A MODEL OF THE POLARIZATION POSITION-ANGLE SWINGS IN BL LACERATAE OBJECTS

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Received 1984 May 29; accepted 1984 August 13

ABSTRACT

The polarization position-angle swings that have been measured in a number of BL Lacertae objects and highly variable quasars are interpreted in terms of shock waves which illuminate (by enhanced synchrotron radiation) successive transverse cross sections of a magnetized, relativistic jet. The jet is assumed to have a nonaxisymmetric magnetic field configuration of the type discussed in the companion paper on the equilibria of force-free jets. For a jet that is viewed at a small angle to the axis, the passage of a shock will give rise to an apparent rotation of the polarization position angle whose amplitude can be substantially larger than 180°. The effects of freely propagating shocks are compared with those of bow shocks which form in front of dense obstacles in the jet, and specific applications to 0727—115 and BL Lacertae are considered. In the case of 0727—115, it is pointed out that the nonuniformity of the swing rate and the apparent oscillations of the degree of polarization could be a consequence of relativistic aberration.

Subject headings: BL Lacertae objects — polarization — radiation mechanisms — radio sources: variable

I. INTRODUCTION

A number of BL Lacertae objects and highly variable quasars have now been found to display rapid swings in the radio polarization position angle (e.g., Ledden and Aller 1979; Altschuler 1980; Aller, Aller, and Hodge 1981; Aller, Hodge, and Aller 1981). The swings often accompany outbursts in the total flux density and are approximately linear in time—this has been the case, for example, in AO 0235 + 164 (Ledden and Aller 1979) and BL Lacertae (Aller, Hodge, and Aller 1981). This behavior, however, is not common to all sources. For instance, in the case of 0727—115 (Aller, Hodge, and Aller 1981), the swing occurred during a period of steady decline in the flux density, and the changes in polarization position angle (P.A.) appeared as a series of jumps.

The inferred association of BL Lac objects with relativistic jets whose emission is beamed toward the observer (which has recently gained support from VLBI observations of apparent superluminal motions [e.g., Phillips and Mutel 1982] as well as from spectral measurements [e.g., Worrall 1984]) led Blandford and Königl (1979) to interpret the polarization P.A. swings as a manifestation of the relativistic aberration effect in dense clumps that are accelerated by a beamed, supersonic jet. However, the maximum swing angle that is consistent with this model is 180°, whereas a number of sources are now known to exhibit apparently continuous swings in excess of this value. For example, Aller, Hodge, and Aller (1981) reported total polarization P.A. rotation of 440° during a 38-day period in BL Lacertae, and possibly as much as 340° over a period of more than 3 years in 0727—115. Thus relativistic aberration evidently does not provide a complete explanation of this phenomenon.

In this paper we propose a new model of polarization P.A. swings which accounts naturally for arbitrarily large apparent rotations, as well as for the observed variations in the characteristics of the swing. This interpretation is based on the force-free equilibrium model of magnetized jets that was presented in the preceding paper (Königl and Choudhuri 1985, hereafter Paper I). In that paper we showed that a magnetically dominated, super-Alfvénic jet which moves supersonically relative to a confining ambient medium while conserving magnetic helicity will have a minimum-energy state that is generally a superposition of an axisymmetric (m = 0) mode and a helical (m = 1) mode (eqs. [6] and [11], respectively, in Paper I). We further demonstrated that the helical mode, whose wavelength along the jet is fixed at ~5 times the jet radius, is energetically favorable and may dominate the field configuration when either the confining pressure or the magnetic flux is sufficiently low (see § IIc in Paper I). In Paper I we found that this model could account for many observed features of the total and the polarized synchrotron emission in resolved jets that are observed at a large angle to the axis. In the present paper we assume that the same model also describes the magnetic field structure in jets which display rapid polarization P.A. swings—namely, unresolved, relativistic jets that are observed at a small angle to the axis.

