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MAGNETICALLY ORDERED JETS FROM PULSARS

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ABSTRACT

We propose that a small fraction ($<10^{-2}$) of the magnetic field energy expelled by a pulsar may leave in an ordered, propagating form. These plasma-loaded field lines leave close to the spin axis and form a high-pressure beam. If the beam encounters no significant opposing pressure it will form a cone, limb brightened in synchrotron emission, until it strikes a supernova shell. Filled supernova remnants will slow the beam, winding up the toroidal field and producing a well-collimated, nearly force-free ($\boldsymbol{J} \parallel \boldsymbol{B}$) equilibrium. Loading of external plasma on field lines enhances stability, and return currents to the pulsar are formed naturally through induction. Such a structure is buoyant in the remnant and may bend, though it should straighten as it leaves the remnant. These types of ordered forms may apply to the jets in the Crab nebula and MSH 15-52. We predict that the Crab jet (*a*) will be a polarized synchrotron source, but probably *not* with a hollow profile and (*b*) will have an H/He ratio typical of the nebula.

Subject headings: nebulae: supernova remnants — pulsars

I. INTRODUCTION

Active galactic nuclei may be powered by black holes, delivering energy into jets through electrodynamic processes. A Kerr black hole in a magnetic field shares many properties of a conductor rotating in a magnetic field. Though black hole magnetospheres lack the extra "organization" present in the fixed rotation frequency of a pulsar, they may produce very well ordered jets.

Supposing that some extragalactic jets are in fact driven by a DC Poynting flux (Phinney 1983; Macdonald and Thorne 1982; Znajek 1978; Lovelace 1976), it seems curious that pulsar magnetospheres, which are better ordered in some ways, do not emit well-collimated and obvious jets. To understand black hole magnetospheres, it would seem useful to understand why or *if* pulsars are jetless. This paper proposes that while most pulsar outflowing energy is indeed disordered, there may be circumstances when a small fraction of the outflow can appear as long-lived, ordered magnetic field structures.

We do not treat in detail how a jet will form since the entire pulsar-nebula coupling problem is so fraught with uncertainty. Instead, in § II we discuss qualitatively why such jets might arise and assume thereafter that consideration of the dynamics of magnetic jets can lead to observable consequences, especially for the Crab jet.

II. ORDERED FLOWS

There is controversy over whether the light cylinder radius, r_L, is crucial to the outflow. As Michel's (1982) excellent review makes clear, Holloway (1973) and Ardavan (1976) claim that disordering electrodynamic transitions can occur. Michel and Tucker (1969), Sharlemann (1974), and others feel that plasma currents induced at the radius where corotation fails (which need not be r_L) will smoothly bend magnetic field lines back through the transition. Here we shall not take sides in this controversy, because the details of the transition zone are of secondary importance. The principal fact is that a unidirectional beam must come from near the rotation axis of a pulsar, for otherwise it would disperse its energy over a large volume. Where the jet is formed (r_L or beyond) is not as important as ensuring that the plasma $\beta \sim 1$ at the transition to an ordered flow.

For the moment let us take the view that r_L is the crucial site where disordering occurs. As an outflowing plasma leaves a rotating neutron star, the ratio of kinetic pressure to magnetic pressure, β, is small in the magnetosphere. At the light cylinder, kinetic pressure may overcome magnetic constraint, and a disordered fluid ("wind") can propagate outward. Near the equatorial region, Holloway (1973) showed that the change in sign of $\boldsymbol{\Omega} \cdot \boldsymbol{B}$ implies large electric fields ("vacuum gaps"). (Here $\boldsymbol{\Omega}$ is the pulsar spin vector and $\boldsymbol{B}$ the magnetospheric field.) This surely helps convert the flow from low β to high β near the light cylinder. However, these electric fields probably do not occur at high latitudes; there disruptions of flows can occur because of dynamical effects.

Disordering of the magnetic field occurs in part because the plasma experiences a rapid rise in β in the last few percent of the way to the light cylinder radius ($r \approx r_L$), due to spinup to higher relativistic energies. A field line leaving at angle θ with respect to the angular velocity vector $\boldsymbol{\Omega}$, experiences a change in β along the distance $r = xr_p$ which varies with θ (r_p is the pulsar radius and x a scaling factor). The ratio β^{-1} ($\partial\beta/\partial x$) $\sim \sin\theta$ for small θ, whereas for the last open field line at $\theta = 0.1$ (aligned rotor), this ratio is ~ 1 (Ruderman and Sutherland 1975). Thus for a nearly aligned rotor the transition is disorderly (rapid change of β for most of the field lines leaving the polar cap) because of this dynamical effect and Holloway's (1973) suggested equatorial gap regions. Rapid reconnection can occur in the outgoing wind. This leads naturally to a "hot" wind in which streaming and thermal pressures exceed magnetic pressure; this allows strong reverse shcks, which cannot occur in a cold wind (Kennel and Coroniti 1983). Seemingly, most of the pulsar energy is delivered to the nebula in this way. Rees and Gunn (1973) argued, and Kennel and Coroniti (1983) showed in more detail, that the Crab nebula boundary conditions probably require an outflowing wind which is not on average magnetically dominated.

