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Dynamo model of double radio sources

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A model of radio sources is proposed in which the magnetised accretion disk of a massive black hole acts as an electric dynamo producing two oppositely-directed beams of ultra-relativistic particles.

WE give here a brief development of a model of double radio sources with the conducting, magnetised accretion disk of a massive black hole acting as an electric dynamo. The electric field created by the rotating disk results in the steady generation of two oppositely-directed, collimated, tenuous beams of ultra-relativistic protons. The output power in the beams is $\sim 10^{44}$ erg s^{-1} and the proton energies are $\lesssim 10^{19}$ eV for a poloidal magnetic field of the disk of 10^3 gauss and a black-hole mass of $10^9 M_\odot$. The beams propagate parallel to the angular momentum vector of the disk. The energy transported by the proton beams is gradually degraded because instabilities lead to the sharing of the beam energy with the ambient plasma. It seems possible that the two-stream interaction can convert the energy of a tenuous proton beam into a relatively dense, collimated beam of relativistic electrons. Energy carried by the electron beams is released as synchrotron radiation in the radio components.

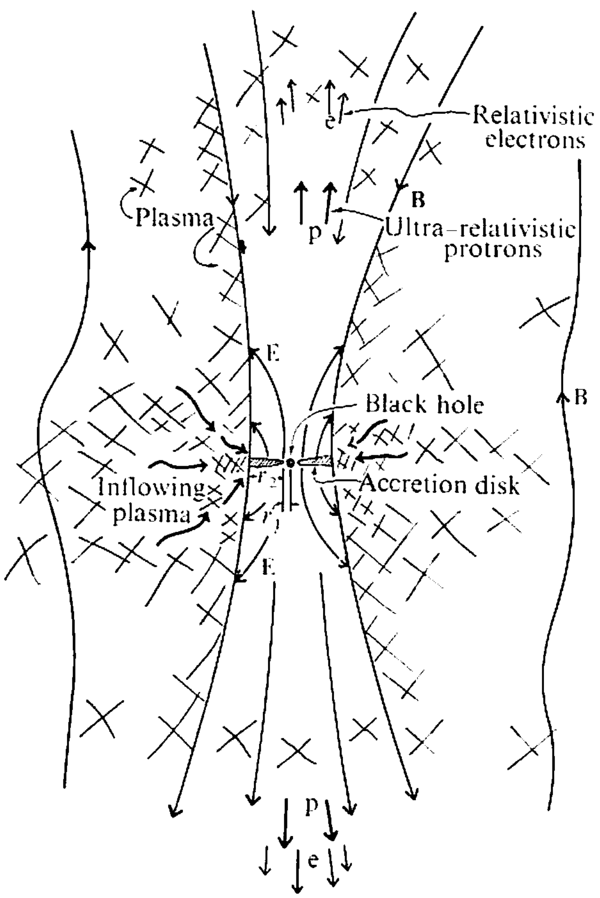

Fig. 1 Schematic drawing of the electric dynamo formed by a conducting magnetised accretion disk. The radio wave emitting regions are discussed in the text.

With this model it seems possible (1) to have a large, steady supply of energy to widely separated small radio components, (2) to obtain alignment and symmetry of the two radio components, and (3) to have a correlation between the axis of the radio components and the direction of the angular momentum of the parent galaxy. The model includes some of the features of the 'twin exhaust' model[1], the spinar or oblique-rotator models[2–5], the magneto-hydrodynamic models[6,7] and the axisymmetric pulsar model[8]. The isotropic production of low energy ($\sim 10^{12}$ eV) cosmic rays from the accretion disk of a massive black hole has been discussed previously[9,10].