We suggest that the observed polarization P.A. swings result from the propagation of strong shock waves along a force-free jet in which the relative magnitude of the m = 1 field component is sufficiently large (e ≥ 1 in the notation of Paper I). Such shocks will "illuminate" (by enhanced synchrotron radiation due to particle acceleration and field amplification behind the shock) successive transverse cross sections of the jet (see Fig. I). The integrated polarized emission from each cross section has a well-defined position angle which is determined essentially by the m = 1 field orientation. Because of the twisted, oppositely directed flux tubes' geometry of the m = 1 mode (see Fig. 1 in Paper I), this position angle varies systematically along the jet. When viewed at a small angle to the jet axis, the propagation of a shock will thus lead to an apparent rotation of the polarization P.A. (Some preliminary ideas along these lines were discussed by O'Dea et al. [1983] in connection with the polarization P.A. swings in 0727—115.)

The swing characteristics predicted by this model are discussed in § II where, in addition to freely propagating shocks,
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Fig. 1.—Schematic representation of the model. A shock wave propagating along a magnetized jet with a nonaxisymmetric equilibrium field configuration will “illuminate” different magnetic field orientations in successive cross sections. This is illustrated by the rotation of the arrows, which represent the transverse components of the magnetic field vectors at the center of the jet. When viewed at a small angle to the jet axis, the propagation of the shock will lead to an apparent rotation of the synchrotron polarization P.A., with each 180° increment corresponding to motion across half a wavelength of the jet’s radius.

we also consider bow shocks attached to dense obstacles in the jet. The model is then applied to the interpretation of observed swings in § III. Our results are summarized in § IV.

II. POLARIZATION POSITION-ANGLE SWINGS IN RELATIVISTIC FORCE-FREE JETS

We first consider the polarization P.A. swings that are induced by relativistic shocks which propagate freely along the jet. In this case, it is necessary to distinguish between the Lorentz factors $\gamma_j$ of the jet and $\gamma_s$ of the shock, where we assume that $\gamma_s \gg \gamma_j \gg 1$. The compression of the jet magnetic field in the shock depends on the Lorentz factor, $\gamma_j = \gamma_s (1 - \beta_j \beta_s)$, that corresponds to the relative motion of the shock and the jet (speeds $\beta_j c$ and $\beta_s c$, respectively, where $c$ is the speed of light). When $\gamma_s$ is much greater than $\gamma_j$, the shock is ultrarelativistic ($\gamma_s \approx \sqrt{2/\gamma_j}$), in which case the compression factor of the transverse field component between the preshock and the postshock frames approaches $\sqrt{2}\gamma_j/y_j$ (e.g., Blandford and McKee 1976); the component parallel to the direction of propagation is, of course, unaffected by the shock.

In the ultrarelativistic limit, the Lorentz factor of the postshock gas (as measured by a stationary observer) is $\gamma_{ps} \approx \gamma_s/(\gamma_j^2)$.

The synchrotron radiation emitted in the postshock frame satisfies $\delta E'_{ps} \cdot (\delta B_{ps} + B_{ps}) = 0$, where $\delta E'$ and $\delta B$ are, respectively, the electric and magnetic fields of the radiation. But the dot product of electric and magnetic fields is a Lorentz invariant, so it vanishes also in the observer’s frame. Hence,

$$\left(\delta E_{obs} - \beta_{ps} \times E_{obs}\right) \cdot \left(\delta B_{obs} + n \times \delta E_{obs}\right) = 0,$$

where $B_{ps}$ is the (normalized) velocity of the postshock gas, and where we substituted $\delta B_{obs} = n \times \delta E_{obs}$ ($n$ being the direction from the emission point to the observer). From this equation one directly obtains the P.A. of the observed electric vector (Blandford and Königl 1979). If we identify the $z$-axis with the axis of the jet, then the observed magnetic field components are given by $B_{ps,z} = B_{ps} \cos \delta$ and $(B_{obs,z}/B_{ps,z}) = (B_{obs,y}/B_{ps,y}) = \gamma_{ps}$. Thus, if the vector $n$ lies in the $y$-$z$ plane and makes an angle $\delta$ to the jet axis, then the observed P.A. $\chi$ with respect to the $y$-$z$ plane is given by

$$\tan \chi = \frac{B_{ps,z}}{B_{ps,x}} \left[\frac{\beta_{ps} \cos \delta}{y_{ps} \gamma_{ps} \tan \delta} - \frac{1}{y_{ps} \gamma_{ps}}\right].$$

