Magnetic field, reconnection, and particle acceleration in extragalactic jets

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Abstract. Extra-galactic radio jets are investigated theoretically taking into account that the jet magnetic field is dragged out from the central rotating source by the jet flow. Thus, magneto-hydrodynamic models of jets are considered with zero net poloidal current and flux, and consequently a predominantly toroidal magnetic field. The magnetic field naturally has a cycindrical neutral layer. Collisionless reconnection of the magnetic field in the vicinity of the neutral layer acts to generate a non-axisymmetric radial magnetic field. In turn, axial shear-stretching of connected toroidal field gives rise to a significant axial magnetic field if the flow energy-density is larger than the energy-density of the magnetic field. This can lead to jets with an apparent longitudinal magnetic field as observed in the Fanaroff-Riley class II jets. In the opposite limit, where the field energy-density is large, the field remains mainly toroidal as observed in Fanaroff-Riley class I jets. Driven collisionless reconnection at neutral layers may lead to acceleration of electrons to relativistic energies in the weak electrostatic field of the neutral layer. A simple model is discussed for particle acceleration at neutral layers in electron/positron and electron/proton plasmas.

Key words: extragalactic jets – magnetohydrodynamics – magnetic field reconnection – particle acceleration

1. Introduction

Magnetic fields have an important role in the appearance and also probably the dynamics of jets in radio galaxies and quasars (see, for example, Begelman et al. 1984; Asseo & Sol 1987). Recent observations have revealed many details in the observed brightness and projected magnetic field patterns (for example, Bridle \& Perley 1984; Perley 1985; Clarke et al. 1986; Owen et al. 1989; Reid et al. 1989). However, the actual three-dimensional structure of the magnetic fields is still unknown. The question of this structure and its variation along and across the jet, remains a fundamental open problem.

Another fundamental problem is the acceleration and/or reacceleration of electrons in jets. It is widely believed that the observed radiation is incoherent synchrotron radiation of relativistic electrons in a weak magnetic field. In some sources, for example, M 87 and 3C-273, reacceleration or in-situ acceleration is deduced from the fact that the synchrotron radiation lifetimes are very much less than the transit times from the central source. The correlation between the regions of high brightness and polarization (Perley 1985; Bridle 1986) suggests that particle reacceleration is connected with the ordering or amplification of magnetic fields. Different mechanisms of particle acceleration have been studied, for example, first-order shock acceleration and second-order Fermi acceleration in magnetohydrodynamical turbulence, magnetic reconnection, etc. but it is not known what mechanism accelerates electrons in jets.

There is a high dispersion in the properties of different jets. They differ in power, morphology, size, brightness and magnetic field patterns (Bridle \& Perley 1984). Their brightness can increase or decrease along the jet, they can be limb-brightened or limb-darkened, in some cases, regions of high brightness are observed at different distances along the jet (for example, Perley et al. 1984). The projected magnetic field is parallel or transverse to the jet, although much more complex structures of mixed parallel and transverse magnetic fields are often observed (Perley 1985; Bridle 1986).

To explain the observed projected patterns of brightness and polarization, simplified models of the three-dimensional magnetic field distribution have been proposed (for example, Chan \& Henriksen 1980; Laing 1981). Laing (1981) considered a jet as a cylinder with uniform density and pressure and assumed that magnetic field was either fully ordered and smoothly varying across the source (that is, uniform or helical field) or partially ordered, a field which was initially random on a small scale and subsequently sheared, stretched, or compressed anisotropically. Laing computed the emission and polarization for different models at different angles between jet axis and the line of sight and concluded that a high degree of linear polarization does not necessarily require an ordered field, but merely that the field be confined to a plane containing the line of sight (see also Laing 1980). Laing pointed out that both fully and partially ordered fields can produce the high fractional polarization observed. Numerical simulations of jets with a randomly tangled magnetic field connected passively with the flow through an ambient medium confirmed the ordering and amplification of magnetic field by compression and shear in the flow (Matthews \& Scheuer 1990a, b). Axisymmetric numerical simulations of jets with ordered passive and active magnetic field have been done by Clarke et al. (1986, 1989) and by Lind et al. (1989). These simulations revealed many details of predicted jet structure, including the formation of oblique shocks in the flows and the

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appearance of a strong annular shock, located near the outer edges of the contact discontinuity separating the shocked jet fluid from the shocked ambient gas (Lind et al. 1989). However, these axisymmetric simulations are artificial in that non-axisymmetric modes of instability are suppressed and the effective Reynolds numbers are very small.

Recent VLA observations of the nearby jets in Cen A (Clarke et al. 1986; hereafter CBF) and M 87 (Owen et al. 1989; hereafter OHC), and VLBI observations of jet in M 87 (Reid et al. 1989) have been done with unprecedented angular resolution and dynamical range. They revealed a complex pattern of filamentary and other fine structure in the jets. Both jets are limb-brightened practically everywhere along all the jet, including the regions of bright knots. Also, the parsec-scale region of the jet in M 87 is limb-brightened (Reid et al. 1989). Often, there is side-to-side brightening. In the regions of bright knots, banana-shaped structures were found in the Cen A jet. Filaments resembling parts of helical structures were found in the M 87 jet.

The observed brightness and polarization distributions must be understood in terms of electron acceleration and motion and the magnetic field configuration. The observed brightness distribution reflects the regions of the most effective acceleration, while the polarization indicates the projected magnetic field. Electrons may be accelerated directly, for example, in the electric field of a neutral layer, or by stochastic electric fields which likely occur near shocks, or in a turbulent plasma (Eilek & Hughes 1990). It is often supposed that acceleration of electrons in jets is due to the Fermi mechanism of particle acceleration (Christiansen et al. 1976; Axford et al. 1977; Krimsky 1977; Bell 1977, 1978; Blandford & Ostriker 1978). In both the cases of shocks and turbulent plasma the acceleration is effective if the particle gyroradius is larger than the size of typical magnetic field inhomogeneities or than the scale of the transition layer of the shock. Both values have a scale of the order of the thermal ion gyroradii, so that the most efficient acceleration is of ions, which are reflected from downstream and upstream matter without "seeing" the transition layer of the shock. As for electrons, it is unclear how (or whether) low-energy electrons can be accelerated by shocks or by turbulent plasma waves (see reviews of Blandford & Eichler 1987; Eilek & Hughes 1990; Jones & Ellison 1991). An electron can be accelerated effectively if its gyroradius is comparable to or larger than the smallest wavelengths of the Alfvén wave spectrum, that is, electron Lorentz factors \( \gamma_e \geq (m_e/m_p)\left(\frac{e\phi}{mc}\right) \), where \( v_a \) is the Alfvén speed, \( c \) the speed of light, \( m_p \) and \( m_e \) the proton and electron masses. This is compatible with the interpretation that galactic cosmic rays arise from strong shock waves in the interstellar medium (for example, Blandford & Ostriker 1978), because the observed energy in the cosmic ray electrons is negligible (<2%) compared with that in the protons. However, it is possible that thermal electrons are accelerated to relativistic energies by whistler waves which may arise in a plasma due to pitch-angle anisotropy driven instabilities (Melrose 1974). The observed limb-brightening of the outer jets in the strong FR II class and limb-brightening of well-resolved nearby jets is often interpreted as particle acceleration at the oblique shocks (for example, Clarke et al. 1986). However, acceleration of electrons at oblique shocks is also inefficient and pre-acceleration of electrons to relativistic energies is also necessary.