Model

Consider the possibility[9,10] that a black hole of mass M_h ($\sim 10^9 M_\odot$) has formed in the centre of a galaxy (or in the nucleus of a quasar) and that the black hole is surrounded by a massive, flat accretion disk (see Fig. 1). Gaseous material accreted on to the disk is supplied from the mass loss of stars, possibly increased by disruptive collisions between stars and/or the close approach of stars to the black hole[11]. It is evident that the angular momentum vector of the disk will have a direct correlation with the angular momentum vector of the parent galaxy. If the mass of the disk is sufficiently large, $M_d \gg M_\odot$ (but $M_d \ll M_h$), then the angular momentum of the disk may have a fixed direction in space for long periods, greater than the lifetime of a radio source. For $M_d \gg M_\odot$, the disk angulai momentum corresponds to an average of the angular momenta of a large number of stars.

The magnetic field

Ionised matter moving into the accretion disk will carry with it a weak ambient magnetic field. The gradual inward radial motion of matter in the disk will amplify the field, tending to establish $B_z(r) \propto \sigma(r)$, where σ is the surface density of the disk, r is the radial distance from the black hole, and the z axis is normal to the plane of the disk. We consider that B_z at some particular epoch has a single polarity throughout the inner part of the disk. Because of differential rotation of the disk, the poloidal field (B_r, 0, B_z) is expected to be axisymmetric. It is, however, possible that the poloidal field is augmented by non-axisymmetric turbulent fluid motions[12] in the disk. Reconnection of the poloidal field lines occurs in the region $r < r_1, |z| < r_1$ (where r_1 is the inner radius of the disk) so that there is no build up of flux through the surface $r < r_1, z = 0$. (For the moment, consider that the space above and below the disk is a vacuum. Thus, although there may be a toroidal field (B_φ) embedded in the disk[9,10], $B_\varphi=0$ in the space above and below the disk.)

The electric field

The material of the disk may be assumed highly conducting so that an observer in a non-rotating (inertial) reference frame sees a strong radial electric field, $E_r(r) = -[u_\varphi(r)/c]B_z(r)$, on the top and bottom surfaces of the disk, where $u_\varphi = (GM_h/r)^{\frac{1}{2}}$ is the Keplerian velocity, assumed much larger than u_r. Thus the potential difference across the disk is

$$V_{12} = -\frac{1}{c}\int_{r_1}^{r_2} \mathrm{d}r u_\varphi(r) B_z(r) \qquad (1)$$

where $r^1 = 6GM_h/c^2 \simeq 10^{14}$ cm $(M_h/10^8 M_\odot)$ is the inner radius of the disk, and $r_2 \gg r_1$ is the effective outer radius where material enters the disk. If we assume $u_r(r) = \varepsilon u_\varphi(r)$, with ε = a constant, $\ll 1$, then $\sigma \propto r^{-\frac{1}{2}}$. Because $B_z(r) \propto \sigma(r)$, we find

$$V_{12} \simeq -\sqrt{(6)}B_*(GM_h/c^2)\ln(r_2/r_1) \qquad (2a)$$

where $B_* = \max(B_z) = B_z(r_1)$. In practical units

$$V_{12} \sim 10^{19}\left(\frac{B_*}{10^3\text{ gauss}}\right)\left(\frac{M_h}{10^8 M_\odot}\right)\ln\left(\frac{r_2}{r_1}\right)\text{V} \qquad (2b)$$

Observe that V_{12} is proportional to r_1, which is well defined, but depends only weakly on r_2.

Magnetic flux and time scales

The magnetic flux through the disk is

$$\Phi \sim 4\times10^{31}\left(\frac{B_*}{10^3\text{ gauss}}\right)\left(\frac{M_h}{10^8 M_\odot}\right)^2\left(\frac{r_2}{r_1}\right)^{3/2}\text{gauss cm}^2 \qquad (3)$$

For $B_* = 10^3$ gauss, $M_h = 10^8 M_\odot$, and $r_2/r_1 = 10$, this flux ($\sim 10^{33}$ gauss cm^2) corresponds to an ordered field of 10^{-6} gauss spread over a region of $\sim$ 10 pc. A characteristic time scale for rapid dynamical variations of the disk is the orbital period at $r = r_1$, $T_d \simeq 5\times10^4(M_h/10^8 M_\odot)$ s. At the other extreme, the time scale for significant changes in the large scale magnetic field structure, once it is established, is $T_B = \mathscr{L}/R$, where $\mathscr{L}$ is the inductance and R the resistance of the configuration. We estimate that T_B is at least as long as the lifetime of the disk because of the large size and high conductivity of the disk.