Suppose the great majority of outflowing field lines experience rapid changes in β at r_L, creating eddies of size $d = fr_L$,

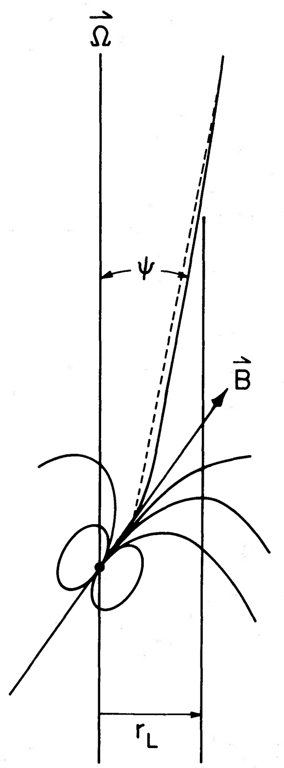

FIG. 1.—Pulsar open field lines emerging from the equatorial region of the light cylinder, r_L, lead to hot plasmas and low magnetic field energy density. Lines emerging close to the spin axis $\boldsymbol{\Omega}$ at angle $\Psi \ll 1$ can be smoothly forced open by corotation forces. This can produce magnetically ordered jetlike equilibria beyond r_L.

with f a fraction. These magnetic structures can reconnect, destroying the stored energy and creating a hot wind, in a magnetic diffusion time $t_r = [\omega_p(r_L) f r_L/c]^2 \nu^{-1}$, where ν is the frequency of electron-positron collisions, and $\omega_p(r_L)$ is the plasma frequency at r_L. For streaming instabilities $\nu \sim g\omega_p$ with $g \leq 0.1$ (Miller 1982). Using this estimate, $\tau_r \sim 3f^2 g^{-1} [n_p(r_L)]^{1/2}$ s. Using even a high value of plasma density $n_p(r_L) \sim 10^8$ cm^{-3}, reconnection occurs within an hour after passing through r_L, for $f = 0.1$, $g = 0.1$. Folding in the spatial dependence $\omega_p \propto r^{-1}$ and letting the eddy size scale linearly with r, then $\tau_r \propto r$, which still implies that most disordered magnetic structures in the bulk of the outflow will disappear close to r_L.

The Crab suggests that orderly flows are not energetically important, but may be necessary to account for the nebular radio radiation. Adding this contribution to the hot wind, which apparently produces the high-energy emission (Kennel and Coroniti 1983), would then give a description of the whole spectrum. Field lines nearly parallel to the spin axis can allow smoother transitions, in β_0 (Fig. 1). Consider a field line crossing the light cylinder radius, r_L, where the magnetic field has a toroidal component, B_t, and a stronger poloidal part, B_p. Using Ruderman-Sutherland parameters (1975), the condition that the relativistic plasma β be unity is (aligned rotor):

$$\beta = 0.1 \frac{P^{13/7} r_{p6}^{4/7}}{B_{12}^{8/7}} \left(1 + \frac{(r/r_L)^2}{\tan^2 \Psi}\right) = 1 \, . \quad (1)$$

Here P is the pulsar period, r_{p6} the pulsar radius in units of 10^6 cm, B_{12} the surface magnetic field in units of 10^{12} gauss, and $\tan \Psi = B_t/B_p$. Field lines of a nearly aligned rotor, making an angle $\tan \psi \approx 0.3P$, should allow a smooth transition from field-dominated, $\beta \ll 1$ flow to $\beta \sim 1$. For the Crab, $\psi \sim 0.01$. This result is dependent on the inclination of the magnetic axis to the spin axis; an orthogonal rotor yields $\psi = 0$. So long as some open field lines are closely aligned with the axis, though, outflow along them can be ordered because $\beta^{-1} \partial\beta/\partial x$ is small. For nearly aligned, longer period pulsars this could include a large fraction of the polar cap surface.

