J. Astrophys. Astr. (1986) **7,** 225–236

The Jet of the Quasar 3C 273

Wolfgang Kundt *Institut für Astrophysik der Universität, Auf dem Hügel 71, D-5300 Bonn, FRG*

Gopal-Krishna *Radio Astronomy Centre, Tata Institute of Fundamental Research, Post Box 1234, Bangalore 560012*

Received 1986 January 1; accepted 1986 June 14

Abstract. New observations of the jet in 3C 273 support and refine our earlier interpretation that (i) the mapped jet is $10^{6\pm0.3}$ yr old and grows at 0.6 to 0.75 times the speed of light, at an average angle θ of $(20 \pm 10)°$ with respect to the line of sight; (ii) its twin is not seen yet because arriving signals were emitted when it was some $10^{0.6\pm0.2}$ times younger; (iii) the fluid moving in the jet is an extremely relativistic $e^{\pm}$-pair plasma, of bulk Lorentz factor $\gamma \gtrsim 10^2$; (iv) the beam has swung in projection through some 10º; and (v) the small excursions (wiggles) of the jet around its average propagation direction result from a self-stabilizing interaction with the nonstatic ambient plasma. All other interpretations of which we are aware depend heavily on the ('beaming') assumption that the jet material radiates isotropically in some (comoving) Lorentz frame, an assumption which we consider unrealistic.

Key words: Jets—relativistic beaming—quasars—3C 273—lobe expansion

1. Introduction

The quasar 3C 273 is of outstanding interest because of its proximity, high lumenosity, hard γ-ray spectrum and because of its one-sided radio-optical-X-ray jet without a detection of 2-sided radio lobes. Conway, *et al.* (1981) and more recently Davis, Muxlow & Conway (1985) have concluded that the jet grows almost relativistically ($\beta_{\text{head}} \simeq 0.74$) and that the observed one-sidedness is intrinsic, *i.e.* that no counterjet exists. For the Lorentz factor γ of bulk motion in the jet they find a decrease from an initial value of $\gamma \simeq 6$ to a final value of $\gamma \simeq 1.5$. Roberts (1984) has argued that the jet is unlikely to consist of an $e^{\pm}$-pair plasma: Its bulk Lorentz factor would have to be high, 'hence' its inclination with respect to the line of sight is so small that its unprotected length would be excessive. Based on this conclusion and a lack of time variability of the γ-ray flux from 3C 273, Morrison, Roberts & Sadun (1984) have discussed the possibility of a hydrogen jet which produces γ-rays on collision with clumped 'warm' entrainments in the outer jet.

These models conflict with our own conclusion (drawn in 1980) that all observed jets are composed of extremely relativistic $e^{\pm}$-pair plasma moving through a thermal channel, with bulk Lorentz factor $\gamma \gtrsim 10^2$, and that one-sidedness is due to relativistic propagation and (generalized) beaming (*cf.* Kundt & Gopal-Krishna 1981, and Section 2).

The discrepancy between our results and those mentioned above can be traced back to the canonical relativistic beaming formula which is based on the assumption of the existence of a Lorentz frame with respect to which the radiation is isotropic (reviewed by Kellermann & PaulinyToth 1981; Begelman, Blandford & Rees 1984). In Section 2 we shall argue that this assumption can be very misleading; we concur with Lind & Blandford (1985) that the observed intensity distributions 'cannot be used to derive a precise value for the flow velocity because of the uncertainty in the [beaming] models'. Section 3 is devoted to a re-discussion of the pros and cons to extreme bulk velocities. In the succeeding sections we present what we think is a consistent description of the beam head dynamics, its size, age, lobe, and fine structure, including a discussion of the unseen counterjet. The uniqueness of 3C 273 is traced back to both its youth and favourable orientation.

2. Relativistic beaming

If the material composing a jet has an isotropic velocity distribution in its (comoving) centre-of-mass frame, then its radiation will likewise be isotropic, and its radiation with respect to any other frame can be obtained by applying a boost. In particular, the bolometric luminosity in the direction θ is proportional to δ^4 where δ is the well-known Doppler factor,

$$\delta := [\gamma(1-\beta\cos\theta)]^{-1} \simeq \left\{ \begin{matrix} 2\gamma/(1+\gamma^2\theta^2) & \text{for } \theta, \gamma^{-1} \ll 1 \\ 1/2\gamma & \text{for } \theta \simeq \pi, \beta \simeq 1 \end{matrix} \right\} \tag{1}$$

with $\gamma := (1-\beta^2)^{-1/2}$.

If, however, the velocity (or energy) distribution in the jet satisfies a power law with suitable cutoffs (to keep all meaningful intergrals finite), then the comoving velocity distribution is far from isotropic: It contains extremely relativivstic particles in the forward hemisphere, but none in the backward hemisphere. In order to see this, insert the formula for the composite Lorentz factor (strictly valid only for parallel velocities)

$$\gamma_{\text{lab}} \simeq \gamma_{\text{bulk}} \cdot \gamma(1 + \boldsymbol{\beta}_{\text{bulk}} \cdot \boldsymbol{\beta}) \tag{2}$$

where $c\beta_{\text{bulk}}$ is the comoving velocity and $c\beta$ the velocity with respect to the centre of mass, into the factorized power-law distribution of the radiating charges (averaged over 3-space)

