexity for the description of the constituents of the jets on a parsec scale, hoping to obtain a more realistic view of the situation and to offer a plausible sketch for the transition zone between the VLBI structures and the large-scale extended radio jets. This will allow us to suggest qualitative answers to observational facts such as the strongly suspected discrepancies in bulk velocity, and sometimes in position angles, between the VLBI jet and the associated large kiloparsec-scale jet, without invoking drastic alteration to the standard model. Merging together the different models for jet formation, basically the radiation pressure in the funnel of a thick accretion disc and the electromagnetic extraction models, the simplest realistic configuration for the region of initial propagation of the jet includes at least two types of flows. Observation of superluminal motion in VLBI jets strongly supports the idea that, in some objects, one of the flows is highly relativistic with bulk Lorentz factors of the order of 10. We propose to describe this flow as a relativistic electron–positron beam, directly responsible for the VLBI structures. We conjecture that the other apparently unavoidable flow is a classical or mildly relativistic wind of electrons and protons coming out from the whole accretion disc. We associate this second flow with the extended radio jets detected on large kiloparsec-scale (Pelletier 1985; Pelletier, Asséo & Sol 1987; Sol, Asséo & Pelletier 1987).

The basic features of our two-flow model therefore include a relativistic beam and a much slower wind, both streaming along a longitudinal magnetic field. The presence of the beam and of the magnetic field are directly grounded on observational facts, while the existence of the wind is inferred from qualitative theoretical arguments. In this context, the fastest instability which is expected to develop is the two-stream (or beam–plasma) instability, fundamental in plasma physics. This paper discusses mainly its impact on the gross properties of the jets, taking into account the fact that the beam is relativistic and that the relative importance of the magnetic effects is likely to vary with the distance from the central engine. Let us just mention that, even in the absence of a wind from the accretion disc, this instability is still important. Indeed, it can potentially destroy the relativistic beam through its interaction with the ambient plasma pervading the initial propagation zone of the jet, namely the intercloud medium partly responsible for the confinement of the clouds of the broad emission-line region (BLR) and for the observed soft X-ray emission. A few authors have already emphasized and investigated the role of the beam-plasma instability in the context of extragalactic jets (Rose *et al.* 1984, 1987; Baker *et al.* 1988). Rose and co-workers have studied the interaction of relativistic beams with interstellar clouds, while Baker and co-workers deal with a configuration somewhat similar to ours, where a relativistic electron beam is injected into the jet plasma. Our results on the global dynamics of the relativistic beams are complementary to those works which mainly focus on the radiative properties of the beam–plasma system and on the consequences of collective emission of electromagnetic waves. In Section 2, we recall previous results that we obtained on the strong beam–plasma interaction (Pelletier, Sol & Asséo 1988) and show how they apply to the problem we are dealing with. In Section 3, we estimate other effects which could influence the beam and lead to its destruction. Section 4 describes the actual relaxation of the beam when it transfers its energy and momentum to the wind or to any ambient medium. Finally, we discuss the astrophysical issues of our two-flow model in Section 5.

2 The strong beam–plasma interaction

2.1 MICROSCOPIC VIEW OF THE INTERACTION

Classical results in plasma physics, and previous work made in the context of relativistic beams interacting with interstellar clouds (Rose *et al.* 1984, 1987), show that collisionless effects are not negligible for the global behaviour of beams of particles ejected from active galactic nuclei (AGN) and quasars. The two-stream instability expected to play a major role will be discussed first. In the absence of magnetic field, this instability excites electrostatic Langmuir waves in the ambient plasma. The presence of a magnetic field introduces the possibility of a cyclotron resonance and new branches of plasma oscillations, but the electrostatic one remains the most important (Mikhailovski 1974). As mentioned in Pelletier *et al.* (1988), this is ensured as long as the plasma frequency ω_p and the electron gyrofrequency ω_c are different. At the resonance $\omega_\mathrm{p}=\omega_\mathrm{c}$, the cyclotron oscillation branch is close to the Langmuir one but no singular behaviour is expected. Furthermore, the zone where equality between ω_p and ω_c occurs is presumably very narrow compared with typical sizes of the beam, because of the gradients in plasma density and magnetic field. Other modes possibly excited by the beam, such as high hybrid waves and whistler waves, have much lower growth rates than the Langmuir mode. As the beam is highly relativistic, quasilinear theory may not be sufficient to describe its behaviour. Indeed, quasilinear theory applies as long as the growth time of the instability is longer than the correlation time, which induces the following condition:

$$\left(\frac{n_\mathrm{b}}{n_\mathrm{p}}\right)^{1/3} \ll \frac{1}{\gamma}, \qquad (1)$$

where n_b and n_p are the beam and ambient plasma densities and γ is the bulk Lorentz factor of the beam (here we have assumed that the dispersion of the Lorentz factors in the beam is smaller than or equal to its bulk value). For high values of γ, condition (1) is not fulfilled, and the maximal growth rate of the Langmuir waves excited by the beam is:

$$\Gamma^{\max}=\frac{\sqrt{3}}{4}\frac{\omega_p}{\gamma}\left(\frac{n_b}{2n_p}\right)^{1/3}, \tag{2}$$

where $\omega_p=(4\pi n_p e^2/m)^{1/2}$ is the plasma frequency, and e and m the charge and the mass of the electron. One possible saturation effect is the trapping of resonant particles in the waves. However, the energy level reached by trapping saturation as given by Thode & Sudan (1973) satisfies the condition for generation of strong turbulence:

$$\frac{W}{n_p T}>(\Delta k_{\parallel}\lambda_D)^2, \tag{3}$$

provided that:

$$\gamma^{3/2}\left(\frac{n_b}{n_p}\right)^{1/3}>\left(\frac{v_T}{c}\right)^2. \tag{4}$$

W is the energy density in the waves, T and $v_T=(T/m)^{1/2}$ the temperature and thermal velocity of the ambient plasma, $\lambda_D=v_T/\omega_p$ the Debye length and $\Delta k_{\parallel}$ the window of the beam plasma instability, which corresponds to the spatial frequency width of the spectrum. We label the projection along and perpendicular to the magnetic field, with indices $\parallel$ and $\perp$, respectively. For classical ambient plasmas with $v_T \ll c$, this condition is easily fulfilled and a strong interaction regime occurs. In such a case the global dynamics of the beam will be controlled by the properties of the strong Langmuir turbulence and the development of the automodulational instability in the ambient plasma.