Suppose now that the point where the pulsar flow "comes to terms" with the nebular pressure lies beyond the light cylinder. A jet can form at this larger radius if it is able to make a $\beta \sim 1$ transition along field lines aligned almost along $\boldsymbol{\Omega}$. The only fundamental requirement is that the winding up of $\boldsymbol{B}$ around $\boldsymbol{\Omega}$ be accompanied by particle pressure along that axis. We know from Kennel and Coroniti (1983) that a hot, high-β wind must dominate the outflow into the Crab nebula, so a picture similar to Figure 2 emerges, with a larger radius $r > r_L$ marking the zone where all but a small jet flow is disordered. The angle between $\boldsymbol{\Omega}$ and $\boldsymbol{B}$ seems noncrucial to this discussion, as long as some field lines can wrap around $\boldsymbol{\Omega}$. If such mild transitions occur, what becomes of the outflowing system beyond r_L?

III. HELICAL EQUILIBRIA

As the ordered field-particle system rises away from the polar regions of the neutron star, B_p drops with cylindrical jet radius r_j as r_j^{-2}, while $B_t \propto r_j^{-1}$. Thus the constraining pressure from B_t will become dominant at a distance along the spin axis $L' = r_L/\tan^2 \Psi \sim 5 \times 10^{10}/P$ cm (Fig. 2). Meanwhile, particle pressure of the electron-positron plasma transverse to the field will decay in seconds due to synchrotron radiation. If energy enters the nebula as a whole at the rate $\dot{E}_n$ through a polar cap with angle θ_c, the jet energy outflow $\dot{E}_j \sim \dot{E}_n(\psi/\theta_c)^2$. Thus only a small fraction $\lesssim 10^{-2}$ of the Poynting energy emitted by the pulsar is tied up in a form which may be ordered. Beyond the distance L', where $B_t \sim B_p$, it is useful to regard the structure as a field-confined beam carrying significant current and filled with an electron-positron plasma in nearly force-free orbits ($\boldsymbol{J} \times \boldsymbol{B} \sim 0$). This class of propagating beams has remarkable stability properties (Benford 1978).

The $\boldsymbol{J} \times \boldsymbol{B}$ force has a narrow peak near the boundary, where the net pressure difference between interior and exterior is supported by the toroidal field. At the core, B_p (nearly parallel to the beam axis) can dominate. A positronic jet may be described this way, since it is not fluid-like if it is cold. Treating jets as fluids requires dissipative scaling lengths which are small compared with macroscopic dimensions. This may not be true for highly relativistic, magnetically dominated flows in which the effective Larmor radius is large due to the nearly force-free field configuration (Benford and Book 1971). This differs from the Blandford and Rees type jets (1974). For example, the Crab pulsar produces a Goldreich-Julian current $I \sim 10^{24}$ cgs $= 3 \times 10^{14}$ A. For $\tan \Psi \sim 10^{-2}$, $I_b \sim 3 \times 10^{12}$ A, about 1% of the current needed to confine jets seen in quasars (Potash and Wartle 1982).

Beyond L', the dominance of B_t will cause the beam to slow its lateral expansion. However, this tendency can be offset by pressures from entrained matter.

a) Entrainment

Eventually this ordered helical structure must leave the evacuated zone near its pulsar and come to pressure equilibrium with nebular material. Magnetically dominated beams

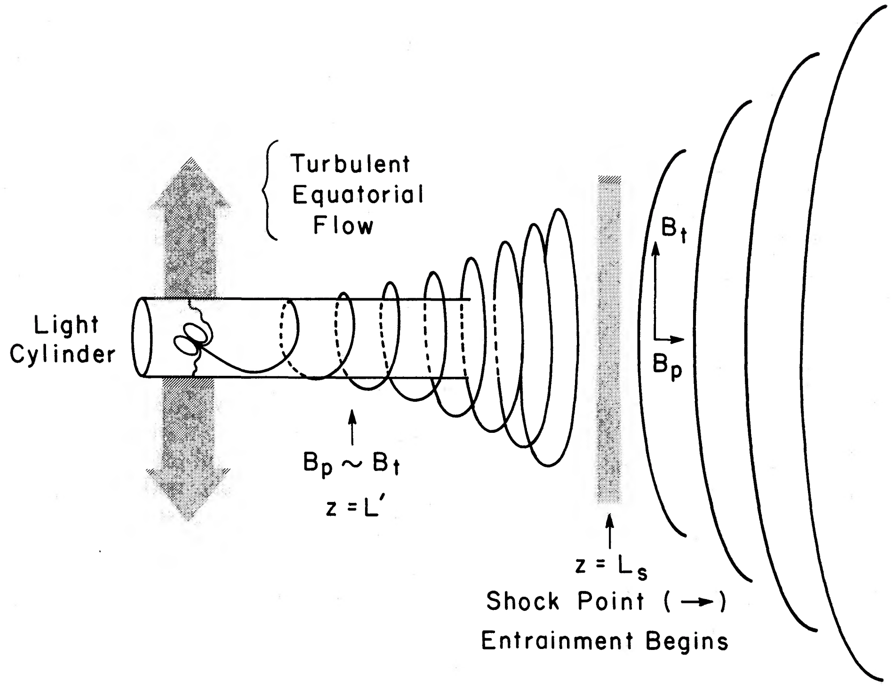

FIG. 2.—Emerging Poynting flux winds up as the jet expands, with toroidal field B_t winning out over polodial field B_p at $z = L'$. The bulk of the outflow leaves in the equatorial region as a hot plasma. At $z = L_s$ nebular matter shocks the jet, expanding it and further winding up the toroidal field as entrainment proceeds.