$$dN = \gamma_{\text{lab}}^{-(2+g)}\, d\gamma_{\text{lab}}\, d^2\omega \tag{3}$$

in which $d^2\omega := f(\zeta) d^2\Omega$ describes the distribution over angles. The density $f(\zeta)$ should incorporate all perturbations of the beam, both by friction on the walls and by volume obstacles; it is strongly forward-peaked in the lab frame (though not as a function of the centre-of-mass quantity ζ). Using $d\gamma = \gamma^3 \beta d\beta$ one gets for $\beta_{\text{bulk}} \simeq 1$:

$$\gamma_{\text{bulk}}^{1+g}\, dN \simeq \gamma^{-(1+g)}(1+\zeta\beta)^{-(2+g)}[\zeta + \beta\gamma^2(1+\zeta\beta)]\, d\beta\, d^2\omega \tag{4}$$

with $\zeta := \cos(\beta_{\text{bulk}}, \beta)$. This distribution behaves as $\gamma^{-(2+g)}\, d\gamma$ for fixed $\zeta > 0$ and large γ, and cuts off near $\gamma = 1$ for $\zeta < 0$ in order to satisfy the centre-of-mass condition $0 = \int dN\gamma\beta$. In particular, the comoving energy in relativistic particles, $\int y^{-(1+g)}\, dy$, converges for (realistic values) $g > 0$, hence corresponds to a small percentage of the total energy for not too-small values of g.

In other words: the comoving velocity distribution of a relativistic beam need by no means be isotropic. Rather, a few particles in forward directions can have large momenta whereas all the others are extremely soft. The beaming pattern from such an ensemble has significant intensities in all forward directions (for suitable $f(\zeta)$) but very low intensities in backward directions. Such a beaming pattern can explain the frequent occurrence of jets in spite of extremely relativistic bulk velocities: With sufficient dynamic range, we see a jet for all angles $\lesssim 90°$. We even see often both the jets (which are needed to feed the symmetrically placed outer lobes), particularly in weak sources where the flow interacts more strongly with the channel walls. For such generalized distributions, the standard estimate $\gamma_{\text{bulk}} \simeq \theta^{-1}$ turns into an inequality: $\gamma_{\text{bulk}} \gtrsim \theta^{-1}$. All the estimates of γ_{bulk} given in the extended literature may have to be reinterpreted in this sense.

3. Pros and cons to high γs

In this communication, we do not want to repeat all the earlier arguments against small bulk velocities (Kundt & Gopal-Krishna 1981). The strongest among them (concerning γ_{bulk}) is our disbelief in the high efficiency of in-situ acceleration to the required ultrarelativistic speeds: Instead of in-situ acceleration, we prefer to think of in-situ deceleration, like for an electron beam illuminating a TV screen.

An independent argument in favour of large bulk velocities is the consideration that the jets emerge from deep inside the broad-line region where thermal emission lines indicate pressure in excess of 1 dyn cm^{-2}. At their termination points, in the heads, the inferred pressures fall mostly short of 10^{-8} dyn cm^{-2}. Consequently, the jet material is squeezed into its channels by an overpressure of more than 10^8 fold. In the absence of excessive friction–which would manifest itself by excessive surface brightness–and after expansion, the channelized plasma must acquire bulk velocities which are equal to its velocities (ordered or random) near the central engine (Landau & Lifshitz 1978). Such velocities are thought to be extremely relativistic, as inferred from the synchrotron spectra.

Now to the cons. It has been argued that extremely relativistic bulk velocities would lead to infinitesimally small scale heights h in the bends of the jets (Begelman, Blandford & Rees 1984; Phinney 1983). The argument is based on the inequality

$$\gamma_j \beta_j / \gamma_s \beta_s \lesssim (R/h)^{1/2} \tag{5}$$

where j, s stand for ‘jet’, ‘sound’ respectively, and R is the curvature radius. It assumes that the ratio of ram pressure to static pressure stays constant throughout a bend. But it is known from both theory and experiments with supersonic jets that ram pressure can converted into static pressure via shocks, such that the cross-section of the beam in the bends does not shrink considerably (Courant & Friedrichs 1976).

Another argument against extremely relativistic bulk velocities is the concern about inverse-Compton losses in the broad-line region: Energetic electrons and positrons would lose most of their energy in collisions with the thermal photon bath (Begelman, Blandford & Rees 1984). There is, however, also the inverse reaction: Photons of energy above 10^{11} eV can produce $e_{\pm}$-pairs on collision with *UV* photons, and so can photons above 10^6 eV on collision with protons and electrons. Moreover, ordered radial flows in the inner part of the broad-line region reduce the collision rate between electrons and photons (*cf.* Kundt 1982). More realistic calculations have to be made before this concern can be considered serious.

4. Beam propagation

Ram-pressure balance in the frame of the contact discontinuity suggests that the heads of strong young jets can advance relativistically into the ambient plasma:

$$\beta_h^2 \beta_j / (\beta_j - \beta_h)^2 = L / A c^3 \rho \tag{6}$$

can exceed unity. Here $c\,\beta_h$ is the speed of the head, $c\,\beta_j$ that of the jet, A its cross section, L is the total power flowing down the jet, and ρ is the ambient mass density, and ram pressure = $\rho_0\, c^2\, \beta^2\, \gamma^2$. Use has been made of Equations (8,10) below for the conversion from the laboratory frame to that of the contact discontinuity (= 'head').