Resulting particle acceleration

Consider the particle acceleration resulting from the disk potential. In the space outside the cylinder $r > r_2$, there is assumed to be plasma, possibly in turbulent motion, but not in rapid rotation. The electrical potential of this plasma may be taken to be zero. The region $r < r_2$ above and below the disk is for the moment assumed a vacuum. We consider in the present study only the case where B_z is anti-parallel to the angular momentum vector of the disk. Thus the potential on the disk surfaces, $V(r,\pm0)$, produces a vacuum field in the space above and below the disk. That is, $\nabla^2 V = 0$, with $V(r,\pm0)$ specified, with $V(r_2, z) = 0$, and with $V(r, |z| \gg r_2) = 0$. The potential $V(r, z)$ implies an electric field which tends to accelerate protons (and other ions) in the $+z$ $(-z)$ direction out of the top (bottom) surface of the disk. The electric field, E, is $\sim (u_\varphi/c)Bz \lesssim 10^5$ V cm^{-1} for $B_* = 10^3$ gauss and $M_h = 10^8 M_\odot$. The electric force on a proton exceeds the gravitational attraction of the black hole by a large factor, $\sim 10^{12}$. The electric field is, however, small compared with the poloidal magnetic field, $E^2/B_z^2 \sim u_\varphi^2/c^2 \ll 1$.

The self-consistent, space-charge-limited flow of protons off the disk surface follows from general dimensional considerations (ref. 13 and E. Ott, T. Antonsen, and R. Lovelace, unpublished) in that the particle motion is extremely relativistic. If I denotes the total current of protons in, say, the $+z$ direction, then

$$I = \kappa c V_{12} \qquad (4a)$$

where κ is a dimensional factor which depends on the geometry through $V(r,\pm0)$. We may assume that κ is not greatly different from unity (ref. 13 and E. Ott, T. Antonsen and R. Lovelace, unpublished). An estimate analogous to equation (4a) occurs in the axisymmetric pulsar model[8]. Using equation (2b), we rewrite equation (4a) as

$$I \sim 3\times10^{17}\kappa\left(\frac{B_*}{10^3\text{ gauss}}\right)\left(\frac{M_h}{10^8 M_\odot}\right)\ln\left(\frac{r_2}{r_1}\right)\text{A} \qquad (4b)$$

Charge neutrality of the disk is maintained by a non-relativistic electron current in the region $r > r_2$. For $r < r_2$ and $|z| > 0$, there is a positive space-charge density J_z/c, where $J_z(r)$ is the beam current density.

The output power in the two beams is, from equations (2b) and (4b)

$$L_R \simeq 2IV_{12} \sim 10^{44}\kappa\left(\frac{B_*}{10^3\text{ gauss}}\right)^2\left(\frac{M_h}{10^8 M_\odot}\right)^2\left[\ln\left(\frac{r_2}{r_1}\right)\right]^2\text{erg s}^{-1} \qquad (5)$$

The power in the beams derives from the infall of matter into the gravitational potential of the black hole. There is a corresponding flow of angular momentum[5] out of the disk carried by the poloidal-magnetic field and self-electric field of the beam and the beam protons. The beam emission may, however, represent only a small dynamical influence on the disk if $L_R \ll L_{tot}$, where L_{tot} is the total luminosity of the disk. In particular, an irregular magnetic field embedded in the disk may be dynamically important[9,10,14]. For comparison, the Eddington limit on the electromagnetic emissions of the disk is $L_E \sim 10^{46}(M_h/10^8 M_\odot)$ erg s^{-1}.