(or "jets") must satisfy a strong physical constraint: transporting magnetic flux outward. An axial shock will form when the beam reaches pressure balance with nebular material, and the slowing of the beam will build up B_t until it controls the dynamics of the downstream flow (Rees and Gunn 1974). Magnetic flows must move near light velocity, $v_j \approx c$, until they entrain material and can slow. This occurs near the shock at distance L_s, where v_j drops abruptly to $c/3^{1/2}$ and entrainment begins. Magnetically dominated winds form weak shocks, do not thermalize significantly, and thus convey energy downstream without dissipation (Kennel and Coroniti 1983). This means little new energy is available for synchrotron radiation beyond the shock site.

Annihilation can occur readily when an electron-positron jet entrains nebular matter. Intuition suggests that the minimum entrainment derives from charge neutralization of the electron-positron plasma by ions. Conserving momentum, the jet with initial density n^* and velocity v^* slows to $v_f = v^*(n^*m/n_f M_f)$, where M_f is the acquired ion mass (hydrogen) and m the electron mass. If the minimum entrainment occurs, $n_f \sim n^*$ and with $v^* \sim c/3^{1/2}$, $v_f \sim 3 \times 10^{-4}c$. Transverse pressure equilibrium,

$$\frac{B_t^2}{8\pi} + P_{\rm ext} = \frac{B_i^2}{8\pi} + P_{\rm rel} + P_{\rm thermal}$$

will adjust during entrainment. Here the internal magnetic field B_i and particle pressures $P_{\rm rel}$ (relativistic) and $P_{\rm thermal}$ act outward against the confining B_t and external gas pressure $P_{\rm ext}$. Let us assume that the two confining terms are approximately equal, since B_t at the jet boundary will compress until it is comparable to the external pressure $P_{\rm ext}$. If the right-hand terms are also comparable, as is usual in equipartition, then the resulting flow is quite complex. We expect entrainment to gradually expand the jet, but a useful solution requires knowledge of how dominant the confining field is upon beam dynamics. Let us assume the jet deceleration winds up B_t so that it remains a dominant pressure term, while thermalization of gradually entrained material forces the jet to expand. In this case $B_t^2 r_j^2 v_j \sim$ const, and using the estimate of v_f above for minimum entrainment,

$$r_j \sim r^* \frac{B_t(r^*)}{B_t(r_j)} \left(\frac{n_f M_f}{n^* m}\right)^{1/2} , \tag{2}$$

where r^* is the jet radius at the shock point. The jet smoothly expands to this radius if entrainment is slow. If entrainment proceeds beyond the minimum level we have estimated, the jet radius will exceed this limit.

These remarks are only indicative, since we have represented the complex physics of entrainment only by invoking momentum conservation. Beyond the entrainment region the jet becomes more fluid-like, though magnetic effects are still important. Emission-line regions in extragalactic jets may arise from similar entrainment (DeYoung 1981).

b) Establishing Magnetic Equilibrium

A current-carrying jet "finds" a return path to the magnetosphere by driving induction fields. At the head of jet the local increase in B_t drives a reverse electric field E_z, distributed over the general area near the jet head (Fig. 3). This E_z will drive the returning currents over an area A_r significantly larger than the jet area A_j, so that the return current $n_r v_r e A_r = n_j v_j A_j e$. This means the energy tied up in returning the current is $E_r = E_j(n_j A_j/n_r A_r) \ll E_j$, where E_j is the jet energy. E_r/E_j is small if the return current is carried in a dense plasma medium, $n_r \gg n_j$, and over a larger area. This is what occurs in laboratory situations. We expect, then, that magnetically ordered jets will terminate in low-energy return flows. This return point will move outward with time as long as the kinetic flux overcomes losses along the way and at the jet head.