In the ultrarelativistic limit, the ratio $B_{ps,z}/y_{ps} B_{ps,x}$ is given in terms of the field components in the jet’s frame by $(y_j/y_x)$

$$(B_{ps,x}/B_{ps,y}) \ll (B_{ps,y}/B_{ps,x}).$$

Hence, the last term in the numerator of equation (2) can be neglected except when $B_{ps,z}$ is very close to zero (and $B_{ps,y} \neq 0$). If, in addition, the angle of observation is sufficiently small ($\delta \ll y_{ps}$), then it is seen that equation (2) reduces almost everywhere to $\tan \chi = -B_{ps,z}/B_{ps,x}$, which is just the nonrelativistic result in the limit $\delta \approx 0$.

The polarized emission from an unresolved shock is obtained by summing up the Stokes parameters over the cross section of the jet in the manner discussed in § IIIa (eqs. [18]–[20]) of Paper I. The P.A. swing is then found by considering the emission from successive cross sections traversed by the shock. Two representative examples, calculated for a pure $m = 1$ force-free field configuration and for $y_{ps} = 5$, are shown in Figure 2 by the dashed and solid curves, which correspond to $\delta = 3°$ and $\delta = 10°$, respectively. These plots illustrate the fact that the apparent rate of change of the polarization P.A. is highly nonuniform when $\delta$ is close to $\cos^{-1}(\beta_{ps})$ (see eq. [2]), even though the velocity of the shock is constant. This is because, for $\cos \delta \approx \beta_{ps}$, the radiation which reaches the observer is emitted in the postshock frame in a direction which is approximately perpendicular to the line of sight (a manifestation of the relativistic aberration effect), so the electric vector rotates in a plane which is roughly parallel to the line of sight. Under these conditions, the apparent swing is generally small except when the electric vector is nearly perpendicular to the $y$-$z$ plane (i.e., except when $[B_{ps,y}/B_{ps,x}] = [B_{ps,y}/B_{ps,x}] > 1$), which maximizes the Doppler effect and hence the apparent brightness of the shock (see Blandford and Königl 1979). An additional feature which is predicted by this model for sources that are viewed at an angle $\cos \delta \approx \cos^{-1}(\beta_{ps})$ is an apparent oscillation of the degree of polarization $P$ (see Fig. 3). In particular, it is predicted that the
minima of \( P \) should coincide with the polarization P.A. jumps exhibited by these sources (cf. Figs. 2 and 3).

Freely propagating shocks of the type that we have considered could develop from velocity fluctuations near the origin of the jet (Rees 1978), and their expected morphology is, in fact, consistent with the observed structure of certain bright emission knots in resolved jets (e.g., Biretta, Owen, and Hardee 1983). However, the shocks that give rise to polarization P.A. swings could also be associated with dense clumps which enter the jet and which are subsequently accelerated by the flow (see Blandford and Königl 1979). In this case, the postshock gas may be assumed to move with the speed \( \beta c \) of the clump. As can be seen from equation (2), the gradual increase of \( \beta c \) will lead to an apparent rotation of the polarization P.A. which will be most rapid when \( \beta_{ps} \) reaches \( \cos \delta \) (for \( [B_{ps1}/B_{ps2}] \neq 0 \)). This relativistic aberration effect was originally proposed as the source of the polarization P.A. swings in BL Lac objects (Blandford and Königl 1979). The amplitude of the swing depends essentially on the ratio \( B_{ps1}/B_{ps2} \) at the time of the most rapid rotation, but it cannot exceed 180°. Hence, as noted in § I, this mechanism cannot explain the much larger swings that have subsequently been measured in some of these sources. However, if the acceleration of the clump is sufficiently slow, then it may be possible for jet material encompassing more than half a wavelength of the \( m = 1 \) mode to pass through the clump shock before the clump attains the jet velocity, which would lead to an apparent swing in excess of 180°. A conservative estimate of the condition for the appearance of such a swing may be obtained by considering the initial stage of the acceleration, when the velocity of the clump is not yet highly relativistic. We imagine that a clump of (transverse) radius \( R_c \) and (upstream) rest-mass density \( \rho_j \) is immersed in a cylindrical jet of radius \( R_j(\gg R_c) \) and proper density \( \rho_j(\ll \rho_j) \). By integrating the equation of motion in the limit \( \beta_c \ll 1 \), we can estimate the acceleration time of the clump to \( \gamma_j \beta_c \approx 1 \) to be