In this paper we consider models where the jet magnetic field is dragged out from the central object by the jet flow. The models imply that the jet magnetic field is predominantly toroidal and that this field naturally has a cylindrical neutral layer where \( B_0 = 0 \). Cylindrical neutral layers for poloidal magnetic field configurations have been considered earlier in applications to planetary magnetospheres (Zelenyi 1979) and laboratory plasma experiments (Marx 1968). However, the neutral layers in the toroidal magnetic field configurations have apparently not been studied previously. In such configurations boundary phenomena connected with spontaneous and driven reconnection of the magnetic field together with shearing by the flow may have important consequences in determining the apparent magnetic field and in accelerating electrons to relativistic energies. Previously, Königl & Choudhuri (1985) and Choudhuri & Königl (1986) have discussed the possibility of electron acceleration in radio jets at neutral layers resulting from resistive tearing. Reconnection and its possible role in accelerating electrons in radio sources has been discussed briefly by a number of authors (Blandford 1983; Begelman et al. 1984; Browne 1985; Ferrari 1984, 1985; Norman 1985; Kirchner 1988). Lesch (1991) has recently discussed the role of reconnection in accelerating electrons in accretion disks.

In Sect. 2, a model of magnetized jets with a cylindrical neutral layer is developed. Reconnection and the influence of shearing of the neutral layer are analyzed in Sect. 3. In Sect. 4, a model of driven collisionless reconnection is developed which leads to a power-law distribution of relativistic electrons. Comparison of predictions of the model with observations is made in Sect. 5. A summary of the results of the paper is given in Sect. 6.

2. Axisymmetric magnetic field

2.1. Basic equations

We use a cylindrical, inertial coordinate system \((r, \phi, z)\), with the \(z\)-axis coincident with the jets axis, and with \(z = 0\) the location of the central object. We first, consider axisymmetric \((\partial / \partial \phi = 0)\), and stationary \((\partial / \partial t = 0)\) jet flows and magnetic field configurations. As a zeroth approximation we neglect dissipative effects and turbulence, and assume that the flow is non-relativistic. We are interested here mainly in the structure of the magnetic field dragged out from the central source by the jet flow.

Thus, the basic equations are those of ideal magnetohydrodynamics (MHD):

\[
V \cdot (\rho V) = 0
\]

\[
V \times B = \frac{4\pi}{c} J,
\]

\[
E + \rho \times B/c = 0,
\]

\[
V \times E = 0,
\]

\[
\rho (V \cdot V) = - \nabla p + J \times B/c,
\]

\[
V \cdot B = 0.
\]

Here, \( V \) is the flow velocity, \( \rho \) is the density, and \( p \) is the pressure of the plasma which may include a contribution due to relativistic particles. We assume that the jet flow speed is much larger than the escape speed for the gravitation potential so that gravitational force can be neglected in Eq. (5).

With the assumptions made,

\[
B = B_r + B_\phi \hat{\phi}.
\]
where \( B_\phi \) is the toroidal magnetic field, and where the poloidal magnetic field is

\[
B_p = \frac{\nabla \times \left( \frac{\Psi \phi}{r} \right)}{r},
\]

\[
= \left( -\frac{1}{r} \frac{\partial \Psi}{\partial z} \right) \hat{\phi} + \left( \frac{1}{r} \frac{\partial \Psi}{\partial r} \right) \hat{z},
\]

with \( \Psi(r, z) = r A_\phi(r, z) \) the magnetic flux function. Note that \((\mathbf{B} \cdot \nabla) \Psi = 0 \) or \( \mathbf{B} \cdot \nabla = 0 \) so that \( \Psi(r, z) = \text{const} \) labels a three-dimensional “flux surface” on which a given set of magnetic field lines lie. Alternatively, \( \Psi(r, z) = \text{const} \) labels the poloidal projection of a field line.

We let \( \alpha(z) \) denote the effective jet radius, and \( z \equiv [r/a(z)]^2 \) a dimensionless radial variable. Then, a simple Ansatz for the poloidal magnetic field is

\[
\Psi(r, z) = \Psi_0(z),
\]

so that

\[
B_\phi(r, z) = \frac{a^{1/2}}{r^2} \frac{dR}{dZ} B_\phi(z),
\]

\[
B_z(r, z) = \frac{1}{r^2} B_z(z).
\]

Here, \( R \equiv a(z)/a_0 \) is the dimensionless jet radius; \( a_0 = a(z=0) \) is the initial jet radius; \( Z = z/a_0 \) is the dimensionless axial distance, and \( B_\phi(z) = (2a^2/3d) \frac{d\Psi_0}{dz} \) is the initial axial magnetic field profile. Note that, it is not necessary to assume that \( dR/dZ \) is small compared with unity.

The poloidal magnetic flux carried by the jet is a constant,

\[
\Phi_p = \int_{z = \text{const}} d^2 x B_p(r, z) = \text{const}.
\]

Equation (2) and (3) imply the “frozen-in” field equation

\[
\nabla \times (\mathbf{e} \times \mathbf{B}) = 0.
\]

In turn the poloidal component of this equation implies that \( e \times \mathbf{B} = 0 \) or \( e_p \propto B_p \), where \( e_p = (v_x, 0, v_z) \) is the poloidal flow velocity. Because \( \nabla \cdot (\rho \mathbf{e}) = \nabla \cdot (\rho e_p) = 0 \), we have in general

\[
\mathbf{e}_p = \frac{F(\Psi)}{4\pi \rho(r, z)} B_p(r, z),
\]

where \( F(\Psi) \) is an arbitrary function of \( \Psi \) (see Lovelace et al. 1986). Therefore, from Eqs. (7),

\[
v_x(r, z) = a^{1/2} \frac{dR}{dZ} v_x(r, z).
\]

Further, we assume the following factorization of the dependence,

\[
v_x(r, z) = v_0 a(z) \Omega(z),
\]

\[
v_z(r, z) = v_x f_x V_z(z),
\]

\[
\rho(r, z) = f_x(z) \bar{\rho}(z).
\]

Here, \( \bar{\rho} \) is the average jet density at \( z \); \( f_x \geq 0 \) is the dimensionless mass-density “profile function” normalized such that \( \int_0^\infty dz f_x = 1 \); \( \alpha(z) \) is the initial jet rotation rate which is equal to the rotation rate of the disk so that \( \Omega(z=0) = 1 \); \( f_x \) is the dimensionless axial-velocity profile function; and \( v_x = \text{const} \) is a reference velocity which we assume to be of the order of \( c_\alpha a_0 \).