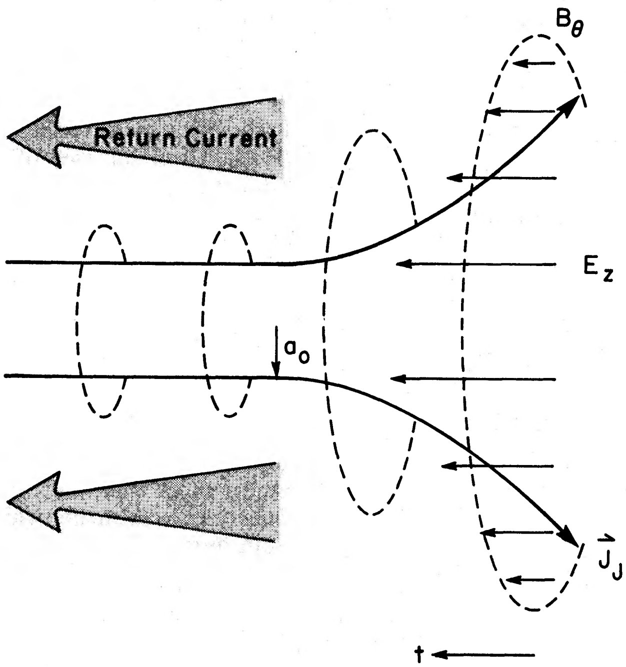

FIG. 3.—Return current generated at the head of a current-carrying jet is distributed over radii larger than the equilibrium jet radius, a.

Consider a model as depicted in Figure 3. Jet current flares radially near the jet head, forming a trumpet-like shape which reflects scattering of the particles by the dense medium, jet thermal pressure, and self-constriction as net current builds up. We assume an "emittance" which is constant in t, where $t = \tau - (z/v_j)$ and τ is time. Then the beam radius a is related to the equilibrium radius a_0 by $a^2 = a_0{}^2(t_r/t)$, where t_r is the rise time of the return current density generated in the ambient plasma, $J_z = J_j[1 - (t/t_r)] + J_0$, with J_0 the equilibrium current density carried by the jet (Fig. 3). The appropriate equilibrium model for a self-pinched jet with scattering is the Bennett distribution, $J_j = I_j/\pi(a^2 + r^2)$. The jet charge is conveyed away in the negligible time $\sim 10^{-4}(v/\omega_p)n_{-4}{}^{-1/2}$ s. The voltage drop at the jet head is an energetically negligible $\sim 100L^*(v/\omega_p)$, with L^* the length over which the drop occurs in parsecs. Thereafter the rising B_θ compresses the jet, driving E_r and E_z. The E_z-driven currents appear both inside the varying jet radius, a, and outside in the denser plasma. We shall neglect the fact that a real jet will evacuate its channel, thus forcing most of the return current to be carried in the outer sheath ("cocoon"). This effect is probably secondary compared with the fact that the inductive, E_z-driven return currents flow throughout the volume beyond $r = a$, with a slow radial dependence:

$$E_z(r, t) \propto \left[2\frac{t}{t_r} - \left(\frac{t}{t_r}\right)^2 - \frac{1}{2}\right] \ln\left(\frac{b^2 + a^2}{r^2 + a^2}\right) + \frac{t}{t_r}\left[\frac{1 - t}{t_r}\right] \frac{a^2(b^2 - r^2)}{(a^2 + b^2)(a^2 + r^2)} .$$

The boundary condition is that $E_z = 0$ at $r = b$, the radius where the return current is fully included. Clearly most of J_r will lie outside the jet for even moderate $b > 2a$.

This kind of solution implies that the jet head *necessarily* drives the return current in a broad cocoon. Typically B_θ will be large at the jet core and small throughout most of the cocoon. This picture applies, also, to magnetically dominated jets, as in the magnetic "bounce" models of Chan and Henriksen (1980).

c) *Instabilities*

As B_t winds up, the classical magnetic hose instability occurs. However, at distances $z < L'$ the expanding, current-carrying flow is not subject to virulent instabilities unless it stops expansion and becomes a diffuse-pinch type configuration. Then, typically it begins to kink. Other perturbations are known to be stabilized within the compass of energy principle calculations (Newcomb 1958). For $z < L'$, lateral expansion draws energy from the driving terms of the general Kelvin-Helmholtz modes (Hardee 1982). Supression of these modes by expansion at an angle dr/dz requires that

$$v_j r_j{}^{-1} dr_j/dz > \gamma_k(\delta B_t/B_t)^2 , \quad (3)$$

where γ_k is the temporal growth rate, $\gamma_k \approx 2\pi v_j(M\lambda)^{-1}$, with λ the axial wavelength and M the Alfven Mach number. The percentage change in the confining toroidal field is $(\delta B_t/B_t) \approx R/\lambda$ for a developed sidewise kink, since the radius of curvature $\sim \lambda$. Modes are stabilized at small amplitude by equation (3) if $dr_j/dz > M^{-1}(r_j/\lambda)^3$. Since growing modes exist only for $\lambda > 2\pi r_j$, we expect expansion at even small rates will nonetheless suppress growth.