For large β_h, only a small portion $\mathscr{L}$ of L can be dissipated so that the radiated power falls short of L. This reduction in the dissipated power is due to the fact that the needed momentum transfer for channel formation consumes almost the total infalling energy. Quantitatively, the mechanical power L_{mech} transferred to the ambient medium through the contact discontinuity exceeds the force $L\beta_j/c$ integrated over the advance rate $c\beta_h$, whence $L_{mech} \gtrsim L\beta_h\,\beta_j$, and

$$\mathscr{L} \leqslant L - L_{mech} \lesssim L(1 - \beta_h) \tag{7}$$

holds for $\beta_j \simeq 1$. Consequently, the dissipated percentage drops to zero in the extremely relativistic limit.

We can also derive a lower bound on the dissipated portion $\mathscr{L}$ from 4-momentum conservation under the assumption of elastic particle reflections in the comoving frame of the head.This result will not be needed explicitly below, hence the reader who is only interested in 3C 273 is advised to skip the rest of this section.

In order to derive the lower bound on $\mathscr{L}$, we work in the free-particle approximation in which extremely relativistic electrons flowing down the jet (with velocity $c\beta_j$) get reflected by the swept-up ambient plasma. In the frame comoving with the head, these electrons have velocities $c\beta_c$ obtained from the relativistic composition law:

$$\beta_c = (\beta_j - \beta_h)/(1 - \beta_j\beta_h) = \beta_j(1 - \beta_h/\beta_j)/(1 - \beta_h\beta_j). \tag{8}$$

In the laboratory frame, the reflected charges have velocities $c\beta_r$ given analogously by

$$\beta_r = (\beta_h - \beta_c)/(1 - \beta_h\beta_c) = -\beta_c(1 - \beta_h/\beta_c)/(1 - \beta_h\beta_c). \tag{9}$$

The corresponding Lorentz factors γ follow successively from Equation (2, 8, 9):

$$\gamma_c = \gamma_j\gamma_h(1 - \beta_j\beta_h) \simeq \gamma_j\gamma_h(1 - \beta_h) \tag{10}$$

$$\gamma_r = \gamma_h\gamma_c(1 - \beta_h\beta_c) \simeq \gamma_j(1 - \beta_h\beta_c)/(1 + \beta_h) \tag{11}$$

in which $y_j \gg 1$ has been assumed. The Lorentz factors of the reflected charges are therefore reduced by the factor $(1 - \beta_h\beta_c)/(1+\ \beta_h)$ compared with those in the jet.

We are now ready to calculate the maximal energy Q per particle (of rest mass m) which is available in the lab frame for dissipation. It follows from 4-momentum conservation during the collision on a heavy collective, of mass M ($\gg m$), whose velocity changes during the collision by $\Delta \beta_h$:

$$m[(\gamma\beta)_j - (\gamma\beta)_r] = M\Delta(\gamma\beta)_h, \tag{12}$$

$$m[\gamma_j - \gamma_r] = M\Delta\gamma_h + Q/c^2 \quad . \tag{13}$$

The identity $\gamma = (1 + \beta^2\gamma^2)^{1/2}$ and equation (12) imply $\Delta\gamma_h \simeq (m/M)\,\beta_h\,[(\gamma\beta)_j - (\gamma\beta)_r]$, whence with Equations (11), (13):

$$Q/\gamma_j mc^2 \simeq (1-\beta_h) - (1-\beta_h\beta_c)(1+\beta_h|\beta_r|)/(1+\beta_h)$$
$$\simeq 0.14\beta_h\,(1-\beta_h)\beta_j\gamma_j^{-2}\,10^{3.68\beta_h}, \tag{14}$$

the latter for $0.1 \lesssim \beta_h \lesssim 0.8$, $10 \lesssim \gamma_j \lesssim 10^4$ (obtained *via* numerical experiments); i.e. the energy Q per incoming particle energy $\gamma_j\,mc^2$ which is available for dissipation in this model is smaller than 0.3 per cent for velocities $c\beta_h$ in the given range whenever $\gamma_j \gtrsim 10^2$. Expression (14) is a lower bound on the dissipated portion $\mathscr{L}/L$.

5. The jet of 3C 273

In Fig. 1 we have drawn our knowledge about the jet, both from radio (Flatters & Conway 1985; Foley & Davis 1985; Davis, Muxlow & Conway 1985), infrared (Henry & Becklin 1984), optical (Lelievre *et. al.* 1984; Röser & Meisenheimer 1986) and X-ray data (D. Harris, personal communication). For a redshift of z = 0.158 and a present Hubble parameter of H_0 = 65 km s^{-1} Mpc^{-1} = $10^{-17.7}$ s^{-1} (Bartel *et. al.* 1985), its (projected) angular size of 23 arcsec corresponds to a projected length of $r_\perp$= 81 kpc $H^{-1}_{-17.7}$. For an inclination angle θ between 10° and 30°, its deprojected length $r = r_\perp/\sin\theta$ falls therefore between 500 and 160 kpc. We favour a value of r in the vicinity of 200 kpc because we expect the jet to be still in its near free expansion stage, yet well outside the halo of the host galaxy whose boundary may be indicated by the inner edge of the 'lobe' discovered by Davis, Muxlow & Conway (1985). For an average propagation speed $c\beta_h$ of the head, the favoured size r translates into an age $t = r/c\beta_h$ between 2×10^6 and 0.5×10^6 yr for $\beta_h \gtrsim 0.6$ (see below). β_h be better not smaller than 0.6.