The toroidal component of the frozen-in field equation implies

\[
\nabla \times (\mathbf{e} \times \mathbf{B}) = \nabla \times (\omega \mathbf{B}_p) = \nabla \times \{ (\omega - \omega_0(a(z)) B_z \}
\]

\[
= -g(\alpha) \frac{a^{1/2} [1 - \Omega(z)] B_z}{R(Z)} V_z(Z),
\]

where \( g = \alpha_0(a_0)/[v_x f_x(a_0)] \geq 0 \) is a dimensionless profile function with \( g = \mathcal{O}(1) \). Notice that we have introduced \( \omega_0 \) in Eq. (10a) so as to give \( (\omega - \omega_0) / v_x \propto \text{const} \) as \( z \to 0 \), that is, at the disk surface where \( \omega = \omega_0 \) and \( v_x = 0 \). The toroidal field (10b) results from the differential twisting (as a function of \( z \)) of the poloidal magnetic field. The equation corresponds to \( B_\theta/B_z = r (\omega - \omega_0) dt/dz \).

We assume that the jet originates from the inner region of an accretion disk around a massive black hole. Thus, it is appropriate to take \( a_0 \sim 6GM/c^2 \approx 10^{14} \text{cm}/10^8 M_\odot \). Consequently,

\[
R \sim 3 \times 10^6 \left( \frac{\theta}{0.1} \right) \left( \frac{z}{1 \text{kpc}} \right) \left( \frac{10^8 M_\odot}{M} \right),
\]

where \( M \) is the black hole mass, and \( \theta \) is the half-angular width of the jet at the axial distance \( z \). For all distances of interest for radio jets (\( z \gg c \)), \( R \) is much larger than unity. The conservation of angular momentum (LBC) then implies that \( |\omega/\alpha_0| = |\Omega| = \mathcal{O}(1/R^2) \ll 1 \). Furthermore, for large \( R \), the axisymmetric poloidal field is negligible compared with the toroidal field. That is, \( |B_\theta/B_z| \sim (V_z/g)/(1/R) \ll 1 \). With the estimate \( B_\theta \sim 10^3 \mathcal{G} [M/(M_\odot/yr)]^{1/2} [10^8 M_\odot/M] \) (Lovelace 1976), we obtain

\[
B_\theta \sim 0.3 \mathcal{G} \left( \frac{\theta}{0.1} \right) \left( \frac{1 \text{kpc}}{z} \right) \left( \frac{M}{M_\odot/yr} \right)^{1/2},
\]

where \( M \) is the mass accretion rate onto the black hole.

We assume that the time-scale on which the jet is set up is short compared with the magnetic diffusion time in the medium through which it propagates through a conducting medium where the magnetic Reynolds number is much larger than unity. Analogous conditions have been studied in laboratory experiments where it is found that plasma beams propagate through a conducting background medium with negligible or small net current in their direction of motion (Miller 1985). Therefore, the jet carries zero net poloidal current,

\[
\int_{z = \text{const}} d^2 x \mathbf{j}_p \approx 0.
\]

Notice however that a small magnetic diffusivity will transport current density out of the jet channel. This non-ideal effect can give the jet a small net poloidal current. In contrast, Benford (1978) has argued that the jets carry a large net poloidal current.

We assume that the poloidal magnetic field of the jet is dragged out of the active galaxy nucleus by the jet flow. Then, because of the large magnetic Reynolds number, the jet carries zero net poloidal magnetic field, that is,

\[
\int_{z = \text{const}} d^2 x B_p = \Phi_p \approx 0.
\]

This is in contrast with some of our earlier work (for example, LBC, LWS) where the net poloidal flux of the jet was assumed to close through the intergalactic medium. Previously, Chiuderi et al. (1989) have argued for MHD jet equilibria with zero net poloidal current, but with non-zero net poloidal flux.
As an illustration of a magnetic field satisfying Eqs. (13) and (14), we take
\[
\Psi_r(z) = \frac{1}{1 + z^2},
\]
(15)
where \( B_\theta = B_z(r \to 0, z = 0) \) provides a reference value for the poloidal field threading the disk, and where \( p \) is a constant. The poloidal and toroidal field components follow from Eqs. (7) and (10). Figure 1 illustrates the radial dependence of the field components for a case where \( \theta_0(a_0 / v_a) = \text{const.} \) and \( f_\phi(z) = \left(1 + z^2 \right)^{-1} \) is the axial velocity profile of the jet with \( q = p - 1 \). Note that the dominant field is toroidal and that the field has a neutral layer at \( r = r_a(z) \) where \( B(r_a, z) = 0 \).

2.2. Conserved fluxes of jet
The ideal MHD Eqs. (1)–(6) imply the conservation of the fluxes of mass, angular momentum and energy of the jet. The conservation of mass gives
\[
\dot{M}_j = \int_{z = \text{const.}} d^2 x \, \rho v_z = \text{const.},
\]
(16a)
For a proton/electron plasma of number density \( n \),
\[
\dot{M}_j \sim 2.2 \times 10^{-2} \frac{M_{\odot}}{\text{yr}} \left( \frac{n}{10^3 \text{ cm}^{-3}} \right) \left( \frac{0.1}{z \text{ kpc}} \right)^2 \left( \frac{v_z}{0.1c} \right). \]
(16b)
Assuming \(|\Omega| \ll 1\), the conservation of angular momentum is
\[
\mathcal{F}_L = \int_{z = \text{const.}} d^2 x r \rho v_z v_\phi - B_\phi B_z / 4\pi,
\]
(17)
where \( k_1 \) and \( k_2 \) are constants with \( k_1 \sim a_0^2 \theta_0 \) and \( k_2 \sim a_0^2 B_z^2 \) \((\theta_0 / v_a)\). The two terms in Eq. (17) can be of comparable magnitude. However, the azimuthal velocity of the matter is negligible, \(|v_\phi / v_a| \ll R^{-1} \ll 1\).