Evidence in support of equations (4) and (5) comes from many laboratory experiments which have demonstrated the production of high-current electron beams[15,16] ($I_e \sim 10^5$ A) and ion beams (P. Dreike, C. Eichenberger, S. Humphries, and R. Sudan, Cornell University Laboratory of Plasma Studies Report 172 (1975)) ($I_p \sim 10^4$ A) in this type of field configuration[15-18] ($E^2/B_z^2 < 1$) at low voltages ($V_{12} \lesssim 10^6$ V) with electric fields $E \lesssim 10^6$ V cm^{-1}. The observed beam particle energies are $\sim$ eV_{12}.

Beam energy/magnetic energy

An important dimensionless parameter for the dynamics of the beams is the ratio, denoted β, of the kinetic energy density of the beam to the energy density of the magnetic field. We first estimate $\beta = \beta_0$ near the disk, $|z| \lesssim r_2$. Note that we have a limit on the proton Lorentz factors γ from equation (2b)

$$\gamma \leqslant \gamma_0 \sim 10^{10}\left(\frac{B_*}{10^3\text{ gauss}}\right)\left(\frac{M_h}{10^8 M_\odot}\right)\ln\left(\frac{r_2}{r_1}\right) \qquad (6)$$

The spectrum of γ values as a function of r is expected to be relatively flat for $\gamma < \gamma_0$. An estimate of the beam density, n_B, follows from equation (4b)

$$n_B \sim 2\times10^{-3}\kappa\left(\frac{r_1}{r_2}\right)^2\ln\left(\frac{r_2}{r_1}\right)\left(\frac{B_*}{10^3\text{ gauss}}\right)\left(\frac{10^8 M_\odot}{M_h}\right)\text{cm}^{-3} \qquad (7)$$

The total number of protons emitted per second is thus $\sim 4\times10^{36}$ for $n_B = 2\times10^{-5}$ cm^{-3}, $r_1 = 10^{14}$ cm, and $r_2/r_1 = 10$.

From equations (6) and (7), we find

$$\beta_0 \quad 8\pi\frac{n_B\gamma_0 m_p c^2}{B_z^2} \sim 0.7\kappa\left[\frac{r_1}{r_2}\ln\left(\frac{r_2}{r_1}\right)\right]^2 \qquad (8)$$

where m_p is the proton rest mass. Evidently, we have $\beta_0 \ll 1$ if $r_2 \gg r_1$. That $\beta_0 \ll 1$, is a result of the fact that $u_\phi^2/c^2 \ll 1$ over most of the surface area of the disk. Thus for the proton acceleration and streaming in the vicinity of the disk, the magnetic field B_z is very strong. To a good approximation the protons follow the magnetic field lines and the configuration is force-free. (Values of β small compared with unity are known to be favourable to the magnetohydrodynamic stability of beams electrically neutralised by plasma[20].)

In the vicinity of the disk there is a small toroidal self-field B_ϕ from the proton beams. We estimate that

$$\frac{B_\phi}{B_z} \simeq \frac{2I}{cr_2B_z} \sim (\beta_0\kappa)^{1/2}$$

is small compared with unity if $\beta_0 \ll 1$. This toroidal magnetic field counteracts the radial space-charge repulsion of the beam and thus tends to collimate the beam.

Divergence of beam

For distances $|z| \gg r_2$ above and below the disk, the proton beams diverge. The outer beam radius is approximated as $a(z) = r_2(1+|z|/H)$, where H is the scale height for the divergence. We assume that processes involved in the initial outward propagation of the beams give $H \gg r_2$. Because the beam particles start off following the field lines, $B_z(z) \propto [a(z)]^{-2}$. The β factor of the beam is proportional to a^2, and, therefore, we have $\beta \simeq 1$ at a distance $|z| \simeq z_* = H(\beta_0)^{-\frac{1}{2}} \gg H \gg r_2$, where β_0 is given in equation (8). For distances $|z| \gg z_*$, where $\beta \gg 1$, the beam divergence comes from the ballistic motion of the protons (in the absence of instabilities). The usual adiabatic energy losses, involved, for example, in the expansion of a cloud of particles, do not occur for a particle beam in ballistic motion.