The strong Langmuir turbulence has been studied by several authors since the pioneering work of Zakharov (1972), but essentially in two limits, (i) in the absence of magnetic field or (ii) for a magnetic field so strong that the situation can be considered as unidimensional. With astrophysical applications in mind, we have investigated in a previous paper the effect of moderate magnetic fields on the evolution of the Langmuir wavepackets (Pelletier *et al.* 1988). We have found that the same partial differential equation, a non-linear Schrödinger equation modified by transverse perturbations (hereafter MNLS equation):

$$\frac{\partial^2}{\partial z^2}\left(i\frac{\partial}{\partial t}+\frac{1}{2}\frac{\partial^2}{\partial z^2}+\frac{\nu}{2}\Delta_{\perp}+|\psi|^2\right)\psi=\frac{1}{2}q\Delta_{\perp}\psi, \tag{5}$$

describes to first order the evolution of the envelope of the electric field component parallel to the magnetic field and controls the growth of the automodulational instability in the two different regimes of strong ($\omega_c>\omega_p$) and weak ($\omega_c<\omega_p$) magnetic field. The non-linear $|\psi|^2$ term is due to the ponderomotive effect, which tends to expel the plasma from the regions of a large energy density of high-frequency waves. The symbols represent reduced variables and have the same meaning as in Pelletier *et al.* (1988), namely z is the coordinate along the magnetic field, t the time, $\omega_c=eB/mc$ the cyclotron frequency; ψ is proportional to the envelope of the electric field component parallel to the magnetic field, $\Delta_{\perp}=\partial^2/\partial x^2+\partial^2/\partial y^2$ represents the transverse Laplacian, $\nu=1$ in the weak (or 0 in the strong) magnetic field case; q is proportional to $\omega_c^2/(\omega_p^2-\omega_c^2)$ and its sign corresponds to the relative strength of the magnetic field, $q>0$ for a weak field, $q<0$ for a strong field. The MNLS equation is valid only under several assumptions and has been established in the limit of small transverse wavenumber $k_\perp$ and small thermal velocity v_T, outside the narrow zone where gyromagnetic resonance $\omega=\omega_p\simeq\omega_c$ occurs. Perturbation analyses of the soliton-like solutions of the type:

$$\psi=a\,\mathrm{sech}\,a(z-z_0)\exp i[u(z-z_0)-\theta] \tag{6}$$

for the MNLS equation give, under the assumption of adiabatic evolution, the equations governing the temporal evolution of the soliton parameters, namely, the amplitude a, the position z_0, the velocity u and the phase θ. The perturbation method of the inverse scattering transform, introduced by Karpman & Maslov (1977) and Kaup & Newell (1978), allows us to take into account the additional terms due to the effects of the beam instability on a large-scale, and of the Landau dissipation on a short scale. The results are basically the following. In the strong magnetic field regime ($q<0$), the soliton-like solutions are stable. The dynamics is similar to the one-dimensional evolution governed by the Zakharov system where no collapse occurs. Thus we can expect the formation of a more or less regular lattice of close-packed solitons at rest in the ambient medium as displayed in the numerical simulation by Shen & Nicholson (1987), provided that the size $1/\Delta k_s$ of the solitons is larger than the Debye length. So as long as $\lambda_D^2\Delta k_s^2\simeq W_s/n_pT\ll 1$ (W_s being the energy density of Langmuir waves in the solitons), the Landau damping probably does not extract much power from the soliton lattice and therefore from the beam. Furthermore, the generation of large ion density fluctuations through the ponderomotive force detunes the beam–plasma interaction, switching-off the beam–plasma instability. Schamel, Lee & Moralès (1976) have shown that a maximum of about 20 per cent of the initial energy of the beam can be transferred to the electrostatic waves. The plasma is then in a state of strong Langmuir turbulence with no driving, and the beam propagates without losing a significant fraction of its energy. This result justifies the possible existence of relativistic beams on large astrophysical scales, provided that the longitudinal magnetic field is strong enough relatively to the ambient plasma density. On the other hand, in the weak magnetic field regime ($q>0$), the soliton-like solutions are unstable. They collapse towards small scales and enter into the Debye sphere where their wave energy is transmitted to the high-energy particle tail of the ambient medium through the Landau effect. No saturation mechanism of the automodulational instability occurs, a continuous transfer of energy from large to small spatial scales (i.e. from small to large wavenumbers) is established, and the beam can lose its energy and momentum at most at a rate controlled by the collapse of the Langmuir wavepackets. This induces a drastic weakening of the beam which is finally destroyed by this process after some relaxation length z_R. A self-similar description of the collapse of the Langmuir wavepackets has allowed us to estimate the average maximal power density Q transferred from the beam to the ambient medium. It is the product of the energy localized in one collapsing wavepacket (related to the initial amplitude of the soliton) by the collapse rate and by the average density of wavepackets. In the presence of an inertial range, a quasi-stationary cascade is set up when the self-modulation and the beam-growth rates are of the same order. This leads in order of magnitude to the general expression (with T expressed in erg):

$$Q\simeq\left(\frac{n_b}{n_p}\right)^{8/9}\gamma^{-8/3}\left(\frac{v_T}{c}\right)^{4/3}\omega_p n_p T \quad \mathrm{erg\,s^{-1}\,cm^{-3}}. \tag{7}$$

2.2 HYDRODYNAMICS OF THE INTERACTION

As long as the beam itself still exists, it can be described as one single fluid and its relaxation length is easily related to the value of Q. We assume the existence of a frame in which the transfer of energy and momentum occurs in a stationary regime. That is to say, the interaction front between the beam and the ambient medium is stable and moves with a speed u_f relative to the AGN which gives us an absolute reference frame. This reduces the generality of our present study since intermittency might arise, but it is expected to give a correct idea of the average parameters. In the absolute reference frame of the AGN, the temporal derivative $\partial/\partial t$ is then just $u_f(\partial/\partial z)$ and, assuming cylindrical symmetry, a constant cross-sectional area for the beam and a non-relativistic value of u_f, mass conservation and kinetic energy loss of the beam are simply:

$$\frac{\partial}{\partial z}[\gamma\rho_b(u_b-u_f)]=0 \tag{8}$$

$$\frac{\partial}{\partial z}[\gamma(\gamma-1)\rho_b(u_b-u_f)c^2]=-Q, \tag{9}$$

where ρ_b is the mass density in the beam and u_b the beam bulk velocity. We therefore define the beam mass flux $J_b=\gamma\rho_b(u_b-u_f)$ which is a constant along the beam. From equation (9) we also define a characteristic minimal relaxation length z_R, obtained when the Lorentz factor of the beam reaches a negligible value compared with its initial value γ_0 and for the initial maximal value of Q, Q_0:

$$z_R=\frac{\gamma_0 J_b c^2}{Q_0}. \tag{10}$$

Our description is justified as long as the beam is unambiguously identified and remains relativistic. When the beam is slowed down so that $u_b-u\simeq v_T$ (u being the ambient plasma velocity), the quasilinear regime takes place. Quasilinear interaction between the beam and the plasma tends to damp the beam completely and to create a plateau in the distribution function. For an homogeneous ambient plasma, this occurs on a relaxation length ($k_0=\mu_b/\omega_p$):

$$z_R(QL)=\frac{n_p mc^3}{\omega_p\varepsilon_0\langle E_\parallel^3\rangle}\frac{\Delta\gamma^2\Delta k_s}{k_0^2}, \tag{11}$$

which corresponds to the very end of the beam. Before that, Q increases smoothly with the distance z, and a typical estimate of the beam relaxation length is obtained for the values of the parameters when the beam enters the weak magnetic field zone:

$$z_R=\gamma^{14/3}\left(\frac{m_e c^2}{T}\right)^{5/3}\left(\frac{n_b}{n_p}\right)^{1/9}\frac{c}{\omega_p}. \tag{12}$$