![Fig. 1. Radial profiles of the magnetic field components and the axial flow velocity of the axisymmetric model jet. For the distances of interest the axial field \( B_z(r, z) \) is negligible compared with the toroidal field \( B_\phi(r, z) \). The magnetic field is zero on a cylindrical neutral layer, \( r = r_a \). For the figure, \( p = 6 \) and \( q = 2 \).](image)

The conservation of energy gives
\[
\mathcal{F}_E = \int_{z = \text{const.}} d^2 x \, \rho v_z \left( \frac{1}{2} v_z^2 + w \right) + \frac{c}{4\pi} \int_{z = \text{const.}} d^2 x \, E_z B_\phi = \text{const.},
\]
(18a)
where \( w = \gamma / (\gamma - 1) (p / \rho) \) is the enthalpy per unit mass of the gas. The perfect conductivity, Eq. (3), gives \( E_z = v_z B_\phi / c \) for \(|\Omega| \ll 1\). We let \( \frac{1}{2} v_z^2 + w = \frac{1}{2} v_z^2 (1 + A) \) with \( A = (v_z / c)^2 + (c / v_z)^2 / (\gamma - 1) \), where \( c \) is the adiabatic sound speed. For a proton/electron plasma, we have
\[
\mathcal{F}_E \approx 3 \times 10^{46} \frac{\text{erg}}{s} \frac{M_{\odot}}{10^{-2} \text{ yr}} \left( \frac{v_z}{0.1c} \right)^2 (1 + A) + 0.1 \left( \frac{10^{-3} \text{ cm}^{-3}}{n} \right) \left( \frac{B_\phi}{0.3 \text{ mG}} \right)^2.
\]
(18b)
Note that the second term of \( \mathcal{F}_E \), the Poynting flux, is positive even though \( B_\phi \) may have both polarities as shown in Fig. 1. The Poynting flux may carry a significant fraction of the jet energy flux as we discuss in the next subsection. Notice that Eqs. (16)–(18) remain applicable to an "impure" electron-positron plasma if the ion component dominates the plasma inertia.

2.3. Radial virial equation
If the jet radius changes relatively slowly with axial distance, \( \text{d}R / \text{d}z \ll 1 \), Eq. (5) for a stationary jet equilibrium simplifies to
\[
0 \approx - \frac{\partial}{\partial r} \left( \frac{B_z^2}{8\pi} \right) - \frac{1}{8\pi r^2} \frac{\partial}{\partial r} (r^2 B_\phi^2).
\]
Multiplying this equation by \( 2\pi r^2 \text{ dr} \) and integrating over the jet cross-section gives
\[
2\pi r^2 \rho \text{exp} = 2 \int_{z = \text{const.}} d^2 x \left( \rho \text{jet} + \frac{B_z^2}{8\pi} \right),
\]
(19)
where \( \rho \text{exp} \) is the pressure of the external medium, and \( r_a(z) \) is the effective outer radius of the jet. In the next section, we argue that the apparent jet radius is the radius of the null of the toroidal magnetic field which occurs at \( r = a(z) \). As Fig. 1 illustrates, the jet flow and magnetic field can extend well outside of \( r = a \) so as to make \( r_a \) significantly larger than \( a \).

Notice that the toroidal magnetic field drops out of Eq. (19). The magnetic pressure \( B_\phi^2 / 8\pi \) is exactly offset by the tension of the toroidal field lines. Thus, the energy-density in the toroidal field can be of the same order of magnitude as the kinetic energy-density of the axial flow. Therefore, the two terms in the energy flux Eq. (18b), may be comparable. Notice that the Poynting flux is usually omitted in the discussions of jets (for example, Bridle 1986).

3. Formation of \( B_\parallel \) by non-axisymmetric reconnection and shearing
3.1. Reconnection in a cylindrical neutral layer
The discussion of Sect. 2 leads to the conclusion that the dominant axisymmetric magnetic field of the jet is toroidal and that
this field has a cylindrical neutral layer at \( r = r_o(z) \) where \( B\eta(r_o, z) = 0 \). Neutral layers are known to be unstable to spontaneous or driven reconnection under different conditions in a variety of laboratory and space plasma configurations. By reconnection we refer here to a change in topology of the magnetic field which leads to field lines connecting spatially separated plasma elements which were not initially connected. Because of the low densities of astrophysical jets the reconnection is collisionless. Episodic collisionless reconnection in the earth’s magnetotail is thought to be responsible for magnetic substorms (Schindler 1974; Galeev & Zelenyi 1976; Hones et al. 1984). Spontaneous reconnection involves perturbations dependent on the coordinates transverse to the current density \( (r, \phi) \), and it is driven by magnetic forces so that the basic speed is the Alfvén velocity. Driven reconnection can occur as a result of secondary, non-axisymmetric flows in the jet or as a result of turbulence in the jet flow. In addition, reconnection may occur as a result of a combination of magnetic forces and velocity-shear instabilities (the Kelvin–Helmholtz instability) in a boundary region of the neutral layer for perturbations which depend on \( r, \phi, \) and \( z \).

The linear stability theory for the onset of reconnection is involved and controversial (see, for example Wang et al. 1990). A treatment of this problem for jets is beyond the scope of the present work. Instead, we assume that reconnection occurs and consider the resulting magnetic field configurations. We first consider the field in the absence of shear in the axial velocity. In the vicinity of the neutral layer the equilibrium field is represented by the Harris (1962) solution

\[
A^\phi(r) = - B_n \Delta \ln \left[ \cosh \left( \frac{r - r_o}{\Delta} \right) \right],
\]

where \( A^\phi \) is the vector potential, \( B_n \) is the equilibrium toroidal field far from the neutral layer, and \( \Delta < r_o \) is the neutral layer thickness. Omitting for the moment a possible \( z \) dependence, the perturbation is taken to be

\[
\delta A^\phi = B_1 \Delta \cos(m\phi)/\cosh \left( \frac{r - r_o}{\Delta} \right),
\]

where \( B_1 \ll B_n \) is the perturbation amplitude, and \( m = 1, 2, \ldots \) is the toroidal mode number. The total vector potential, \( A^\phi + \delta A^\phi \), evidently corresponds to a magnetic field with “0” points at \( r = r_o \) and \( \phi = 0, 2\pi/m, 4\pi/m, \ldots \) and “X” points at \( r_o \) and \( \phi = \pi/m, 3\pi/m, \ldots \). There are \( m \) magnetic islands surrounding each “0” point with the radial extent of an island being \( r \approx r_o \pm \Delta(2B_1/B_n)^{1/2} \). Figure 2 illustrates the field topology resulting from reconnection. Note that the radial width of the islands has been exaggerated. The time-scale for spontaneous reconnection is roughly \( \tau_r = r_o/(m\nu_A) \), where \( \nu_A = |\partial B^\phi/\partial r| \Delta(4\pi \rho)^{-1/2} \) is the Alfvén speed.