Pinching modes will grow similarly to those in gas-confined jets (Cohn 1983). These can threaten a magnetically confined jet which has stopped lateral expansion, yielding growth lengths $\sim 25r_j$ for flows with $M \lesssim 5$ (Ferrari, Massaglia, and Trussoni 1982). There are a variety of unstable reflection and short wavelength modes available as well (Hardee 1982; Ray 1982). Pinching can compress the confined gas considerably, leading to enhanced synchrotron luminosity, perhaps by factors as great as 1000 as proposed by Achterberg, Blandford, and Goldreich (1983) for Sco X-1. Again, expansion can supress these modes.

Once the beam comes into pressure balance with the surrounding nebular material and begins entrainment, the heavy cocoon threaded by the jet self-field B_t lowers growth rates significantly, allowing survival far beyond the shock point L_s. The growth lengths are extended by a factor $(M_c/M_j)^{1/2}$ with M_j the jet mass per unit length and M_c the threaded cocoon mass per unit length. Factors $M_c/M_j \geq 100$ are not unlikely (Benford 1981). Since the jet will initially be supported by relativistic particles, it will be quite light in comparison with the surrounding supernova remnant. This ensures significant cocoon stabilization but also means buoyancy can warp the jet and perhaps affect growth rates (Achterberg 1982).

Expansion of factors of 10 or more, plus entrainment of ≥ 10 times the jet's initial mass can stabilize all modes for times ~ 1000 years.

IV. IMPLIED MORPHOLOGY

Two basic pictures emerge from our argument, depending upon whether the outflowing beam encounters appreciable pressure (and entrainment) during its passage through the supernova remnant.

a) *Crab-like Case (Filled Remnant)*

Fast, Crab-like pulsars with surface fields $\sim 10^{12}$ gauss probably have electron-positron densities exceeding the Ruderman-Sutherland values. This means Ψ will exceed the $\sim 10^{-2}$ estimated in § II, and the opening angle of the beam will be larger. Also, polar cap models yield average outflow densities, while real pulsars may preferentially load some field lines with extra plasma pressure (Arons 1981). These uncertainties can change the observed opening angle of the beam, but the point L' where B_t becomes dominant will still be $\ll 1$ light year.

Beyond the shock point, L_s, the jet radius r_j will evolve from competing pressures and entrainment. The tendency will be toward larger radii as the equilibria evolve, until the falloff in external pressure at the nebular boundary allows constant-r propagation. Generally, the following should appear:

1. Sharp edges, since the maximum toroidal pressure occurs at the beam boundary.
2. A buildup of entrained matter at the edge, where it is held in by the magnetic pressure. Seen in thermal (optical) lines, this will cause limb brightening.
3. Relatively little synchrotron radiation from original jet electrons, since electron-positrons will have radiated away their transverse energy and the equilibrium will have evolved to $\boldsymbol{J} \parallel \boldsymbol{B}$. Whatever synchrotron radiation does appear should also be limb brightened. The beam head may yield some radiation as well. However, if nebular electrons are entrained they will dominate and yield a filled synchrotron source.
4. Two jets occur generally, one from each polar cap, though one could be obscured by our perspective of the nebula. Though beams may bend from pressure and buoyancy in the nebula, they should straighten upon exiting.

For a surface field of 10^{12} gauss and Crab parameters, B_t at the distance L' (where it exceeds B_p) is ~ 0.5 gauss. If the beam continues to expand laterally at the same rate beyond L' (i.e., a cone), as it convects the magnetic stresses outward, it will still be quite thin, 10^{-2} lt-yr, at a distance 0.1 pc. At this point, L_s, it will confront an MHD shock, as in the wisps of the Crab. Here the self-field of the jet is $\sim 6 \times 10^{-3}$ gauss, about 20 times greater than the observed Crab nebular field. Thereafter B_t winds up due to slowing from entrainment, as described in § III. It seems plausible that the confining field, compressed between trapped entrained matter and external pressures, will maintain approximately the same strength. If so, assuming $B_t(r^*) \sim B_t(r_f)$ in equation (2), yields $r_j \sim r^*(M_f/m)^{1/2} \sim 0.5$ lt-yr, for the minimal entrainment $n_f \approx n^*$. The growth length for kinking perturbations will be reduced by further expansion during entrainment, plus the cocoon loading factor $(M_c/M_j)^{1/2}$ discussed in § III*c*. Growth lengths of several parsecs are then plausible, using equation (2) for the radius after entrainment and the estimate that $> 10 r_j$ is needed for significant growth after lateral expansion from entrainment has stopped.