On the VLBI scale, superluminal expansions have been observed with $\beta_\perp = 8\ H^{-1}_{-17.7}$ (Unwin *et. al.* 1985), suggesting relativistic motion at an angle between 10° and 20° with respect to the line of sight, depending on its bulk Lorentz factor:

$$\beta_\perp/\beta = \sin\theta/(1-\beta\cos\theta) \simeq 2/\theta, \tag{15}$$

the latter for $\gamma \gg 1$. But there is no guarantee that the inferred VLBI angle agrees with the large-scale inclination angle of the jet to within less than 10°, say. We are therefore still left with a considerable uncertainty concerning the largescale θ. Acomparatively short—and therefore young—jet implies a large θ (> 20°). Note that according to our ideas expressed in Section 2, we do not accept estimates of θ based on beaming arguments. An independent estimate, based on the width of the lobe, will be presented in the following section.

In order to assess the propagation speed $c\beta_h$ of the jet's head, we evaluate Equation (6):

$$\beta_h^2/(1-\beta_h)^2 \simeq L/Ac^3\rho = L_{45.5}/A_{43}\rho_{-29} \tag{16}$$

where we have inserted a total power of $L = 10^{45.5}$ erg s^{-1} flowing down the jet, a jet

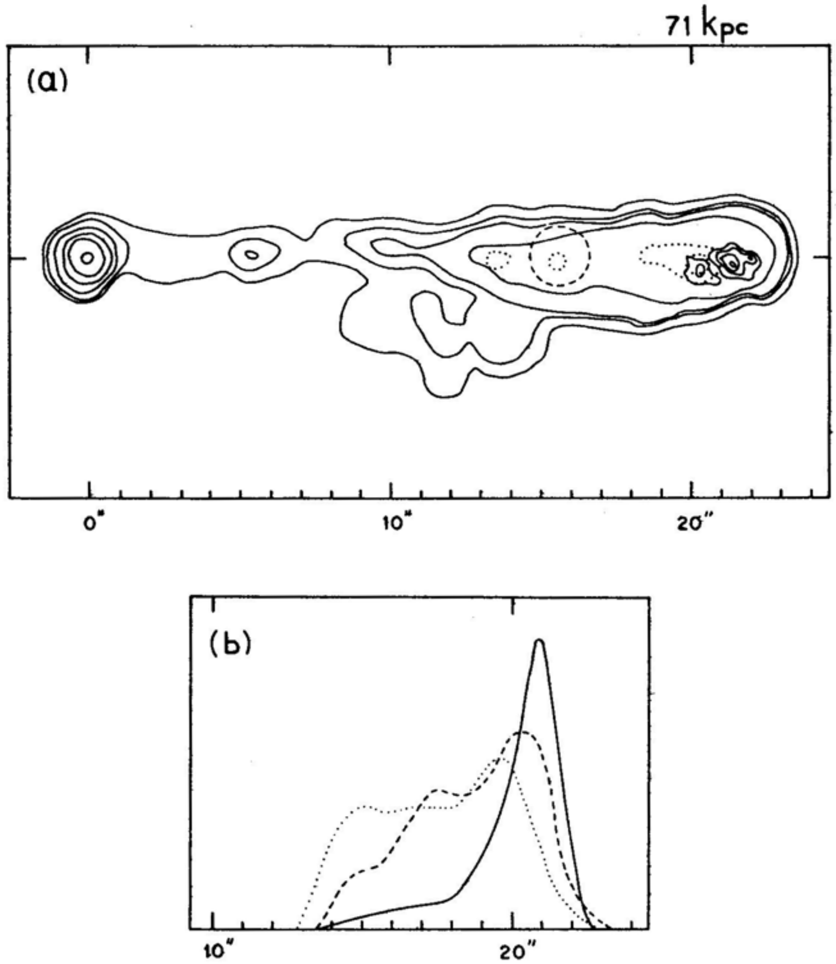

Figure 1. (a) Schematic of a radio-optical-X-ray overlay map of 3C273, rotated in position angle through 48°. Continuous lines are 4×10^8 Hz radio isophotes, reproduced from Foley & Davis (1985) and Davis, Muxlow & Conway (1985); dotted lines are optical isophotes, taken from Röser & Meisenheimer (1985); and broken lines are X-ray isophotes as communicated to the authors by D. Harris. The (upper) linear scale is based on H_0 = 65 km s^{-1} Mpc^{-1}. (b) The corresponding flux profiles along the ridge of the jet, where this time the broken curve represents the infrared profile (taken from Röser & Meisenheimer, 1985). All fluxes are plotted linearly but on different scales. The beam width for the radio profile is 1 arcsec, that for the optical-infrared profiles 2.4 arcsec.

cross-section (at termination) A of order $(\mathrm{kpc})^2$ and a (low) intergalactic hydrogen density of $\rho = 10^{-29}$ g cm^{-3} (*cf.* the optical map by Wyckoff, Wehinger & Gehren 1981, and anizares, Gordon & Fabian 1983). The jet's present radiated power (if only mildly beamed) is no larger than 10^{44} erg s^{-1}. But Equations (14) and (7) tell us that the radiated fraction can be small for a near-relativistic head velocity, and most of the synchrotron radiation is expected somewhat retarded, pending on the prevailing magnetic field strengths (Equation 19 below). Also, other well-studied sources suggest that the power fed into the jets by an AGN is of order 1 per cent, corresponding to

$L \lesssim 10^{46}$ ergs^{-1} (Bezler *et al.* 1984). For $L_{45.5}/\rho_{-29} \leqslant 3$, Equation (16) implies $\beta^2 \leqslant 3(1-\beta^2)^2$, whence $\beta_\mathrm{h} \leqslant 0.63$. For comparison, the equations $\beta^2/(1-\beta)^2 = 1, 10$ are solved by $\beta = 0.50, 0.76$ respectively.