### 3.2. Shearing of reconnected field

The jet will in general have shear in the axial velocity across the neutral layer, \( \partial v_z/\partial r \neq 0 \). This shear will act to axially stretch the field pattern of Fig. 2. We assume for the moment that the stretched field does not affect the flow so that \( v \approx v_z(r) \hat{z} \). Thus,

\[
\frac{\partial B}{\partial t} = \nabla \times (v \times B),
\]

or

\[
\frac{\partial B}{\partial t} \approx B_n \left( \frac{\partial v_z}{\partial r} \right)_n \hat{z}. \tag{20}
\]

Thus, \( B_n \) is independent of time and \( B_n(r, \phi, t) = t \left( \frac{\partial v_z}{\partial r} \right)_n B(r, \phi) \), where the \( n \) subscript indicates evaluation at \( r = r_o \). Figure 3 illustrates the growth of the axial field which results from the velocity shear. As in Fig. 2, the radial width of the field structure has been exaggerated.

The growth of the axial field will of course be limited once it becomes dynamically significant for the flow. This limit is reached when

\[
B_n^2 \approx 4\pi \rho \left( \frac{\partial v_z}{\partial r} \right)_n (\Delta r)^2, \tag{21}
\]

where \( \Delta r \) is the radial half-width of the magnetic islands. The relative significance of the axial field can be measured by comparing it with the axisymmetric toroidal field within a distance \( \Delta r \).
of the neutral layer,

\[
(B^2_{\text{z}})_{\text{max}} \equiv \frac{4\pi \rho_0 \left( \frac{\partial v_z}{\partial r} \right)_{\text{max}}^2}{(\partial B_{\text{z}}/\partial r)_{\text{max}}^2} \equiv \beta_n.
\]  

(22)

For comparison, the ratio of the energy-density of the flow to that in the toroidal magnetic field is

\[
\beta \equiv 4\pi \int \frac{\partial^2 \rho v_z^2}{\partial^2 \mathbf{B}_5^2}.
\]  

(23)

If the radial profiles of \( \rho, v_z, \) and \( B_5 \) are similar for different jets, then \( \beta_n = K \beta \) with \( K = \text{const.} \) A rough estimate is: \( K \sim (\Delta r/r_\text{c})^2 \), where \( \Delta r \) is the above mentioned neutral layer thickness, and \( r_\text{c} \) is the length-scale characterizing the velocity shear near \( r_\text{c} \). We assume that the neutral layer is thin in the respect that \( K \ll 1 \). For a “strong” toroidal magnetic field we have \( \beta_n \ll 1 \). In this limit the axial magnetic field due to stretching is negligible. The time-scale for spontaneous reconnection \( t_s \) is less than shear time-scale \( t_{\text{sh}} = (\partial v_z/\partial r)^{-1} \). In the opposite limit, the toroidal field is “weak” in the sense that \( \beta_n \gtrsim 1 \), and field stretching near the neutral layer can give an axial field comparable to or larger than the toroidal field. In this limit the time-scale for spontaneous reconnection may be of the order of the shearing time-scale. If \( t_s > t_{\text{sh}} \), reconnection may still occur if it is driven by secondary or turbulent flow, or if the perturbations have an axial dependence, the Kelvin–Helmholtz instability. For \( \beta_n \gtrsim 1 \), notice that there can be reconnection between regions of oppositely directed axial magnetic field. The transformation of toroidal field into axial field by shear stretching of random magnetic loops has been discussed previously by Begelman et al. (1985).

Independent of the value of \( \beta_n \), reconnection of the toroidal field across the neutral layer can give an important off-diagonal contribution to the momentum flux-density, \( T_{\text{e}z} = -\langle \delta B_{\text{e}z} \delta B_{\text{e}z} \rangle / 4\pi \), which acts as a “magnetic friction” or viscosity. An analogous magnetic friction has been discussed in connection with the anomalous viscosity in accretion disks (see, for example, Eardley & Lightman 1975; Coroniti 1981).

### 3.3. Helical structures

Reconnection of the magnetic field can occur as a result of the Kelvin–Helmholtz instability due to velocity shear (Ray & Erskovich 1983; Choudhury & Lovelace 1986) in the vicinity of the neutral layer. This instability is driven by the free energy associated with the velocity shear as well the free energy of the magnetic field. Thus, the perturbations depend in general on \( r, \phi, \) and \( z \). The linear perturbations have the general form of helical structures, \( \delta B_z = f_z(r, z, \phi, k_z) \), where \( k_z = 2\pi/\lambda_z \) with \( \lambda_z \) is the axial wavelength of the perturbation. Figure 4 illustrates the nature of a perturbation. The angle between the centerline of the perturbation or structure and the \( z \)-axis is \( \theta = \tan^{-1}(k_z r_\text{c}) \).

The shear in the axial velocity causes this structure to rotate about its centerline at an angular rate \( \omega_k \sin(\theta)/2 \), where \( \omega_k = -\langle \partial v_z/\partial r \rangle_{\text{max}} \) is the vorticity of the flow at \( r = r_\text{c} \). An initial, purely toroidal magnetic field passing through the structure will rotate about the centerline at the rate \( \omega_k \sin(\theta)/2 \). As a result of this rotation the average direction of the magnetic field within structure will coincide with structures centerline.

4. Driven collisionless reconnection and electron acceleration

The direct acceleration – as opposed to Fermi multiple bounce acceleration – of electrons and ions in electrostatic and inductive electric fields has been investigated by many authors over the last several decades. The high conductivity expected for the plasma in jet implies that \( E + \mathbf{v} \times \mathbf{B}/c = 0 \) or \( E \cdot B = 0 \) so that the electric field appears to cause only a drift perpendicular to \( B \) rather than any acceleration. One broad group of direct acceleration models treats time-dependent fields and involves an important role for the inductive electric field. The episodic acceleration of electrons and ions observed to be associated with substorms of the earth’s magnetotail have been interpreted in terms of spontaneous, “explosive” collisionless reconnection of the thought to be thick neutral layer (thickness \( \sim \) ion gyro radius) (Galeev et al. 1978; Galeev 1979, 1984; Zelenyi et al. 1984; Taktakishvili & Zelenyi 1988). The model time dependences of the electric and magnetic fields of the neutral layer have been used to derive the energy distributions of energetic test particle electrons and ions. However, a universal distribution is not found and the time dependences of the fields and the particle acceleration remain to be confirmed by particle simulations. Particle-in-cell simulations of spontaneous collisionless reconnection of thin neutral layers (thickness \( \sim \) electron gyro radius) (Hewett et al. 1988) do not show significant particle acceleration. The reconnection experiments of Stenzel & Gekelman (1979) do not mention the presence of energetic particles.