It is tempting to connect such a magnetically confined beam with the "jet" projecting from the Crab (van den Bergh 1970; Davidson 1979; Gull and Fesen 1982). The jet does not point at the pulsar or the nebular center of expansion, suggesting significant bending. In our view the [O III] emission which makes it conspicuous must be entrained nebular material confined by the toroidal field (De Young 1981). The ability of the jet to withstand bending and retain sharp, straight edges argues that magnetic fields have put "backbone" in the system and tend to straighten it after the pressure gradients of the nebula have fallen behind. This might not be the case if the jet were a dense gas cloud left long ago by the presupernova star, and blown by an interstellar wind until it no longer pointed at the nebular center (Blandford *et. al.* 1983). It is difficult to see how such a cloud could remain so well collimated while being bodily moved.

Morphology.—Our estimates of $r_j \sim 0.5$ lt-yr, $v_f \sim 100$ km s^{-1} after entrainment are in rough agreement with the Gull and Fesen (1982) observations of jet radius ≈ 0.7 lt-yr and a radial velocity of $+150 \pm 100$ km s^{-1}. This velocity agreement implies that the jet does not lie nearly in the plane of the sky. Thus any counterjet can easily be blocked from our view of the nebula itself. Since we base our estimate on a guess that the confining magnetic field B_t changes little in magnitude during jet expansion and expulsion from the nebula, as well as the assumption of minimal entrainment, no truly quantitative comparison of r_i and v_i is possible.

The sharpness of the jet boundary we attribute to the compression of the confining fields to a thin layer, as is common in nearly force-free equilibria (Fig. 3). Gull and Fesen note "a tubelike structure rather than a gaseous sheet or filled cylinder," and "the jet's emission appears clumpy, consisting of several small discrete condensations." Such uneven emission implies illumination of clouds in the jet walls. These may represent sporadic entrainment, or density enhancements which are preferentially illuminated. Within our picture, relativistic electrons must deposit this energy, quite possibly through anomalous (non-Coulomb) mechanisms, as are appropriate for streams penetrating denser plasmas. Such processes are quite parameter-dependent (Hammer *et al.* 1978), so some spottiness in energy deposition seems likely.

Since we do not require the jet to be pointed at any special angle, its apparent length is approximately its true length, about a parsec.

Emission.—Fundamental to our view is that the original outflowing plasma is only midly thermalized at the shock point L_s, since the magnetically dominated shock is weak. The X-ray torus lying in the (supposed) equatorial region of the pulsar testifies to rapid thermalization and higher particle energies there (Swinbank and Pooley 1979). Lower electron-positron energies and weak shocks (along the jet; i.e., the pulsar spin axis) yield less synchrotron radiation overall. Thus $\gamma \sim 10^2$–10^3 particles travel along the jet, losing energy in thermal fashion to ions. Their synchrotron emission in the jet toroidal fields will be submerged in the general nebular emission until the jet is free.

Highly relativistic electrons ($\gamma > 10^5$) from the nebula may be entrained along with protons. They will dominate the radio image of the jet. Should it be resolved, we expect a *filled* source, since entrained electrons will not participate in the thin-walled current-carrying profile of Figure 3. However, since the jet is only 1 lt-yr thick, the jet would be a weak source even if it entrained electrons up to the nebular density.

There is ample energy in the outflow to power the [O III] line emission from clouds at ~ 10 eV; the required efficiency is less than 1%.

Since the final velocity after entrainment ~ 100 km s^{-1}, electron velocities in swept-up nebular matter will not exceed ~ 10 eV. As usual, the cooling function in this temperature region allows thermal instabilities with cooling times ~ 100 yr. Thus bright, spotty emission is favored.

This outflowing jet was set up in the earliest days of the pulsar, when it presumably moved at $v_j \sim c$, and has outraced the nebula. By now entrainment has slowed its net motion and the nebula is catching up. The hypothesized energetic, hydrogen-dominated remnant beyond the nebula can affect jet stability, but only if its perturbations are large and exceed v_A, the Alfvén velocity in the jet magnetic fields. For the minimal entrainment of equation (2), $v_A \gtrsim c$, so the jet reacts swiftly.

Heating of entrained matter by mildly relativistic electrons explains naturally why there is an overlap between the last continuum emission contour and the base of the jet. (The jet optical emission is swamped by stronger emission deeper in the nebula, however). This view of a magnetic structure, squeezed between the interior beam thermalization and outer pressures, is compatible with Gull and Fesen's remark that "its strong [O III] emission may indicate a high degree of ionization and low or average electron densities relative to the Crab nebula's bright filaments ($N_e < 10^3$ cm^{-3})." Indeed, though the jet was powerful enough to carve a path through the outflowing supernova ejecta, our density estimates show it was incapable of accelerating dense matter, compatible with the Trimble (1968) observation of no unusual proper motion of any of the brighter filaments near the jet.