Lower limits on β_h follow from the fact that Davis, Muxlow & Conway (1985) have not detected the counterjet (yet). We do not believe in intrinsic one-sidedness, observationally because of the overwhelming number of extended double sources, and theoretically because the viability of models which prevent the high-pressure relativistic plasma in the broad-line region from escaping to both sides of the feeding disc is far from established. But the counter jet would have been observed unless its head has not yet reached the boundary of the host galaxy's halo (where enhanced dissipation sets in, as evidenced by the approaching jet, see Fig. 1). Its 'present' age t_+ is restricted by retarded visibility:

$$t_+/t_- = (1-\beta_\mathrm{h}\cos\theta)/(1+\beta_\mathrm{h}\cos\theta) \tag{17}$$

in which $t_- = r/c\beta_\mathrm{h}$ is of order 10^6 yr. Postulating the counterjet to be shorter by at least a factor of 0.4 (*cf.* Fig. 1) means postulating t_+/t_- to be (roughly) $\leqslant 0.4$, whence $\beta_\mathrm{h}\cos\theta \geqslant 0.43$, or $\beta_\mathrm{h} \geqslant 0.49$ for $\theta \leqslant 30°$. This leaves us with the solid lower bound $\beta_\mathrm{h} \gtrsim 0.5$. On the other hand, the above constraint $\beta_\mathrm{h} \lesssim 0.75$ implies $t_+/t_- \gtrsim 0.16 = 10^{-0.8}$.

In deriving these constraints, we have tacitly assumed that β_h does not vary significantly during the lifetime of the young jet, an assumption that is suggested by Equation (16): the younger (shorter) jet has to push through an initially higher mass density ρ of the halo, yet with a smaller cross-section A ($\sim r^2$). Consequently, the 'velocity function' $L/A\rho c^3$—and hence β_h—are expected almost independent of r.

Why is the counterjet invisible? According to the canonical beaming formula, the forward-backward contrast in spectral flux $S_\nu \sim \nu^{-\alpha}$ is given by (*e.g.* Phinney 1985):

$$(\delta_-/\delta_+)^{2+\alpha} \simeq [2/(1-\beta_\mathrm{h}\cos\theta)]^{2+\alpha} \geqslant 10^2 \tag{18}$$

for $\alpha = 0.8$, $\theta \leqslant 20°$, and $\beta_\mathrm{h} \geqslant 0.65$. But as discussed in Section 2, the real contrast is expected much higher for a power-law distribution in energies for which hardly any relativistic charges are deflected into the backward hemisphere: charges moving with $\langle\beta_z\rangle \gtrsim$; 0.5, $\gamma \gtrsim 10^2$ do not radiate into directions k with $k_z < 0$ unless the distribution is relativistically hot. A quantitative estimate would depend on the (unknown) effective temperature and turbulence of the shocked beam plasma.

A possible worry concerns the invisibility of the counterjet's cocoon whose energetic charges can be expected to have relaxed, on an average, to an almost isotropic velocity distribution. Their number per beam length is, however, small in proportion to $A\rho(1-\beta_\mathrm{h})$ (according to Equations 7, 14) due to the beam's slimness and fast propagation, and their radiation is further reduced by a somewhat relaxed pressure when compared with the beam head. Moreover, there are enhanced inverse-Compton losses in the galactic bulge, as well as collisional losses to neutral hydrogen atoms which can traverse the magnetized channel walls. As we do not see the inner cocoon of the approaching jet, we need not expect to see that of the receding jet.

Two further constraints on the head velocity $c\beta_\mathrm{h}$ are worth mentioning. Within our model, the hard tail of the power-law distribution in particle energies degrades within the head, as inferred from the fact that the foremost optical hotspot lags behind the foremost radio hotspot. For an energy-density spectral index of 1.6, the share in the total energy of the degrading charges is of order 2 per cent, hence $\beta_\mathrm{h} < 0.98$ must hold according to Equation (7). Secondly, we compare the observed lengthscale of spectral softening in the head (Fig. 1) with the expected lengthscale for synchrotron losses. For the break frequency ν_b of synchrotron-emitting electrons in a transverse magnetic field $B_\perp$ one has:

$$\nu_b = 4\pi m_e ce/\sigma_T B_\perp^3 t^2 = 10^{14.6}\ \mathrm{Hz}\ (B_{\perp,-5})^3 t_{12}^2, \tag{19}$$

where $t := \gamma/\dot{\gamma}$ is the e^{-1} folding timescale of particle energies. An ageing of the optical-emitting charges on the (deprojected) lengthscale of 10 kpc= 10^{12} *s.c* (which is traversed almost relativistically) therefore corresponds to an average transverse magnetic field of 10^{-5} G, a small fraction of the inferred equipartition field strength of $\lesssim 1$ mG near the front end of the jet. We shall expand on this interpretation in Section 7.

The presence and morphology of the radio lobe will impose a seventh constraint on β_h which we discuss in the following section.

6. Swinging beams and the lobe of 3C 273

The jet of 3C 273 is much narrower than the lobes of most extragalactic radio sources. We interpret this narrowness as an indication of its young age: the cocoon has not yet had time to expand into a lobe. If the cocoon consists of extremely relativistic pair plasma, this plasma wants to expand at its sound speed $c_s = c/\sqrt{3}$ as long as its pressure exceeds $\rho c^2 = 10^{-8.5}$ dyn cm^{-2} ρ_{-29}. Put differently, the sideways expansion of a lobe (at speed $c\beta_1$) is expected not to a fall much behind its forward expansion:

$$\beta_l = \min \begin{Bmatrix} \beta_h \, (p_{\text{lobe}}/p_{\text{hotspot}})^{1/2} \\ 1/\sqrt{3} \end{Bmatrix} \tag{20}$$

where p_{hotspot} is the (highest) stagnation pressure and p_{lobe} is the more relaxed pressure in the extended lobe.