Related to the ideas of Galeev and Zelenyi are the studies of explosive coalescence of magnetic islands (Sakai & Ohswa 1987; Tajima & Sakai 1989). This work develops simple analytic two-fluid (electrons and ions) MHD equations for the coalescence of two regions of oppositely directed magnetic field using the earlier one-fluid MHD approach of Imshennik & Syrovatskii (1967). They obtain approximate agreement with their MHD simulations. However, the determination of the energy distributions of energetic particles relies as yet on the test particle approach and this indicates a distribution with an exponential dependence on energy (Syrovatskii 1981). Also, in the category of time-dependent particle acceleration models is the magnetic pinch model (Troubnikov 1990). Troubnikov obtains analytic solutions to the equations of relativistic MHD for the pinching of a current carrying filament and finds a distribution of accelerated protons in approximate agreement with the observed cosmic ray distribution. However, the generalization of this work to the collisionless plasma conditions of astrophysical jets remains to be done.
A second group of acceleration models considers time-independent or quasi-stationary configurations. We mention but do not discuss here the possibility of particle acceleration in relativistic double layers (see, for example, Borovsky 1986; Raadu 1989). Because \( \mathbf{E} \cdot \mathbf{B} = 0 \), most of the time-independent models consider the regions near neutral points of the plasma flow where \( v = 0 \). The models involve reconnection and/or annihilation of the magnetic field driven by the plasma flow. The driven reconnection is in contrast with spontaneous reconnection which appears as an instability without plasma inflow. Over the last several decades many theoretical and simulation studies have been made of driven reconnection based on the equations of resistive MHD (Sweet 1958; Parker 1963; Petschek 1964; Syrovatskii 1971; Vasyliunas 1975; Biskamp 1986). Most of the volume of the plasma is described by the equations of ideal MHD while in only a small volume near the stagnation point is the resistivity important.

Less theoretical attention has been directed to driven collisionless reconnection or annihilation (Alfvén 1968; Dessler 1968, 1971; Speiser 1970; Cowley 1971, 1973; Vasyliunas 1980; Lesch 1991). For describing the Alfvén (1968) model, consider a uniform planar neutral layer with current density, \( \mathbf{J} = J_y \mathbf{e}_z \) with \( J_y \leq 0 \) so that the magnetic field is \( \mathbf{B} = B_y \mathbf{e}_z \), with \( B_y (y \geq 0) \geq 0 \). In driven reconnection, plasma inflows towards the neutral layer, \( v_y (y \geq 0) \geq 0 \). Thus, there is an associated electric field \( E_z = -\frac{1}{c} \frac{d}{dy} \mathbf{B} \times \mathbf{B} = +v_y(c/B_y)B_x \mathbf{e}_z \), where \( E_x = v_y B_x/c = \text{const.} > 0 \) for quasi-stationary conditions. When an electron or ion flows for within a distance of order \( \delta_y = [m_i c^2 |\varphi_e(\partial B_r/\partial y)|]^{1/2} \), its motion becomes non-adiabatic with the \( E_z \) field causing systematic acceleration in the (+z) directions or \(-z\) motions become if \( n_y \geq 0 \) (Alfvén 1968). The particles are magnetically confined in the \( y \)-direction. Here, \( m_i \) is the particle mass, \( v_{ij} \) is the mean thermal speed in the \( (y, z) \) plane, and \( q_i \) is the particle charge. In a self-consistent model, the non-adiabatic motion of the charges in the \( \pm z \)-directions produces the neutral layer current density \( J_y (y) \). Thus, the magnitude of the surface current density \( J_y (x, y) \) increases as the \( z \)-length of the system, \( L_z \), increases for a given inflow speed. Consequently, there is an upper bound on \( L_z \), or on the energy which can be gained by an electron or ion (Vasyliunas 1980). In terms of the electron Lorentz factor, \( \gamma_e \), this limit is

\[
\gamma_e - 1 = \int_{E_{min}}^{E_{max}} dE_e/m_e c^2 \leq \frac{B_x^2}{4 \pi \rho c^2} \left( \frac{m_e}{m_n} \right),
\]

where \( \rho \) is the plasma mass density and \( e \) is the magnitude of the electron charge. As a numerical illustration, for an ion speed of \( v_i = 0.1c \) and a magnetic field of \( B_z = 0.3 \text{ mG} \), we have \( E_x \approx 10^{-3} \text{ Vcm}^{-1} \) so that an electron is accelerated to \( \gamma_e \approx 2000 \) over the short distance \( L_z \approx 10^{12} \text{ cm} \). This acceleration is compatible with inequality (24) if the Alfvén speed, \( v_A = (B_x^2/4\pi \rho)^{1/2} \), is larger than \( c \).

The one-dimensional Alfvén (1968) model of driven collisionless reconnection is over simplified in the respect that the current-density is independent of \( x \). Furthermore, the model is inconsistent if inequality (24) is violated. Thus, we assume that the current density varies with \( x \) so that there are \( 0 \) and \( X \) points along the \( x \)-axis. This plasma problem is two dimensional and an analogy can be made with the classic MHD reconnection problem treated by Sweet, Parker, and others. As in the MHD reconnection problem, most of the volume of the plasma can be treated using ideal MHD. We replace the small volume where resistivity is important in the classic problem with the small region in which

\[
\psi_d (x, y) = \frac{B_y}{2a} (y^2 - x^2) - \frac{B_y}{12a^3} (x^4 - 6x^2 y^2 + y^4) + \ldots,
\]

where \( B_y \) is the typical field strength and \( a \) is an “outer” length scale. With the current sheet included, the flux function is \( \psi = \psi_a + \psi_{\text{sheet}} \). For \( d \lesssim a \) and \( \delta_j \ll d \), the field in the vicinity of the current sheet is

\[
B_d (x, y, z) \sim \frac{2\pi c}{H (y)} J_d (x),
\]

where

\[
B_d (x, y = 0) = \frac{B_x}{a} + \frac{B_y}{c} x^2 + \frac{2}{3} \frac{p}{c} \int_{-d}^{d} dx' J_d (x').
\]

Fig. 5. Geometry of driven collisionless reconnection. The high energy particles are expelled from the current sheet at \( x = \pm d \). The thick dashed lines crossing the \( x \)-axis beyond the current sheet indicate possible standing shock waves.