Termination and Alignment.—The misalignment of the jet with the pulsar position by $\sim 15°$ is explained by Blandford *et al.* as arising from an interstellar wind which blows the jet to the side without rippling or destroying its straight walls. Though such a benign process seems unlikely to leave straight sides intact, we must seek an even more radical means of deflection, since our outflowing jet must have deviated from its ejection along the pulsar spin axis—which is often assumed to be aligned along the long axis of the nebula itself (Trimble 1968). An equilibrium with pressure balance, as in § III*a*, with magnetic pressure comparable to P_{ext}, will respond to the nebular pressure gradient. A transverse gradient comparable to the nebular pressure P_n over a scale height h of the nebular radius can turn a jet by $\tan \phi \sim P_n R/P_j h$ over a distance along the jet R. Since $P_j \sim P_n$, over a scale height the jet can bend $\sim 90°$. Thus whatever the initial ejection direction (i.e., whatever the orientation of $\boldsymbol{\Omega}$), a gradient such as assumed by Kennel and Coroniti (1983), for example, will provide the needed bending. The other pulsar polar axis may yield a magnetically confined outflow as well. This would appear as a counterjet (if it is not obscured by the nebula), but *not* necessarily exactly at 180° from the apparent jet, since buoyancy forces can deflect it along a different course. Thus a search for a counterjet, while interesting, may not be definitive in falsifying our model.

The fact that the jet extends beyond the remnant implies that it had a higher velocity earlier, and outran the nebula. However, at its present velocity as given by Gull and Fesen, it will be overtaken and perhaps destroyed. We can crudely estimate the jet pressure as in § III, but we cannot reliably say that it is capable of withstanding the pressure gradients now existing in the nebula.

Models of the Crab jet which rely on Rayleigh-Taylor instability in the nebular walls, such as Bychkov's (1975) and Kundt's (1983) imply that a unique unstable spot occurred, blowing matter out in a spectacularly ordered fashion. Unless a dominant confining field is also ejected, as Kundt supposes, it is difficult to see how the jet can be so straight. Otherwise, typically a flow will open at an angle $\sim M^{-1}$, with M the Mach number. If we imagine the flow is caused by a hole opening sidewise in the nebular shell, to appear straight the hole must be opening at relativistic speeds; but then the flow velocity must be relativistic as well, which is not observed.

b) Shell-like Case

Here no entrained nebular matter slows the beam or affects its expanding, conelike structure, limb brightened in the synchrotron radiation. If there is some weak nebular pressure, $P_n(R)$, a single pressure argument shows that if $P_n \propto R^{-n}$, the beam cone opening angle will scale as $R^{(n/2)-1}$, so if $n = 2$, θ is constant.

The limb-brightened cone will finally strike the shell, depositing relativistic particles and large magnetic fields in a thin, circular region. This signature, as well as the cone itself, can be disguised by other shell emission and projection effects, but it may correspond to a morphology seen recently in MSH 15-52 by Manchester (1983). The cone may be visible in the radio synchrotron region, since it will not be masked by nebular emission. The foregoing arguments about stability and evolution should hold as well. Beyond the shell, evolution of the jet should follow as in the Crab-like case. The shell wall is, however, clearly the place to search for prominent energy deposition and a clear signature.

V. CONCLUSIONS

We have explored the qualitative implications of jet emission from pulsar magnetospheres. Though a variety of morphologies can emerge, we focused on the Crab jet of Gull and Fesen as a possible example. This leads to two predictions for the Crab jet:

1. *Thermal emission.* We have argued here that this thermal yet highly ordered feature may arise in a magnetically confined jet. Since entrained matter probably accounts for the warmed clouds carried in the jet, the jet hydrogen/helium ratio should be the same as the nebula's. Also, a counterjet may be visible if our viewing angle is appropriate. This stands in contrast with the "in the wake of the red star" model of Blandford *et al.*, which predicts an enhanced hydrogen/helium ratio and no counterjet.

2. *Synchrotron emission.* If the thin jet edges can be resolved by the VLA, this would be compatible with our magnetic jet model. However, it is more likely that nebular electrons of $\gamma > 10^5$ will dominate the emission, giving a filled profile. The helical magnetic configuration will yield a substantial polarization. The required particle energies and magnetic fields would clearly imply some origin in the nebula, a clear contrast with the star-wake picture.

This paper speculates on a possible new class of astrophysical objects, magnetically ordered pulsar jets, not because data demand it but because they would provide a link to our understanding of extragalactic jets. As well, direct evidence of the anisotropic nature of pulsar-to-remnant coupling would provide valuable incentives to attack the full global problem of pulsar dynamics.

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