This relaxation length is extremely dependent on the value of the Lorentz factor. Nevertheless, a maximal value of γ can be fixed of the order of $\gamma_{max}=10$ in the context of the standard model for superluminal sources. Conversely, the density ratio has no influence on z_R and is taken here to be of the order of 0.01. Constraints on the value of the ambient temperature and density are difficult to assess. Tenable values are in the range 10^4–10^8 K for the temperature and 10^{-3}–10^2 cm^{-3} for the density, and typically lead to relaxation lengths much smaller than usual sizes of VLBI jets. As an example, for $T=10^6$ K and $n_p=10^{-2}$ cm^{-3}, a set of parameters which ensures a reasonable value for the pressure of thermal material inside the jet, the relaxation length is as small as 0.05 pc for $\gamma=10$. The beam relaxation zone appears therefore more likely to be related to some radio knot or small region of the VLBI jet rather than to the overall VLBI structure.

We then suggest the following scenario. The beam goes through the strong magnetic field regions without being noticeably altered and starts losing energy as soon as it enters a weak magnetic field zone. One can define a critical value for the magnetic field $B_c = (4\pi mc^2 n_p)^{1/2} = 3.2 \times 10^{-3}\sqrt{n_p}$, above which the beam is stable (and below which it is unstable). To ensure that the beam can emerge from its formation zone, it is necessary to avoid its destruction by the development of strong Langmuir turbulence. This requires that $B/B_c > 1$ inside the beam zone close to the AGN. Indeed for the fiducial values proposed by Rees (1984):

$$B_{\rm Edd}^{\rm ambient} \simeq 4 \times 10^4 M_8^{-1/2},$$
$$n_{\rm Edd}^{\rm ambient} \simeq 10^{11} M_8^{-1}, \tag{13}$$

we obtain a ratio $B/B_c = 40$, independent of the mass $10^8\,M_8 M_\odot$ of the central black hole. But those values do not apply to the jet itself, however. For a magnetic field purely longitudinal, a simple model of magnetic flux conservation gives a magnetic field inversely proportional to S, the cross-sectional area of the jet. Mass flux conservation of the particles in the wind leads to the same variation for the density n_p. Therefore B/B_c decreases as $1/\sqrt{S}$ with the distance from the AGN if the jet widens, as expected on average. We then advance as a first-order model that the beam is destroyed in a small relaxation zone at the distance z_c from the nucleus such that $B = B_c$, that is to say $B/B_c \simeq 1$ is reached at the end of the VLBI structure due to the primary beam of particles with relativistic bulk velocity.

The detailed observational results of Walker, Benson & Unwin (1987) on 3C120 allow a check on the overall coherency and relevance of our proposition. From combined VLBI and VLA data, Walker *et al.* can derive the evolution of several physical parameters along the jet of 3C120. In particular, from a fit of the emissivity versus the transverse width of the jet, they obtain, under the usual minimum-energy assumption, the variation of the magnetic field:

$$B = 4.1 \times 10^{-5} r^{-0.97} \quad \text{G}, \tag{14}$$

where r is the measured full width at half maximum (FWHM) of the jet in arcseconds. The absence of any significant Faraday rotation gives an upper limit on the ambient plasma density inside the emitting jet:

$$n_p \lesssim 9.9 \times 10^{-4} r^{-0.03} \quad \text{cm}^{-3}, \tag{15}$$

assuming a uniform magnetic field mainly orientated along the line-of-sight (which is justified since (i) polarization data show that the magnetic field is longitudinal and (ii) the jet makes a small angle with the line-of-sight as superluminal motions are detected). These two observed laws imply that the magnetic field is strong in the first part of the jet close to the AGN and reaches its critical value for $r = r_c \gtrsim 0.4$ arcsec $\equiv 180$ pc (1 arcsec $\equiv 460$ pc), which corresponds to $z = z_c \gtrsim 3$ arcsec $\equiv 1.4$ kpc. This distance is amazingly close to the location of the 4-arcsec knot, a 'rather curious structure' described by Walker *et al.* at about 4 arcsec (1.6 kpc) from the AGN. At such a distance, the relaxation length of the beam is in the range of 5–500 pc for γ in the range of 4 (the minimum required in the standard superluminal scenario) to 10. Those values are compatible with typical longitudinal sizes of the 4-arcsec knot. From the expression for the density of radiating particles in the beam given by Walker *et al.*, $n_r = 6.7 \times 10^{-8} r^{-2.43}$ in the observer's frame, we deduce the proper beam density $n_b \simeq 0.7 \times 10^{-6}/\gamma f$ at the location of the 4-arcsec knot (f being the fraction of radiating particles in the beam). The power of the beam can be estimated as $P = Qz_R S = \gamma J_b S c^2 \simeq 2\gamma^2 \times 10^{40}$ erg s^{-1}. The radio luminosity radiated in the 4-arcsec knot (of the order of 10^{39} erg s^{-1}, without including any Doppler enhancement effect) is therefore just a small fraction of the total power of the beam. All this strongly suggests that the knot corresponds to the expected relaxation zone of the beam, as dissipation of the beam energy is likely to enhance the radio emission in this zone. Several properties of this curious 4-arcsec knot find a natural explanation in this framework and will be discussed in more details in Section 5. Let us just mention that radio maps of other superluminal sources such as NRAO 140, 4C39.25, 3C279, 3C345, 3C454.3, 3C245, 1642 + 690 and 1928 + 738 show the presence of somewhat similar knots at a few arcsec from the central radio component (Browne *et al.* 1982; Browne 1987; Simon *et al.* 1987; Hough & Readhead 1987).

3 Other beam dissipation mechanisms

Several effects other than the strong beam–plasma interaction can potentially affect and even destroy the relativistic beams after a characteristic relaxation length. These effects must be taken into account, especially in the strong magnetic field zone where the strong beam–plasma instability saturates. As relativistic beams are actually observed on the scale of several parsecs, their total destruction by any effect must be avoided, that is to say, the associated relaxation length must remain always larger than the effective length of the beams, or other compensating effects such as acceleration or confinement mechanisms must be efficient. This induces unavoidable constraints on the physical parameters concerned. In the following we recall and comment on the relaxation length of different kinds of possible beam dissipation. Just as for the case of cosmic rays, radiative losses are expected to have more influence than pitch angle scattering for energetic particles (Ginzburg 1979). Our list includes the most obvious effects but is not exhaustive. We assume the other constraints fulfilled *a priori* to ensure that the beams can escape from their formation zone to the VLBI scale, and refer the reader to Rose *et al.* (1984, 1987) for interaction with interstellar clouds.