Radiation effects

In the initial outward propagation of the beam there may be significant curvature radiation. If ρ denotes the instantaneous radius of curvature of a proton trajectory, then the typical energy of the curvature photon is

$$\left(\frac{3}{2}\right)\gamma^3\left(\frac{hc}{\rho}\right) \sim 3\times10^{10}\left(\frac{\gamma}{10^{10}}\right)^3\left(\frac{10^{15}\text{cm}}{\rho}\right)\text{eV}$$

The rate at which a proton emits energy is

$$\frac{d\gamma}{dz} = \frac{2r_e}{3}\left(\frac{m_e}{m_p}\right)\frac{\gamma^4}{\rho^2}$$

where r_e is the classical electron radius, and m_e is the electron rest mass. Thus the path length for a proton to lose half its initial energy ($\gamma m_p c^2$) is

$$l_{cur} \sim 2\times10^{16}\left(\frac{10^{10}}{\gamma}\right)^3\left(\frac{\bar\rho}{10^{15}\text{ cm}}\right)^2 \text{cm}$$

where $\bar\rho$ is an average of the radius of curvature. In the initial stages of formation of a source we may have $l_{cur} \lesssim \bar\rho \simeq r_2$ if $B_* M_h \gtrsim 10^{11}$ gauss $M_\odot$. One effect of the curvature photons may be that of expelling matter and thus straightening the magnetic field in two conical regions of small solid angle defined by the velocity distribution of the protons at $|z| < r_2$. The effect is, however, inefficient because the production and scattering cross sections for the curvature photons interacting with electrons and protons are small ($\lesssim 2\times10^{-26}$ cm^2). The curvature photon production of electron–positron pairs off the magnetic field[21] is negligible.

The beam protons may interact with low energy photons[22] emitted by the accretion disk[9,10] and surrounding plasma to give pions and electron–positron pairs. The proton energy threshold for pion production is

$$\sim 10^{17}(1-\cos\Theta)^{-1}\left(\frac{1\text{ eV}}{\mathscr{E}}\right)\text{eV}$$

where Θ is the angle between the proton and photon momenta, and $\mathscr{E}$ is the average photon energy. The threshold for e^-/e^+ production is a factor 1/135 smaller. We estimate that the proton energy loss is small if

$$L_{ph} < 10^{46}\left(\frac{\mathscr{E}}{1\text{ eV}}\right)\left(\frac{R}{10^{17}\text{ cm}}\right)\text{erg s}^{-1}$$

where L_{ph} is the photon luminosity, and R is the effective radius of the photon-emitting region.

Counter flow instability

For distances $|z|$ larger than a certain value z_{neut} (discussed below), the proton beams become both charge and current neutralised by plasma in the beam channels. In the outward propagation of a beam, plasma enters the beam front, and, furthermore, plasma may leak radially into the beam channel because of instabilities (particularly for $|z| > z_*$, where $\beta > 1$). Thus the presence of free electrons on magnetic field lines connecting with the accretion disk presents the possibility of electrons counter-flowing and partially 'shorting-out' the proton accelerating potential. This shorting would result in an ohmic dissipation (and radiation) of a fraction (~1/2) of the power $2IV_{12}$ in a small region of size $\lesssim r_2$ centred on the black hole with the remaining fraction of the power in the two proton beams. There is, however, a definite asymmetry between the proton acceleration and the counter acceleration of electrons.