The half width b of the lobe must therefore fall short of the length Δr of the outer jet by β_l/β_h times the ratio t_l/t_h of their look-back ages. From $b/\Delta r \simeq \beta_l t_l/\beta_h t_h$, $\beta_h\ t_h$, cos $\theta = t_h - t_l$ (because of the finite light-travel time) and $\Delta r_\perp = \Delta r \sin\theta$, we get for the projected length ratio:

$$b/\Delta r_\perp \lesssim \beta_l(1-\beta_h\cos\theta)/\beta_h \sin\theta \leqslant \beta_l/\beta_h \tag{21}$$

for $\beta_h \geqslant 0.7$, $\theta \leqslant 20°$. The observed ratio $b/\Delta r_\perp$ of half width to projected length in the map by Davis, Muxlow & Conway (1985) is close to 0.13. We infer that the expansion speed of the lobe satisfies $\beta_1 \gtrsim 0.13\ \beta_h$ for $\beta_h \geqslant 0.7$. When compared with Equation (20), this estimate disfavours values of β_h below 0.6.

In this interpretation, we have ignored the fact that the 'lobe' mapped by Davis, Muxlow & Conway (1985) is offset from the present jet by order of 2*b.* As a transverse motion through 2*b* corresponds to a speed of $\gtrsim$ 0.3 *c,* the beam must have swung sideways almost relativistically, through some 10°. Of course, the sideways speed of a beam is only a phase velocity: It is realised by a bending of its inner part in conjunction with the grinding (ramming) of a new channel. Such a bending can be achieved by the Keplerian rotation of the inner halo of the host galaxy, at $r \gtrsim 1$ kpc, which tends to freeze the channel walls. It is clearly mapped in Fig. 3b of Davis, Muxlow & Conway (1985). (Note that we discard the hard-beam model in which individual blobs perform a quasi-straight-line motion. In the soft-beam model, a precession of the central engine could only influence the innermost part (< 1 kpc) of a jet.) When the direction of a beam is changed by almost corotation with the galactic halo, its succeeding segment is dragged across the (relativistic) lobe of the earlier channel for a long time where it finds little resistance. The resistance is resumed when it leaves the former lobe and has to ram a new channel.We interpretthis to happen atadistance of some 10 arcsec from the core.

This interpretation of the lobe offset from the jet by swinging was forced upon us by the (plausible) assumption that the jet propagates almost relativistically so that the observed sideways offset corresponds likewise to almost relativistic speed. No (intergalactic) wind can blow that fast. We would not feel confident about it had not a similar lobe offset been observed in Cyg A (for a slower propagation speed of the head, $\beta_h \simeq 0.1$; Perley, Dreher & Cowan 1985) and also in the form of 'emission bridges' in many sensitively mapped extended double sources (Leahy & Williams 1984). These emission bridges can occasionally take the shape of semi-circles, like in 3C111. In other cases they assume *Z* shape, like in 3C 52 or in the sources discussed by Ekers *et. al.* (1978).

Sources of Z shape have been often interpreted in terms of a precessing beam, based on the hard-beam model, or in terms of a deflected backflow. We prefer to interpret them as due to a swinging (soft) beam whose inner parts are dragged along by galactic rotation. This interpretation is similar to one given by Wirth, Smarr &Gallagher (1982) for the jets of close pairs of galaxies, with the distinction that we reject non-relativistic beam velocities and hence all forces on the beam other than entrainment of their channel walls by the respective galactic haloes.

Even more difficult to understand are the semi-circular sources. They have been interpreted as blown by a sideways intergalactic wind; but the required velocities are highly supersonic, comparable to or even larger than the head speed! Clearly, in all these cases the beam must have headed earlier into the directions of the 'older' lobes (as judged by their spectrum) and subsequently been swung into those of the younger lobes. A semi-circular morphology can result when the source starts out as a head-tail galaxy whose extragalactic beam segments are swept downstream by the intergalactic wind. When at some time in the past the output of the central engine strengthened, the beam got stiffer and straightened into the presently observed antipodal shape.

Why is no lobe observed inside the galactic halo of 3C 273, at projected separations $\lesssim$ 10 arcsec? The electron spillover which leads to the formation of a cocoon or lobe can only be kept small by a low ambient density, so that little momentum has to be transferred. We therefore interpret the missing inner cocoon as due to a cosmic-ray halo, or sector of the halo through which the jet has propagated. The cosmic rays generated by the young stellar population(concentrated in the disc,like in Cen A) drive the galactic fountain and 'boil out' through the halo, possibly in counter-stream to infalling hot cluster gas.

7. Structure of the beam head

As already indicated in Section 5, we interpret the spectral softening in the front end of the jet as due to synchrotron losses (in-situ deceleration; The optical radiation stops before the linear polarization (E-vector) flips from transverse to parallel, *i.e.,* before the transverse magnetic field is significantly enhanced.