It seems that the only way to overcome the Compton drag problem is to imagine some acceleration mechanism taking place as long as the radiation field is too strong close to the AGN (Phinney 1987). This is a way to solve one fundamental problem of the jet formation models. Here we just estimate the inverse Compton losses at the VLBI scale, once the relativistic beams have already been produced. The relaxation length of the beam due to inverse Compton losses is given by:

$$z_{\rm R}({\rm IC}) = \frac{3}{4}\frac{mc^2}{\gamma\sigma_{\rm T}w_{\rm ph}} \simeq \frac{0.3}{\gamma w_{\rm ph}} \quad {\rm pc}, \tag{16}$$

where $\sigma_{\rm T}$ is the Thomson cross-section and $w_{\rm ph}$ the radiation field energy density in erg cm^{-3}. At the VLBI scale, the radiation field is essentially dominated by the emission of the AGN and can be expressed as:

$$w_{\rm ph} = L/4\pi cz^2 \simeq 2.9\times10^{-3} z_{\rm pc}^{-2}, \tag{17}$$

where L, the luminosity of the AGN, has been taken as $\sim 10^{46}$ erg s^{-1}, close to the Eddington value. The relaxation length is then $z_{\rm R}({\rm IC}) \sim 100 z_{\rm pc}^2/\gamma$, far larger than typical VLBI jet size. Thus, although implying drastic constraints to the jet formation models in the acceleration zone in the neighbourhood of the AGN (typically for $z_{\rm pc} < \gamma/100$), inverse Compton losses appear negligible above the parsec scale.

Radiative synchrotron losses become effective in slowing down the bulk velocity of the beam after a relaxation length:

$$z_{\rm R}({\rm S}) = \frac{9}{4}\frac{m^3c^6}{e^4\gamma B_\perp^2} \simeq 10/\gamma B_\perp^2 \ {\rm pc}, \tag{18}$$

where $B_\perp$ is the magnetic field component transverse to the beam. Actual propagation of the beam therefore requires $\gamma B_\perp^2 \ll 10/z_{pc}$, which is to be expected, especially in a model where the magnetic field is longitudinal. On the other hand, a fraction f of the beam particles acquire high individual Lorentz factors, $\gamma_r = 9 \times 10^{-4} \nu^{1/2} B_\perp^{-1/2}$, and are responsible for the synchrotron emission observed. With the parameters of 3C120, particles with $\gamma \sim 10^4$ radiate at 15 GHz in a field $B = 10^{-4}$ G and lose their energy in 10^5 pc, so there is no need for reacceleration of particles along the beam.

Relativistic bremsstrahlung in the ambient plasma is another kind of energy loss, with a relaxation length (without screening effect):

$$z_R(\mathrm{RB}) = mc^3/7 \times 10^{-23}\, n_p(\log \gamma + 0.36) \simeq \frac{3 \times 10^8}{n_p} \quad \mathrm{pc} \tag{19}$$

(Ginzburg 1979). This effect requires $n_p \ll 3 \times 10^8\, z_{pc}^{-1}$ to be negligible, which is easily fulfilled for tenable values of the parameters or in the case of 3C120, for instance.

Ionization losses can occur if the beam crosses neutral hydrogen clouds. The associated relaxation length is then:

$$z_R(\mathrm{IL}) = \frac{\gamma mc^3}{3.5 \times 10^{-31}\, n_H(3 \log \gamma + 18.8)} \simeq 10^{15}\, \gamma/n_H \quad \mathrm{pc}, \tag{20}$$

where n_H is the neutral hydrogen density. Ionization losses are therefore not expected to play a significant role at the VLBI scale except in case of interaction with a cloud of the BLR (Rose *et al.* 1984, 1987).

Diffusion on turbulent Alfvén waves might also destroy the relativistic beams. The diffusion occurs only for particles above an energy cut-off of $m_p v_A c$, where m_p is the proton mass and v_A the Alfvén speed, so there is no diffusive effect as long as $\gamma \lesssim 10^3$.

For an electron–positron beam, the electron–positron annihilation mechanism must be mentioned. The cross-section for annihilation with a free electron at rest is:

$$\sigma \simeq \frac{\pi r_0^2}{\gamma} [\ln(2\gamma) - 1] \tag{21}$$

for $\gamma \gg 1$ (γmc^2 being the energy of the positron and $r_0 = e^2/mc^2$ the classical electron radius; Lang 1980), and the $e^- e^+$ relaxation length is given in the observer's frame by:

$$z_R(e^- e^+) = 1/\gamma n_b \sigma \simeq \frac{10^6}{n_b} \quad \mathrm{pc}, \tag{22}$$

where n_b is the proper beam density. The $e^- e^+$ relaxation therefore occurs on scales typically larger than a megaparsec and is insignificant.

We just mention the possibility of energy loss by the Joule effect if the beam is electrically charged and if there is a return current. Radiation by other modes such as high hybrid waves at the synchrotron resonance and whistler waves at the Cherenkov resonance should also be considered, although their growth rates are slow. Moreover, high hybrid modes can be easily stabilized by the Landau effect, and whistler modes are not expected to propagate since their wavelengths are larger than the mean free path between solitons of the lattice. In any case, none of the above effects appears competitive with the strong beam–plasma interaction to destroy the relativistic beams at the VLBI scale between a parsec and a kiloparsec. The only noticeable exception is the inverse Compton effect which occurs in the formation zone of the beams at a scale below the VLBI one, and which must be overcome by some appropriate acceleration mechanism. We are therefore confident that other dissipation mechanisms of the relativistic beams do not drastically alter the first-order model proposed in Section 3.

4 Description of the beam relaxation

The transfer of energy and momentum from the beam to the wind (i.e. to the ambient plasma) which occurs when the beam reaches a weak magnetic field zone results in acceleration and heating of the wind. It can be described by the equations of hydrodynamics for two coupled fluids. Indeed, only the suprathermal particles of the ambient plasma can receive energy through wave–particle interaction by the Landau effect when the collapsing Langmuir wavepackets enter the Debye sphere. Acoustic waves, however, are emitted as a by-product of the collapse mechanism when it reaches the supersonic regime, and an ionic turbulence is generated. Diffusion on this turbulence of all the particles of the ambient plasma, the thermal as well as the suprathermal ones, isotropizes the distribution function of the plasma electrons due to pitch angle scattering. Indeed, it is noteworthy that the beam tends to generate a current when driving strong Langmuir turbulence, since it loses energy and momentum in accelerating first ambient electrons, inducing a drift velocity between electrons and protons of the ambient plasma. As strong Langmuir turbulence produces intense density fluctuations, an anomalous resistivity limits the current. The effective collision frequency ν_* gives the rate of isotropization of the electron distribution in the frame of the ions (Papadopoulos & Coffey 1974; Roubaud 1984):

$$\nu_* = \sqrt{\frac{\pi}{2}}\,\omega_{\rm p}\int \frac{w_{\rm f}(k)}{n_0 T}\,\frac{k_{\rm D}}{k}\,dk \simeq \omega_{\rm p}\,\frac{W_{\rm f}}{n_0 T} \tag{23}$$