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electron or ion in the sheet are:

\[
\frac{d}{dt} (m_{j} v_{zj}) = q_{j} [E_{xj} - v_{xj} B_{yj} / c], \tag{27a}
\]

\[
\frac{d}{dt} (m_{j} v_{yj}) = q_{j} [E_{zj} + v_{zj} B_{yj} / c], \tag{27b}
\]

\[
\frac{d}{dt} (m_{j} v_{zj}) = q_{j} [E_{xj} - (v_{xj} B_{yj} - v_{xj} B_{yj}) / c], \tag{27c}
\]

where \( \gamma_{j} \) is the particle Lorentz factor. The field components \( B_{x}, B_{y}, \) and \( E_{x} \) have been discussed previously. The \( (E_{x}, E_{z}) \) field is the weak ambipolar electrostatic field which arises from the condition of net charge neutrality of the current sheet. For reconnection in an electron/positron plasma we would have \( E_{x} = 0 \).

The particle motion in the \( y \)- and \( z \)-directions is similar to that in the Alfvén (1968) model. The \( y \)-motion is magnetically confined to the current sheet by the \( B_{y} \) field. The motion in the \( \pm z \)-directions undergoes a steady acceleration due to the \( E_{z} \) field. In the \( x \)-direction the particles are expelled from the sheet \((|x| > d)\) by the \( v_{xj} B_{yj} \) force. For an electron/proton plasma an electron is expelled much more rapidly than an ion. This asymmetry gives rise to the above mentioned ambipolar field \( E_{x} \). A self consistent treatment of \( E_{x} \) is beyond the scope of the present work. However, an approximate treatment can be obtained by requiring that \( E_{x} \) be such as to make the \( x \)-velocity of an ion equal to that of an electron. Assuming in addition non-relativistic ion motion, we obtain a modified \( x \)-equation of motion for the electrons,

\[
\frac{d}{dt} [(m_{i} \gamma_{i} + m_{e}) v_{xe}] = -q_{e} v_{xe} B_{y} / c. \tag{28}
\]

Figure 6 shows sample electron orbits in the current sheet. For particles starting from the same initial \( x \), the energy gained by an ion is less than that gained by an electron for \( \gamma_{e} < m_{i} / m_{e} \).

An electron which enters the current sheet at, say, \( x_{i} \), is found, by numerical integration of Eqs. (27) and (28), to exit from the sheet with a final energy \( E_{e} = m_{e} c^{2} \gamma_{e}(x_{o}) \). The inflowing plasma is uniform so that the number of electrons \( dN_{e} \) accelerated to Lorentz factors between \( \gamma_{e} \) and \( \gamma_{e} + d\gamma_{e} \) is simply \( dN_{e} \propto d\gamma_{e} \). Thus, the electron energy distributions is \( dN_{e} / d\gamma_{e} \propto (d\gamma_{e} / d\gamma_{e})^{-1} \). Figure 7 shows the electron energy distributions for an electron/proton plasma and for an electron/positron plasma. The approximate power-law form of the distributions depends on the axial length of the reconnection region, \( L_{x} \), being sufficiently long, \( E(\gamma_{e}) \max (m_{e} c^{2} \gamma_{e}) \). However, the distributions are not significantly dependent on other quantities such as the neutral layer thickness, \( E_{n} / B_{n}, B_{i} / B_{n} \), etc. Of course, the results of our model calculations need to be confirmed and studied further by self-consistent computer simulations.

As noted previously (for example, Galeev et al. 1978; Burkhart et al. 1990, 1991), the time a particle remains in the current sheet is finite, and this has the role of an effective collision time which in turn gives an effective resistivity. Under some conditions the particle motion in the neutral layer is stochastic (Chen & Palmadesso 1986; Buchner & Zelenyi 1986). This may result in a stochastic collisionless conductivity (Horton & Tajima 1990) and give rise to power-law energy spectra (Mima et al. 1991).

5. Comparison of model with observations

The derived jet model has a predominantly toroidal magnetic field and a cylindrical neutral layer where reconnection and stretching of the magnetic field occurs. Here, we compare this model with observations of the polarization and intensity patterns of jets.

Jets can be roughly divided into two main groups depending on the total power of radio galaxy (Fanaorff & Rilely 1974; Bridle & Perley 1984). Those of them with total 1.4 GHz spectral powers exceeding \( 10^{25} \) W/Hz are generally one-sided, of linear morphology, with the edge-brightened lobes and show hot spots at the extremities (FR type II). Well-resolved FR II jets reveal a limb-brightening all along the jet (CBA 1986; OHC 1989). The weaker sources are usually two-sided, more diffuse, more distorted, without hot spots, predominantly center-brightened (FR

![Fig. 6. Representative electron orbits in the plane of the current sheet. The electrons drift into the sheet from the \( \pm y \)-directions, are accelerated by the \( E_{z} \) field, and are subsequently expelled from the sheet, \(|x| > d\)](image)

![Fig. 7. Electron energy distributions for driven collisionless reconnection in an electron/proton plasma (1) and in an electron/positron plasma (2) obtained from Eqs. (27) and (28). The electron energy is \( E_{e} = m_{e} c^{2} \gamma_{e} \), where \( \gamma_{e} \) is the Lorentz factor](image)
type I). The distribution of apparent magnetic fields \( B_\parallel \) also depends on the spectral power (Bridle 1986):

(i) In the strong sources (FR II) longitudinal magnetic field predominates all along the jet; sometimes it becomes transverse in the regions of bright knots. The polarization \( p \) is usually a minimum at the center of the jet and may constitute \( >30-40% \) near the edge. If jet is bent, the polarization in the outer region of the bend is usually higher and can be more than 50%.

(ii) In the weak sources (FR I), the magnetic field is usually transverse to the jet along the jet, excluding the base region, where the field may be longitudinal. Often, \( B_\perp \) is longitudinal near one or both jet edges, or in the outer region of a bent jet. Polarization of the transverse field region is usually a maximum, \( 60-70\% \), at the center of the jet. However, polarization of the longitudinal field regions at the edges of the jet is usually less, \( \sim 10\% \).