Consider the two-stream interaction: let n_p denote the density of cold electrons (speeds $\ll c$) within the beam channel at some distance z. Then the dispersion relation for electrostatic waves with frequency ω and wavevector **k** is

$$0 = 1-\frac{\omega_p^2}{\omega^2}-\frac{\omega_B^2\psi^2}{(\omega-kv_B\cos\psi)^2} \qquad (9)$$

where **k** is at an angle ψ to the beam velocity $\mathbf{v}_B$, with $\gamma^{-1} \ll \psi \ll 1$, and where $\omega_p = (4\pi e^2 n_p/m_e)^{\frac{1}{2}}$, $\omega_B = (4\pi e^2 n_B/\gamma m_p)^{\frac{1}{2}}$, $k = |\mathbf{k}|$, and $v_B = |\mathbf{v}_B| = c(1-1/2\gamma^2)$ is the local proton speed. It is evident from equation (9) that the proton beam may be considered weak in the limit where $n_p \gg (m_e/\gamma m_p)n_B\psi^2$. In this limit a simple form of the two-stream instability occurs[23]. The frequency of the unstable wave is written as $\omega = \omega_p + \delta$, with $|\delta| \ll \omega_p$, and $k = \omega_p(v_B\cos\psi)^{-1}$ so that

$$\text{Im}(\delta) = \frac{\sqrt{3}}{2}\left(\frac{\omega_B^2\omega_p}{2}\right)^{1/3}(\psi)^{2/3}$$

or

$$\text{Im}(\delta) \sim 3\times10^{-5}\left(\frac{10^{10}}{\gamma}\right)^{1/3}\left(\frac{pc}{a}\right)^{2/3}\left(\frac{n_p}{10^{-3}\text{ cm}^{-3}}\right)^{1/6}\psi^{2/3}\text{ s}^{-1} \qquad (10)$$

where we have taken

$$n_B = 10^{-12}\left(\frac{pc}{a(z)}\right)^2\text{cm}^{-3}$$

with $a(z)$ the beam radius. The characteristic wavelength of the unstable wave is

$$\lambda \simeq \frac{2\pi c}{\omega_p} \sim 10^8 \text{ cm} \left(\frac{10^{-3} \text{ cm}^{-3}}{n_p} \right)^{1/2}$$

The phase velocity of the unstable wave

$$v_\phi = c \left[1 - \frac{\psi^2}{2} - \tfrac{1}{2} \left(\frac{\omega_B{}^2 \psi^2}{2\omega_p{}^2} \right)^{1/3} \right] < v_B \qquad (11)$$

is quite close to the speed of light. As an illustration of equation (10), for $a = 10^{-2}$ pc ($=3\times10^{16}$ cm), $\gamma = 10^{10}$ and $n_p = 10^{-3}$ cm^{-3}, an e-folding of a wave with $\psi = 10^{-3}$ occurs over a distance of $\sim 10^{-3}$ pc ($=3\times10^{15}$ cm). Thus an unstable wave may grow considerably before it convects out of the beam channel.

Because of the two-stream instability, plasma waves of finite but small amplitude are expected to coexist with the plasma in the beam channel[24,25]. One important effect of the waves is to produce an anomalous electrical resistivity (which does not depend on binary collisions). This resistivity acts to prevent electrons from partially shorting the proton accelerating potential. A self-consistent, non-linear theory of the interaction would determine the neutralisation length $z_{\rm neut}$. It is of course possible that the nonlinear theory would allow only a quasi-periodic acceleration[21] of protons with a time scale $\sim z_{\rm neut}/c$.

Linear instability of the proton beams may also involve the circularly polarised Alfvén and/or cyclotron plasma waves propagating parallel to the magnetic field. The interaction of streaming cosmic rays with Alfvén waves has been studied[26,27], but apparently only under conditions where the cosmic rays are nearly isotropic.

For distances $|z| > z_{\rm neut}$, the magnetic field is the 'stretched-out' poloidal field of the accretion disk, and is nearly parallel to the local proton velocity. For example, if the magnetic flux through the disk is 10^{33} gauss cm^2 (from equation (3)) then

$$B_z \sim 10^{-4} \left[\frac{pc}{a(z)} \right]^2 \text{ gauss} \qquad (12)$$

where $a(z)$ is the beam radius. The proton gyro-radii implied by equation (12) are $\simeq 10^2 pc \sin(\theta)\ (\gamma/10^{10})(a/pc)^2$, where $\theta \ll 1$ is the pitch angle.