where $W_{\rm f} = \int w_{\rm f}(k)\,dk$ is the energy density in the density fluctuations. The corresponding relaxation length is very short since one can reasonably take $\nu_* \simeq 10^{-2}$ or $10^{-3}\,\omega_{\rm p}$. So isotropization of the ambient plasma by pitch angle scattering occurs on a characteristic length:

$$z(\mathrm{I}) \simeq 100c/\omega_{\rm p}, \tag{24}$$

much smaller than $z_{\rm R}$, the effective relaxation length of the beam obtained from strong beam–plasma interaction. The ambient plasma can then be considered as one-single fluid entrained by the beam over all the beam relaxation zone. If the wind remains non-relativistic, mass flux conservation in the beam and in the wind, kinetic energy loss of the beam, and total energy and momentum conservation expressed in the frame of the AGN, under the same assumptions as in Section 2, provide the following closed system of equations:

$$J_{\rm b} = \gamma\rho_{\rm b}(u_{\rm b} - u_{\rm f}) = \text{constant} \tag{25}$$

$$J = \rho(u - u_{\rm f}) = \text{constant} \tag{26}$$

$$J_{\rm b}c^2\,\frac{d\gamma}{dz} = -Q \tag{27}$$

$$A = \frac{1}{2}Ju^2 + \frac{5}{2}pu - \frac{3}{2}pu_{\rm f} + \gamma c^2 J_{\rm b} = \text{constant} \tag{28}$$

$$F = Ju + p + \gamma u_{\rm b} J_{\rm b} = \text{constant}, \tag{29}$$

where ρ, u and $p = n_{\rm p}T$ are the wind density, velocity and classical pressure, with adiabatic index of 5/3. Combining (28) and (29) we obtain the equation for the plasma velocity:

$$2Ju^2-\frac{5}{2}(F-\gamma u_b J_b+\frac{3}{5}u_f J)u+\frac{3}{2}u_f(F-\gamma u_b J_b)+A-\gamma c^2 J_b=0. \quad (30)$$

Asymptotic values (denoted by the index ∞) are reached at the end of the relaxation zone where $u_{b_\infty}=u_\infty$. The asymptotic velocity u_∞ can be deduced from (30). Our description, which assumes non-relativistic values for u and u_f, provides regular solutions without shocks as long as $(J+J_b)/J_b \leqslant 25\gamma_0^2/32(\gamma_0-1)$, and leads to values of u_∞ higher than $2c/5$. Entrainment of the ambient wind by the beam can therefore be quite efficient. Such mildly relativistic values for the asymptotic wind speed, if one identifies them with the large-scale jet velocity, are high enough to explain the one-sidedness of large-scale jets by a Doppler beaming effect.

The derivative of the equation (30):

$$\frac{1}{2}[3J(u-u_f)-5p]\frac{du}{dz}+Q\left(1-\frac{5}{2}\frac{u}{u_b}+\frac{3}{2}\frac{u_f}{u_b}\right)=0 \quad (31)$$

gives the relation at any critical point (denoted with index $*$),

$$u_*-u_f=\frac{5p_*}{3J}=c_s-u_f, \quad (32)$$

where c_s is the sound velocity, and the condition for the existence of a transonic solution:

$$u_*-u_f=\frac{2}{5}(u_{b_*}-u_f). \quad (33)$$

So the following equation must be fulfilled at the sonic point:

$$u_{b_*}\left(1+\frac{25}{16}\gamma_*\frac{J_b}{J}\right)=\frac{25}{16}\frac{F}{J}-\frac{9}{16}u_f. \quad (34)$$

Let us introduce the new variables

$$Y=\frac{8J(u-u_f)}{5(F-Ju_f-\gamma J_b u_b)} \quad (35)$$

and

$$Z=\frac{32J[A-\gamma J_b c^2+Ju_f^2/2-(F-\gamma J_b u_b)u_f]}{25(F-\gamma J_b u_b-Ju_f)^2}. \quad (36)$$

Then equation (30) is equivalent to

$$Y^2-2Y+Z=0 \quad (37)$$

and the conditions (32)–(34) for a transonic solution can be written as

$$Z_*=Y_*=1 \text{ and } \frac{\partial Z_*}{\partial u_b}=0.$$

The flow is subsonic for $Y<1$ and supersonic for $Y>1$. The subsonic branch is $Y_-=1-\sqrt{1-Z}$ and the supersonic one $Y_+=1+\sqrt{1-Z}$. Clearly these solutions exist only if $Z<1$. Fig. 1 illustrates the flow diagram. The flow is transonic if it starts with an initial value Y_1 and Y_2 such that $Y_2=2-Y_1$. In the former case, it starts subsonic and becomes supersonic at the critical

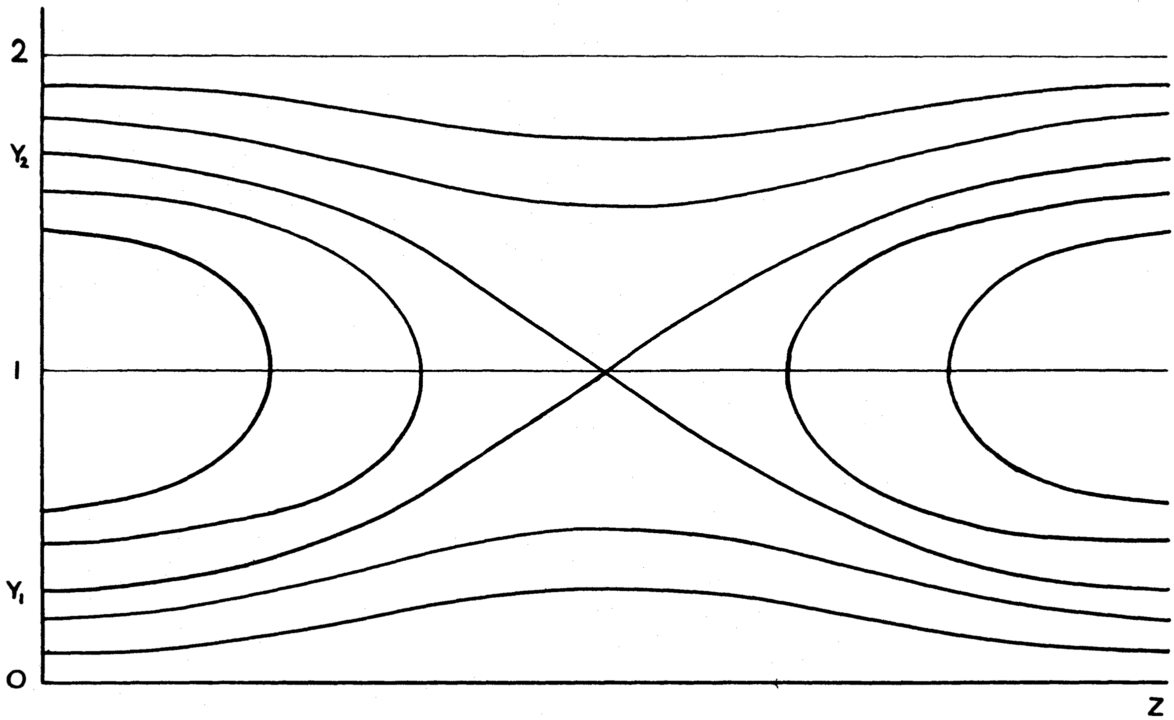

Figure 1. Illustration of the flow diagram. Variation of $Y=8J(u-u_f)/5(F-Ju_f-\gamma J_b u_b)$ versus the distance z to the central engine (or any increasing function of z, such as $\gamma_0-\gamma$ or $u_{b_0}-u_b$), for the different types of flows.