The brightest radio hotspot is preceded by yet another radio hotspot, at projected separation $\simeq$ 0.6 arcsec, of width 0.5 arcsec (Foley & Davis 1985).Apparently, the mapped head has an almost cylindrical geometry, with a deprojected ratio $\gtrsim$ 0.3 of front width to length, where the foremost radio hotspot marks the position of the highest static pressure. This hotspot should be identified with the front end of the jet, which touches the contact discontinuity versus the ambient IGM. We infer that the conversion of ram pressure $p_{\rm ram}$ into static pressure p takes place between these two hotspots, throughout some 3 kpc. This distance is larger in proportion than the thickness of more familiar (solar-system) shock layers, yet similar to the size of hotspots in double radio sources (Miley 1980). The geometry may resemble that of the inner shock layer of the Crab nebula, for whose thickness Kundt & Krotscheck (1980) derived some 10^{17} cm based on the assumption that the optical wisps are a laser. Note that the beam plasma is extremely thin (and collisionless) and that the beam head has not had much time to relax.The bulk velocity of the jet material takes some 3 kpc for being decelerated from β_j to β_h.

The question then arises of which of the two radio hotspots has the higher intrinsic luminosity. If the pair plasma near the front end had an approximately isotropic velocity distribution, then the beaming formula (1) would predict that observers at angles θ smaller than θ_c = arcsin $\sqrt{2/(\gamma_h + 1)}$ = 66° (for β_h = 0.7) see more than average luminosity, in proportion to $\delta^{2.5+\alpha}$ for S_ν (Equation 1). In reality, beaming will be much narrower; ordered relativistic streaming (at $\beta_j \simeq 1$) would yield θ_c = arccos β_h = 45° for β_h = 0.7, because individual charges move at an average angle θ_c with respect to the bulk flow. Beaming is expected to be even narrower for the (less decelerated) brighter radio hotspot, so that its intrinsic luminosity is still higher than that of the foremost hotspot. (Note that in this prediction we disagree with the analysis by Flatters & Conway 1985). A correct evaluation of beaming depends sensitively not only on the involved bulk Lorentz factor and inclination angle but also on the intrinsic velocity distribution.

An important observation is the presence of an X-ray hotspot in the outer jet. If it is emitted in the same electromagnetic fields as the optical radiation, the corresponding synchrotron-radiating electrons (and positrons) would have to be so energetic that they could not have survived the journey from the central engine to their present location, because of the ubiquitous 3K background radiation, see Equation (19) with $B_\perp$ (background) = $10^{-5.4}$ G $(1 + z)^2$. The electrons responsible for the optical jet just barely make it. (We interpret essentially the whole optical radiation from the jet as synchrotron radiation, despite its low linear polarization in the inner part: linear polarization can be wiped out by superposition of the radiation from various charges passing at different angles through the line of sight.) Instead, we interpret the X-ray hotspot as inverse Compton radiation on the infrared/optical photon bath, which is more strongly beamed than the emitted synchrotron radiation because it is a volume phenomenon, rather than an interaction with the walls, produced by particles of Lorentz factor $\gamma_j \gtrsim 10^2$. We only see the radiation from those parts of the wiggling jet whose moving charges approach the line of sight within less than a degree. Note that a very similar beam-head structure—that of 3C33—has been recently mapped by Meisenheimer & Roser (1986).

8. Wiggling is self-stabilized

The jet of 3C 273 is not straight, nor is any other well-mapped extragalactic jet on any length-scale: there are wiggles throughout. We have interpreted the innermost bend of 3C273, on the scale of its galactic halo, as due to entrainment of its channel walls by the galactic rotation. It is then clear that the front end of the young jet must curve backwards, because ram pressures on a swinging beam grow quadratically with the lever arm (in a uniform medium), hence freeze the outer parts of a channel. But why the jet swing back again and oscillate around its initial direction, in a remarkably stable manner?

A partial answer to this question has been given by Hughes & Allen (1985) who hold small-scale clumping of the ambient plasma at the beam head responsible for small deflections. But why are the deflections stable, *i.e.* do not lead to an ever-increasing deviation? The answer is probably contained in the time-dependent, inhomogeneous manner in which the 'swept-up' ambient plasma reacts on the beam. This reaction is expected similar to the behaviour of the channel-wall material in the chimney of the Crab which appears not to thermalize the shock energy (Kundt 1983). When the propagating head pushes it at high speed, it continues to recede for some time after the head's sweeping until it is slowed down by the unperturbed medium; now the locally generated overpressure relaxes explosively. But sound speed of the IGM, of order $\gtrsim 10^8$ cm s^{-1}, is some 1 per cent of the head's speed, hence this explosive recoil sets in slowly and long after channel formation.

As a consequence of the recession of the newly swept ambient plasma, a fresh bend tends to straighten because successive generations of pushing charges arrive at increasingly larger distances from the local curvature centre. This sideways offset of the beam is always in the opposite direction to the former bend; it tends to deflect the beam head away from its swept-up buffer, initiating the next bend.

In this scenario, sideways offsets of the beam of order $d \tan \phi$ are expected after the traversal of a distance d at (small) inclination angle ϕ with respect to its former direction. When such an offset gets comparable to the beam width b, the beam head is deflected again. This naive model suggests for the typical distances d between bends

$$d = b \text{ ctg (22) } \phi, \qquad (22)$$

in satisfactory agreement with the observations. In other words: each bend of a beam is likely to trigger the formation of a counter-bend because of time-dependent pressure fluctuations in the channel wall plasma.