The two types of jets may be explained in terms of the proposed model by the different possible values of the ratio of the kinetic energy-density and the toroidal field-energy-density of the jet flow near the neutral layer:

(i) If the flow kinetic energy is large relative to the toroidal field energy, \( \beta_\parallel \gtrsim 1 \) in Eq. (22), then the magnetic field of the jet will consist of two components: a strong axial magnetic field, in a thin annular cylindrical volume layer, and weak toroidal magnetic field filling the remaining volume of the jet. Electrons, accelerated by driven reconnection at the neutral layer will spread along the longitudinal magnetic field and to some extent radially across the jet to the region of weak toroidal field. If the radial diffusion is \( \ll \Delta r \) in the radiative lifetime, then the apparent magnetic field direction deduced from synchrotron radiation will be axial. The synchrotron emission will be from an annular cylindrical volume, and thus limb-brightening is expected. This limit is pertinent to the straight jets in strong sources which exhibit predominantly longitudinal magnetic field (FR II). This limit is compatible with the flows being supersonic with respect to the fast magnetoosonic speed.

(ii) In the opposite limit, the kinetic energy is small relative to the toroidal field energy and \( \beta_\perp \ll 1 \). In this limit, the magnetic field is toroidal throughout the volume of the jet, including the neutral layer. Electrons, accelerated by driven reconnection of the toroidal magnetic field, spread along the toroidal field lines, and they may diffuse radially to the central parts of the jet. If the time of radial diffusion is less than the radiative lifetime, then jet will be observed as center-brightened. In the opposite case it will be observed as limb-brightened. We suggest that the \( \beta_\parallel \ll 1 \) limit may correspond to the FR I type jets which exhibit predominantly transverse magnetic field and high polarization, approaching sometimes the theoretically possible value of 70%. This limit is compatible with flows being sonic or subsonic with respect to the fast magnetoosonic speed.

Some FR type I jets have a base region with longitudinal magnetic field. In these sources, the initial flow energy of the jet may be sufficiently high compared with toroidal magnetic field energy that stretching and formation of a longitudinal component of the magnetic field occurs in the beginning of the jet. However, in the outer regions of the jet \( \beta_\parallel \ll 1 \) and the magnetic field remains toroidal. Some FR type I jets have regions of longitudinal magnetic field at the edges of the jet. This may be explained as a small shearing of the toroidal magnetic field near the neutral layer.

Some FR type II jets exhibit bright knots with transverse magnetic field (Perley et al. 1984). These features have been interpreted as shocks in the supersonic jet flow (see, for example, Begelman et al. 1984). However, recent high resolution observations of nearby FR II jets (CBF 1986; OHC 1989) have shown that these knots look like filaments, or parts of helical structures. Thus, an explanation of the knots as shocks, at least in the M 87 jet, seems less probable than an explanation in terms of instabilities. For example, the Kelvin–Helmholtz instability in the magnetized shear layer, discussed in the Sect. 3.4, can lead to formation of helical magnetic field structures in the shear layer of the jet. Reconnection may be enhanced in these structures giving an increased brightness.

In both FR I and FR II jets, bends in the jet are characterized by enhanced brightness and longitudinal magnetic field in the outer part of the bends. The increase of polarization at the outer region of the bend may be explained by the larger velocity gradient and the larger stretching of magnetic field lines in the shear layer. In addition, the centrifugal force acting on the jet matter in the outer part of the bend can act to enhance the driven reconnection on this side of the jet. Reconnection may also be enhanced by secondary poloidal flow, which results from the bend.

Most jets have a power-law radio spectrum \( S_\nu \propto \nu^{-x} \) with spectral index \( x \sim 0.5-0.7 \). Often, the spectral index is the same all along the jet. This suggests that the electrons are accelerated by the same process equally at different distances along the jet. However, in the optical and X-ray bands the spectrum is steeper, \( \alpha_\text{ox} \sim 1.1-1.5 \) (Giacconi et al. 1978; Harris & Stern 1989) and \( \alpha_\text{ox} \) appears to increase with distance along the jet (Biretta et al. 1991). In the Cen A jet, the spatial distribution of spectral index \( x \) along and across the jet was obtained in the radio band (CBF 1986). In the bright features the spectrum is relatively flat, \( x \sim 0.6-0.7 \), whereas it is much more steeper in the low-brightness regions between them, \( x \sim 1.6 \). Across a bright feature, the spectral index is lowest at the edge of the jet, whereas the brightness is a maximum. This indicates that acceleration of electrons occurs at the edge of the jet in Cen A and that the time-scale for synchrotron losses is less than the time of radial motion of electrons across the jet. This is consistent with the present model where electrons are accelerated via driven reconnection near the neutral layer at the jet's edge.

6. Summary

1. This work has developed a model of non-relativistic magnetized jets with the magnetic field dragged out from the rotating disk of the compact object by the jet flow. The outflowing plasma has two regions of oppositely directed magnetic field, twisted by the disk rotation. Thus, the jet has zero net poloidal magnetic flux and zero net poloidal current. At large distances the toroidal magnetic field dominates, and the jet magnetic field consists of an inner region (small \( r \)) with, say, \( B_\parallel < 0 \), and an outer region with \( B_\parallel > 0 \), with a cylindrical neutral layer with \( B = 0 \) separating the regions.

2. Magnetic reconnection at the neutral layer may occur due to instability or it may be driven by the secondary jet flow. Driven collisionless reconnection at the neutral layer may lead to the acceleration of electrons to relativistic energies. This mechanism converts magnetic field and jet flow energy to energy of relativistic electrons. A simple model is developed for particle
acceleration in electron/position and electron/proton plasmas which predicts power-law energy distributions.

3. Velocity shearing of the neutral layer is taken into account. Magnetic inhomogeneities (for example, magnetic islands), formed during the reconnection, are stretched by the jet flow and form elongated loops in the axial direction. This process leads to formation of axial component of the magnetic field in the boundary layer of the jet. If stretching energy is comparable with the energy of toroidal magnetic field, then electrons, accelerated at sites of reconnection, will move along the stretched magnetic lines and the apparent magnetic field will be longitudinal. We suggest, that this limit correspond to jets in powerful sources (FR type II). In the opposite limit, stretching is unimportant and the toroidal magnetic field dominates throughout the volume of the jet. This limit may correspond to weak sources with mainly transverse magnetic field (FR type I), although there may be field stretching near the bases of these jets.

4. Different features such as bright knots in FR type II jets, bends in different types of jets, and many features such as helical magnetic field structures of the well-resolved jet in M 87 may be explained by superposition of reconnection with velocity shearing in a boundary region of the neutral layer.

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