point, and conversely in the latter case. Regular flows exist for initial values of Y_0 such that $Y_2<Y_0<2$ and are entirely supersonic, or for Y_0 such that $0<Y_0<Y_1$ and are entirely subsonic. For Y_0 between Y_1 and Y_2, no regular stationary flow exists and shocks develop. So it is important to know the initial conditions which lead to a transonic flow. For that purpose it is convenient to write Z in the form:

$$Z=\frac{32}{25}\frac{J}{J_b}\frac{(\gamma_0-\gamma)c^2-u_f(\gamma_0 u_{b_0}-\gamma u_b)+u_1^2}{(\gamma_0 u_{b_0}-\gamma u_b+u_2)^2}, \tag{38}$$

where

$$u_1^2=\frac{u_0-u_f}{2J_b}[J(u_0-u_f)+5p_0] \tag{39}$$

$$u_2=\frac{J}{J_b}(u_0-u_f)+\frac{p_0}{J_b}. \tag{40}$$

The terms in u_f and in u_1^2 in (38) are essential for the initial conditions but become rapidly negligible, since u_0 and u_f are assumed non-relativistic. Thus:

$$Z_0=\frac{32}{25}\frac{J}{J_b}\left(\frac{u_1}{u_2}\right)^2<1 \tag{41}$$

and

$$Z_\infty\simeq\frac{32}{25}\frac{J}{\gamma_0 J_b}<1 \tag{42}$$

are important constraints on the parameters if one wants to obtain regular flows avoiding shock formation. Straightforward algebra shows that, if $\gamma_0 \gg 1$, the critical point, when it exists, is still relativistic. The transonic conditions $\partial Z_*/\partial u_b = 0$ and $Z_* = 1$ lead to

$$\gamma_0 - \gamma_* \simeq \frac{u_2}{c} - \frac{2u_1^2}{c(c-u_f)} \tag{43}$$

and

$$\left(\gamma_0 - \gamma_* + \frac{u_2}{c}\right)^2 \simeq \frac{32}{25}\frac{J}{J_b}\left(\gamma_0 - \gamma_* + \frac{u_1^2}{c^2}\right). \tag{44}$$

It turns out that $u_1^2/c^2 \ll u_2/c$, which can be checked *a posteriori*. So we find $\gamma_0 - \gamma_* \simeq u_2/c$ with $u_2 = (8J/25J_b)\,c$, or from relation (40), we obtain the condition for a transonic flow:

$$J(u_0 - u_f) + p_0 = \frac{8}{25} Jc, \tag{45}$$

allowing two possibilities for the interaction front speed:

$$u_0 - u_f = \frac{8}{25} c \tag{46}$$

or

$$u_0 - u_f = \frac{25}{8}\frac{p_0}{\rho_0 c} = \frac{15}{8}\frac{c_s^2}{c}, \tag{47}$$

where c_s is the sound speed, $c_s^2 = 5p_0/3\rho_0$. The interaction front is then slower than the initial value of the wind velocity relative to the central AGN. It slightly proceeds up the ambient flow which carries it. Now, from relation (35):

$$Y_0 = \frac{8J(u_0 - u_f)}{5[J(u_0 - u_f) + p_0]}. \tag{48}$$

The values of Y_0 leading to a transonic solution are thus $Y = 5(u_0 - u_f)/c$. Hence, when (45) is fulfilled with $u_0 - u_f = 8c/25 > c/5$, $Y_0 = 8/5$, the flow is transonic, starting supersonic and becoming subsonic. This corresponds to a case where the sound speed increases more rapidly than the plasma velocity; heating of the plasma by the beam is more efficient than entrainment. When (45) is fulfilled with $u_0 - u_f = 15c_s^2/8c < c/5$, $Y_0 = 75c_s^2/8c^2$, we obtain a transonic flow starting subsonic and becoming supersonic. The flow is entirely subsonic if $u_0 - u_f < c/5$ and $p_0 > 8Jc/25 - J(u_0 - u_f) > 3Jc/25$ and entirely supersonic when the reverse inequalities occur. We note that $p_0 > 0$ implies $u_0 - u_f < 8c/25$ for the supersonic solutions. In the other cases, a shock takes place. The asymptotic velocity u_∞ is such that:

$$J(u_\infty - u_f) \simeq \frac{5}{8}\gamma_0 J_b u_{b_0}(1 \pm \sqrt{1 - Z_\infty}) \tag{49}$$

with Z_∞ given by (42). The asymptotic pressure p_∞ is thus

$$p_\infty = [p_0 + \gamma_0 J_b c - u_\infty J_b + J(u_0 - u_f)][1 - \frac{5}{8}(1 \pm \sqrt{1 - Z_\infty})]. \tag{50}$$

So another important constraint comes from the requirement that $p_\infty > p_0$, which requires $16/25 < Z_\infty < 1$ for a supersonic velocity. The asymptotic velocity reaches high values, since

$$\frac{u_\infty}{c} \simeq \frac{4}{5}\frac{1 \pm \sqrt{1 - Z_\infty}}{Z_\infty}, \tag{51}$$

and the asymptotic behaviour of supersonic solutions would even require a relativistic treatment since $u_\infty \gtrsim 4c/5$. With the set of parameters adopted for 3C120, one can show that the constraints on Z_∞ which allow a transonic solution starting subsonic and becoming supersonic are fulfilled if one assumes $T \simeq 6\times10^7$ K for $\gamma_0=4$ or $T \simeq 9\times10^8$ K for $\gamma_0=10$. The possibility of such transonic solutions shows the ability of relativistic beams to transform a subsonic flow of ambient plasma into a supersonic wind or jet. Furthermore, the relaxation of the beam does not necessarily induce a decollimation process. On the contrary, the efficient entrainment of the electrons of the ambient wind induces a drift velocity between the ambient electrons and protons. This leads to the generation of an electric current limited by the development of an anomalous resistivity. This current creates a subsequent azimuthal magnetic field which possibly reinforces a pre-existing confining magnetic field. Magnetic field production due to drift velocities has been recently considered by Rose (1987), but in the opposite case where electrons are slowed with respect to the protons.