The electron beams

It is possible that the energy loss per unit distance of the proton beam is determined by the two-stream interaction. A small fraction of the cold electrons in the beam channel get trapped and untrapped in the unstable plasma waves and are thereby accelerated to relativistic energies[24,25]. At the same time the proton spectrum is shifted to lower energies. A nonlinear theory is required to predict the energy spectra of the electrons and protons. It is, however, relevant that equation (11) implies that an electron trapped with the phase velocity of an unstable plasma wave has a Lorentz factor

$$\gamma_w = \Gamma(\psi^{2/3} + \Gamma^2\psi^2)^{-1/2}$$

$$\Gamma = 2^{1/6} \left(\frac{\omega_p}{\omega_B} \right)^{1/3} \sim 5\times10^3 \left(\frac{n_p}{10^{-3} \text{ cm}^{-3}} \right)^{1/6} \left(\frac{\gamma}{10^{10}} \right)^{1/6} \left(\frac{a}{pc} \right)^{1/3} \qquad (13)$$

where we have taken $n_B = 10^{-12}(pc/a)^2$ cm. The energetic electrons are collimated with the proton beams, having small pitch angles and small synchrotron losses. An upper limit on the range of electron Lorentz factors may however occur because the pitch angles are not negligibly small. For small pitch angles the synchrotron radiation is collimated with the proton beams.

The radio components

The collimated propagation of the oppositely travelling beams of electrons terminates in two regions of size a_R at $z = \pm z_R$, which are the radio components. In these regions the magnetic field of the beam channel connects with a non-uniform irregular field. Thus relativistic electrons arriving at $\pm z_R$ rapidly lose their energy by incoherent synchrotron emission. The collimated propagation of protons with gyro-radii $\gg a_R$ continues beyond z_R, and these protons may continue to accelerate cold electrons by the two-stream interaction. The observed values of z_R extend from <1 pc to ~ 1 Mpc. Consider first the small sources with, say, $z_R < 10^2$ pc. The magnetic field is that carried by the beam (equation (12)). The momentum of the beam electrons corresponds to a pressure $\lesssim P_R \equiv (L_R/\pi c)\,[a(z)]^{-2}$, where L_R is given by equation (5). The electron pressure compresses the external medium and gives rise to shock wave structures[28,1] at $z = \pm z_R$. These structures may move with relativistic velocities in the $\pm z$ directions if $P_R > n_{\rm ext}\, m_p\, c^2$, or $L_R > 10^{12}\,(a/pc)^2\,(n_{\rm ext}/10^{-3}\text{ cm}^{-3})$ erg s^{-1}, where $n_{\rm ext}$ is the ambient density of the external medium. For the extended sources with $z_R > 10^2$ pc, the location of the shock structures (and associated radio emitting regions) is determined by the pressure balance conditions[26,1] $P_R \gtrsim n_{\rm ext} m_p U_z{}^2$, where U_z ($<c$) is the velocity of the structures in the $\pm z$ directions into the external medium. For the large sources, the magnetic field carried by the beam is exceedingly weak. Amplification of the field may, however, occur by field line stretching at the expense of the energy of the turbulent motions of the plasma and relativistic electrons. Subsequent to the formation of an extended double source, a second pair of oppositely travelling shock wave structures may be generated, for example, by an instability involving the accretion disk. The secondary shock structures will be aligned with the disk angular momentum and thus approximately aligned with the extended double.

The high energy protons which reach the large radio components escape into extragalactic space without appreciable synchrotron losses if the magnetic field strengths are $< 10^{-3}$ gauss. It appears possible that the accumulation over time of the high energy protons lost from nearby ($\lesssim 300$ Mpc) dynamo sources has produced the high-energy end (10^{16}–10^{20} eV) of the observed cosmic ray spectrum.

An analytic solution to the electrodynamics of a magnetised accretion disk has recently come to our attention[29].

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