At times long compared to the head's sweeping, the bends are in danger of sharpening *via* centrifugal forces, like the meanders of a river. Now the created overpressure of the relaxing channel-wall plasma serves as a buffer, counteracting centrifugal grinding. nly in this way can we understand the rema rkable straightness of a (soft) beam. 3-D numerical simulations of jets should therefore be careful in modelling the time-dependent back reaction of the channel-wall material on the flow

9. Conclusions

Within our modified framework of relativistic beaming, we have explained the one-sidedness of the jet in 3C273 as due to almost relativistic propagation, with a head speed in the vicinity of 0.7 *c*, at some variable angle q between 10° and 30° with respect to the line of sight. The counterjet is too young to be visible: Most of its radiation is still beamed away from us. 3C273 is thereby exceptional both for its youth and orientation: age of $\simeq 10^6$ yr is some 10^2 times less than the estimated ages of typical sources, and an inclination angle of $\lesssim 30°$ of one of the two jets corresponds to $\lesssim 10$ per cent of the full spherical angle. Consequently, there should be no more than one source in a thousand similar to 3C 273.

How can future observations test this model? An obvious possibility is offered by a VLBI measurement of $\beta_\perp$, the transverse growth rate of the jet, which is expected marginally superluminal: $\beta_\perp = 10^{-0.2\pm0.2}$ (*cf.* Equation 15). According to Perley (1983), such a measurement might be possible in the foreseeable future. Another test would be to discover the (lobe of the) counterjet at a higher dynamic range.

Acknowledgements

W. K. wants to thank R. Cowsik, S. Falle, R. Porcas and P. Scheuer for discussions and T.I.F.R. and R.R.I. for their hospitality during his stay at Bangalore. Both authors are grateful to R. Nityananda for discussions, to Moxa H. for her enthusiastic typing and to a helpful referee for constructive criticism.

References

Bartel, N , Rogers, A. E. E., Shapiro, I. I., Gorenstein, M. V., Gwinn, C. R., Marcaide, J. M., Weiler, K. W. 1985, *Nature,* **318**, 25.
Begelman, M. C, Blandford, R. D., Rees, M. J. 1984, *Rev. Mod. Phys.,* **56**, 255.
Bezler, M., Kendziorra, E., Staubert, R , Hasinger, G., Pietsch, W., Reppin, C., Trümper, J., Voges, W. 1984, *Astr. Astrophys.,* **136,** 351.
Canizares, C. R., Stewart, C. G., Fabian, A. C. 1983, *Astrophys. J.,* **272,** 449.
Conway, R. G., Davis, R. J., Foley, A. R., Ray, T. P. 1981, *Nature,* 2**94**, 540.
Courant, R, Friedrichs, K. O. 1976, *Supersonic Flow and Shock Waves,* Springer, Berlin.
Davis, R. J., Muxlow, T. W. B., Conway, R. G. 1985, *Nature,* **318**, 343.
Ekers, R. D., Fanti, R, Lari, C., Parma, P. 1978, *Nature,* **276**, 588.
Flatters, C, Conway, R. G. 1985, *Nature,* **314**, 425.
Foley, A. R, Davis, R. J. 1985, *Mon. Not. R. astr. Soc.,* **216**, 679.
Hughes, P. A., Allen, A. J. 1985, *Mon. Not. R. astr. Soc.,* **214**, 399.
Henry, J. P., Becklin, E. E., Telesco, C. M. 1984, *Astrophys. J.,* **280**, 98.
Kellermann, K. I., PaulinyToth, I. I. K. 1981, *Ann. R. astr. Astrophys.,* **19**, 373.
Kundt, W. 1982, *Mitt astr. Ges.,* **57**, 65.
Kundt, W. 1983, *Astr. Astrophys.,* **121**, L15.
Kundt, W. 1984, *J. Astrophys. Astr.* **5**, 277.
Kundt, W., GopalKrishna 1981, *Astrophys. Space Sci.,* **75**, 257.
Kundt, W., Krotscheck, E. 1980, *Astr. Astrophys.,* **83**, 1.
Landau, L. D., Lifshitz, E. M. 1978, 6, Chapter X, Section 90.
Leahy, J. P., Williams, A. G. 1984, *Mon. Not. R. astr. Soc.,* **210**, 929.
Lelievre, G., Nieto, J. L., Horville, D., Renard, L., Servan, B. 1984, *Astr. Astrophys.,* **138**, 49
Lind, K. R, Blandford, R. D. 1985, *Astrophys. J.* **295**, 358
Meisenheimer, K., Röser, H.-J. 1986, *Nature,* **319,** 459
Miley, G. K. 1980, *Ann. R. Astr. Astrophys.,* **18,** 165.
Morrison, P. Roberts, D., Sadun, A. 1984, *Astrophys. J.,* **280,** 483
Perley, R. A. 1983 in *IAU Symp.* 110: *VLBI and Compact Radio Sources,* Eds R. Fanti, K. Kellermann & G. Setti, D. Reidel, Dordrecht, p. 153.
Perley, R. A., Dreher, J. W., Cowan, J. 1985, *Astrophys. J.,* **285**, L35.
Phinney, E. S. 1985, in *Proc. Santa Cruz Summer Workshop,* Ed. J. S. Miller, Univ. Sci. Books
R berts, D. 1984, *Astrophys. J.,* **285**, 64.
Röser, H.-J., Meisenheimer, K. 1986, *Astr. Astrophys.,* **154**, 15.
Unwin, S. C., Cohen, M. H., Biretta, J. A., Pearson, T. J., Seielstad, G. A., Walker, R. C , Simon, R. S., Linfield, R. P. 1985, *Astrophys. J.,* **289**, 109.
Wirth, A., Smarr, L., Gallagher, J. S. 1982, *Astr. J.,* **87**, 602.
Wyckoff, S., Wehinger, P. A., Gehren, T. 1981, *Astrophys. J.,* **247**, 750.