To analyse this phenomenon, we assume in the following that, when there is no beam relaxation, the wind can be described by ideal axisymmetric MHD. The magnetic and flow surfaces are the same since $\mathbf{u}_p$ and $\mathbf{B}_p$, the poloidal components of the wind velocity and of the magnetic field, are parallel. Then S, the cross-sectional area of the wind, corresponds also to the section of a magnetic tube. Conservation of the magnetic flux in the wind implies $B_z S=$ constant and, for small variation of S and neglecting rotation effects, the electric current j close to the axis is proportional to B_z. Therefore $jS=$ constant. The current generated by the beam relaxation, j_R, modifies Ohm's law and the ideal MHD current j_0, with a subsequent change of the section of the wind such that

$$S_R = \frac{S_0}{1 + j_R/j_0}. \tag{52}$$

The relaxation zone corresponds to a recollimation region if j_R and j_0 have the same sign. Due to its origin, the generated current j_R is always pointing towards the central source, so the relaxation of the beam reinforces the magnetic collimation effect provided that the MHD magnetic axis and the rotation axis of the accretion disc and of the wind have the same direction. If the two axes are antiparallel, the MHD confinement is less efficient in the relaxation zone, without being necessarily entirely destroyed. Estimates of $j_0 \simeq cB/4\pi R$ and of $j_R \simeq Q/n_p e\eta c$, where R is the jet radius and $\eta = n_p \nu$ the anomalous resistivity, with $\nu = x\omega_p$ the anomalous collision frequency, lead to the ratio:

$$\frac{j_R}{j_0} \simeq \frac{R}{z_R}\frac{\omega_p}{\omega_c}\frac{\gamma_0 n_b}{x n_p}, \tag{53}$$

which can be of the order of R/z_R for typical values of x in the range 10^{-2}–10^{-3}. The relative importance of the additional current j_R is then highly dependent on the Lorentz factor γ_0. In the case of the 4-arcsec knot of 3C120, j_R/j_0 is in the range 0.03–3 for γ_0 between 4 and 10, which illustrates the possibility of a non-negligible recollimation effect due to j_R. A further comment is that, in the context of MHD models for jets and when the magnetic and rotation axes are parallel, the relaxation phenomenon helps to stabilize the beam confinement in the strong magnetic field zone. Indeed, any sudden increase of the cross-sectional area of the beam induces a decrease of the ratio B/B_c because of magnetic flux conservation. Relaxation starts if B/B_c reaches unity, which, as we just saw, intensifies the magnetic confinement and reduces the section. This can be of interest since on one hand magnetic confinement is known to be easily unstable, and on the other hand the external pressures deduced from X-ray observations of

AGN are often smaller than the internal minimum pressures deduced from equipartition inside the VLBI jets, therefore prohibiting the confinement of VLBI jets by external thermal pressure.

When a shock develops instead of a regular flow, one deduces the Rankine–Hugoniot-like relations from equations (25), (26), (28) and (29). Specifying with index 1 the upstream (and 2 the downstream) physical quantities, one can easily show that, in the case where the initial beam can be considered as a perturbative flow superimposed on the plasma flow, the compression ratio r, defined by $u_1 = ru_2$, is simply:

$$r \simeq \frac{4}{1+5p_1/J_1 u_1} - \frac{32}{3}\frac{\gamma_1 J_{b_1}}{J_1}\frac{c^2}{u_1^2}. \tag{54}$$

Thus the additional beam flow J_{b_1} leads to a decrease of the compression ratio. The pressure downstream of the shock, however, is usually greater than the upstream pressure.

In cases both of shock or regular flow, increase of the pressure corresponds to plasma heating, especially heating of the electrons of the suprathermal tail. This induces subsequent radiation by the energetic suprathermal electrons and qualitatively explains the enhancement of radio emission expected at the relaxation zone of the beam. Assuming a distribution function in $(\xi-1)\, n_{\rm tail}\gamma^{-\xi}$ for the suprathermal tail ($n_{\rm tail}$ being the density of particles in the tail, typically of the order of 10^{-3}–$10^{-1}\, n_{\rm p}$), one can express the pressure as the integral of the internal energy:

$$p_\infty = \frac{1}{3} m_{\rm e} c^2 \int_1^{\gamma_{\max}} (\xi-1)\, n_{\rm tail}\gamma^{1-\xi} d\gamma, \tag{55}$$

where $\gamma_{\max}$ is the high cut-off for individual particle Lorentz factor. For $p_\infty \simeq \gamma_0^2 n_{\rm b} m_{\rm e} c^2$ one obtains:

$$\gamma_{\max}^{2-\xi} - 1 = \frac{3(2-\eta)\gamma_0^2 n_{\rm b}}{(\eta-1)\, n_{\rm tail}}. \tag{56}$$

For ξ slightly smaller than 2, $\gamma_{\max}^{2-\xi} \simeq \gamma_0^2 n_{\rm b}/n_{\rm tail}$, and the Lorentz factors can reach values high enough to cover the whole range of observed radio frequencies through synchrotron emission. With the parameters of 3C120, one can check that the corresponding expected luminosity can be high enough to account for the radio emission enhancement at the location of the 4-arcsec knot.

In the model proposed in Section 2, relaxation of the beam occurs at the end of the VLBI structure. The beam is then gently dissipated without inducing strong discontinuities or disruption of the jet pattern. Heating and acceleration up to mildly relativistic speeds of the wind or of any present ambient plasma take place, locally strengthening the radio emission and supplying high-velocity material to the large-scale jet. The reachable speeds are high enough to allow some Doppler enhancement of the large-scale jets. They are, however, smaller than the highly relativistic speeds required to explain the superluminal motion and are thus more compatible with radio data and the usual description of large-scale jets.

This model may be modified by the occurrence of density and magnetic field fluctuations, for instance when a cloud of the broad or narrow line regions crosses the beam. Increase of the ambient density gives rise to beam relaxation. It can be only partial if the size of the cloud is smaller than the relaxation length $z_{\rm R}$. The cloud is then probably swept by entrainment. This is one way to explain the huge discrepancy between the upper limits on $n_{\rm int}$, the density internal to the jet deduced from limits on Faraday rotation and depolarization measures (Wardle *et al.*

1986; Walker *et al.* 1987), and n_{ext}, the expected external density as deduced from arguments on the intercloud medium in the broad- and narrow-line regions (Collin-Souffrin & Lasota 1988). It is worth mentioning that, contrary to usual shocks, this type of interaction between the beam and a cloud can correspond to a recollimation zone and does not destroy the beam confinement.

5 Astrophysical implications of the two-flow model

The model proposed in this paper emphasizes the likely existence of two populations of particles and the important role played by an ambient plasma or a slow wind of particles issuing from the accretion disc for the description of extragalactic radio jets. It is not drastically different from the usual picture and still has to face the same basic problems, although some of the difficulties seem to be easier to solve in the context of the two-flow model. The question of jet formation might be more readily overcome, since it is now only a small fraction of the particles which need to reach a highly relativistic bulk velocity. With the parameters of 3C120 taken at the 4-arcsec knot location, the mass carried by the wind is slightly less than, or of the order of, $(u/c) M_\odot$ yr^{-1}, but it is only of $4\times10^{-7} M_\odot$ yr^{-1} for the beam associated with the VLBI jet. Another puzzle of extragalactic jets finds here a natural explanation: highly relativistic motions are deduced at VLBI scale from observation of superluminal sources, while there is evidence for slower classical or mildly relativistic bulk velocity at the scale of VLA maps. This situation is typically the case for 3C120, where superluminal motion is detected at the VLBI scale and where the absence of severe changes in surface brightness at bending of the large-scale jet seems to rule out beaming effects and highly relativistic speeds above the kiloparsec-scale. The change in velocity must nevertheless occur without disrupting the jet or producing strong discontinuities in its shape. This is precisely what happens in the relaxation zone studied in Section 4, when the beam reaches the weak magnetic field zone and transfers its energy and momentum to the ambient plasma. The population of particles with highly relativistic bulk velocity disappears smoothly, leaving only the classical or mildly relativistic wind to make up the large-scale jet. Furthermore, instead of disturbing the jet structure, this phenomenon can help to stabilize the confinement mechanism and to reinforce the collimation. A last observational point can be qualitatively accounted for in the frame of the two-flow model; it is the discrepancy often observed between the position angle of the jets at small and at large scale (Simon *et al.* 1987). The beam and the jet can, indeed, be issued from different parts of the central engine, namely the funnel of a thick accretion disc for the beam and the outer edges of the disc for the wind. They are therefore not necessarily emitted in exactly the same direction, since the inner and outer parts of the accretion disc can have different rotation axes, due for instance to the Lense–Thirring effect.

We argue that the peculiar 4-arcsec knot of 3C120 is a possible example of a zone of beam relaxation through collapses of Langmuir wavepackets. We saw in Section 2, using the expression of Walker *et al.* (1987) for the minimum-energy magnetic field, and assuming for the density of the ambient plasma a value close to their upper limit deduced from the absence of detectable Faraday rotation effects, that this 4-arcsec knot is located at the place where the plasma frequency ω_p becomes comparable to the cyclotron frequency ω_c, that is to say where the collapses are expected to occur with subsequent ambient plasma heating and radio emission enhancement. The polarization direction remains perpendicular to the jet over all the knot, which is consistent with our relaxation picture but in contrast with what would be expected if it were due to an usual shock, or to projection and beaming effects when the curving jet passes through the line-of-sight. Significant bending of the jet occurs just beyond the 4-arcsec knot, also suggestive of a change in the flow regime. Moreover, Walker *et al.* (1987)

point out that the half-width of the jet narrows at the knot, which can be interpreted as the recollimation effect described in our relaxation model. Our proposal predicts that very relativistic bulk motion still exists at 1.6 kpc from the nucleus of 3C120, with a possible superluminal effect depending on the beam orientation. This is consistent with recent VLBI results by Benson *et al.* (1988), who possibly detected superluminal motion up to 0.1 arcsec from the core and showed that the VLBI jet does not slow down quickly. It does not, however, necessarily require highly relativistic bulk velocity beyond the 4-arcsec knot.

Various observational tests of the two-flow model can be envisaged. Any direct evidence of the wind, for instance from significant Faraday rotation and depolarization measures, or from detection of very faint non-superluminal radio emission at VLBI scale, would strengthen the likelihood of such types of model. Non-superluminal components have been already detected in superluminal sources such as 4C39.25, 0836 + 71, 3C395, and possibly 0153 + 74 (Waak *et al.* 1985; Schalinski *et al.* 1988; Shaffer *et al.* 1987; Simon *et al.* 1987; Witzel *et al.* 1988), although they cannot be easily ascribed to any slow wind of particles, since they are as bright as the superluminal components and are therefore difficult to interpret in the absence of significant Doppler boosting in the standard model of superluminal sources (Blandford & Königl 1979). Data on the variation along the jets of the magnetic field and of the density of the ambient medium would be of particular interest to check whether features similar to the 4-arcsec knot of 3C120 are present in other sources such as 4C39.25 or 3C454.3 at the location where equality between plasma frequency and cyclotron frequency occurs. Conversely, our two-flow description should allow constraints to be placed on the physical parameters of the jets, such as the Lorentz factor of the beam, which is strongly related to the length of the relaxation zone, but any definite conclusion would still be premature, pending further investigation of the widespread validity of the model.

Two different morphological types of the relaxation zone and of the jet beyond that zone can be expected according as the rotation axis of the central engine and the large-scale magnetic axis are parallel or antiparallel, since the relaxation favours deconfinement in the latter case, while enhancement of the confinement is expected in the former one. If the direction of the magnetic field is reversed on both sides of the accretion disc, as in the model of odd field symmetry studied by Lovelace, Wang & Sulkanen (1987), this effect should be considered as one source of asymmetry of the jet structure. Matter would be ejected on both sides of the AGN, keeping the overall symmetry of the problem of jet formation and of the large-scale structure of radio sources, but the VLBI jets and the relaxation zones would be detected essentially on the side where reconfinement is efficient, with subsequent amplification of the radio emission and collimation of the jets. The superluminal source 4C39.25, with a jet and a compact hotspot on one side and a more relaxed counterjet on the other, could be an illustration of this phenomenon (Browne *et al.* 1982). The new latitude brought by the two-flow model allows qualitative explanations of the different types of extragalactic radio sources. For instance, compact radio sources with no well-organized large-scale jet can be interpreted as cases where transonic or supersonic flows do not occur, and where the relative importance of the wind is small. Conversely, two main families of jets have been identified, one associated with low-luminosity sources and the other with high-luminosity sources. Let us just suggest one possible origin of this difference, the presence or the absence of highly relativistic beams coming from the AGN. Pair creation and efficient acceleration mechanism could be at work only in high-luminosity sources.

Focusing on the role of the magnetic effects and of the beam–plasma interaction, which seem unavoidable in the surroundings of AGN and quasars, we have presented here one possible picture of the first kiloparsecs of extragalactic jets. Our main result is to show that, due to a strong magnetic field, highly relativistic beams of particles are not destroyed by the beam–plasma interaction in the strong turbulence regime. This warrants their existence over several hundred parsecs from the AGN and strengthens the idea that coherent radio emission should be considered for VLBI jets (Baker *et al.* 1988). Coherent and incoherent Compton radiation from relativistic electrons, scattered by Langmuir waves at radio frequencies of order $\gamma^2\omega_p$, could be emitted by the VLBI jet before the end of the relaxation zone, this emission being superposed on the synchrotron emission. Fluctuations of the Langmuir waves should then be taken into account as a new source of variability. Estimates of the physical parameters of the jets are model-dependent, however. This allows other descriptions of the jets, implying the same type of beam–plasma interaction but at another level. An interesting alternative to the present model is to assume that beams with bulk Lorentz factors of the order of 10^3 are generated far enough from the central engine and from any strong radiation field to avoid Compton drag. The observed superluminal motion is then ascribed to the motion of some interaction front. These beams can provide enough energy to power the large-scale radio structure by entrainment of the ambient interstellar medium, even if no additional wind is issuing from the accretion disc. Another alternative model, beyond the scope of the present work, invokes the possibility of intermittent beam relaxation, which could explain the knotty structure of the VLBI jets.

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