\[
t_{acc} \approx 0.9 \frac{\rho_j}{\rho_j} \frac{R_c}{\gamma_j \beta_c^2 },
\]

where the numerical coefficient corresponds to an isothermal clump model to and adiabatic indices of 5/3 and 4/3 for the clump and the postshock gas, respectively (cf. Blandford and Königl 1979). A rough criterion for the occurrence of a \( \gtrsim 180° \) swing is that the distance (\( \sim 0.5ct_{acc} \)) traversed by the clump during the time \( t_{acc} \) should be larger than half a wavelength \( \lambda \) of the \( m = 1 \) mode as measured in the observer’s frame. Substituting for \( \lambda \) from equation (15b) in Paper I and taking into account the Lorentz contraction of the wavelength in the stationary frame, the condition becomes \( t_{acc} \gtrsim \pi R_j/1.25c\gamma_j \). Hence,

\[
\left( \frac{\rho_j}{\rho_j} \frac{R_c}{R_j} \right) \gtrsim 5.5\gamma_j ,
\]

which implies that the initial column density of the clump must be substantially larger than the transverse column density of the jet. (According to our numerical calculations, however, the coefficient on the right-hand side of equation (4) decreases to a value \( \lesssim 1 \) if \( t_{acc} \) is taken to be the total acceleration time to the velocity of the jet.)

## III. Applications

The model proposed in this paper can, in principle, account for arbitrarily large apparent polarization P.A. swings. In practice, the angular extent of the observed rotation is limited by the distance along the jet over which the shock remains strong enough to provide the necessary “illumination” of the nonaxisymmetric field configuration. The idea that the temporal swings correspond to a spatial variation of the mean polarization P.A. along the jet receives indirect support from optical and infrared polarization measurements of several BL Lac objects (e.g., Sitko, Stein, and Schmidt 1984), which indicate a systematic change in the polarization P.A. with wavelength. Since the jets associated with these sources are probably inhomogeneous (e.g., Blandford and Königl 1979), it is likely that the peak emission at different wavelengths is dominated by different regions in the jet, so that the observed dependence on wavelength reflects a variation with distance along the jet. Although we have focused on a specific field geometry, it is conceivable that other nonaxisymmetric field configurations could also account for the observed swing events. However, the quasi-periodic behavior of a source like 0727—115 as well as the general physical arguments given in Paper I strongly suggest that it is the \( m = 1 \) force-free mode which is, in fact, involved. It is worth pointing out, though, that other variants of this model could mimic the effect of the \( m = 1 \) field component. For example, one might imagine a jet with a directed longitudinal field (but random transverse field) which is being perturbed from its straight trajectory towards the observer by a helical Kelvin-Helmholtz instability, thereby giving rise to a projected normal field component which rotates with distance from the origin. In view of the constraints imposed by relativistic beaming, we regard this scenario as being less probable than the one we have adopted. Nevertheless, this example illustrates the possibility of producing shock-induced periodic swings even in jets which are not magnetic-pressure dominated.

The model outlined in this paper also provides a natural explanation of the apparent difference in the time evolution of the swing in sources like BL Lacertae and 0727—115 (see § I). As we demonstrated in Figure 2, the measured swing rate depends sensitively on the observation angle \( \delta \); in particular, for \( \delta \) near \( \cos^{-1}(\beta_{ps}) \), the polarization P.A. exhibits a series of jumps even in the case of constant shock velocity. This behavior is illustrated again in Figure 4, where we have used our model to fit the data in 0727—115. (In this figure, we have arbitrarily chosen the value of \( \gamma_{ps} \), and then adjusted the value of \( \delta \) for best results. However, other combinations of these
In particular, the 440° polarization P.A. swing that was measured in BL Lacertae in 1980 occurred during a major radio outburst (Aller, Hodge, and Aller 1981). Although the angular rotation rate during this swing was not strictly constant, the average value implies a time interval \( t_{180} \approx 15.5 \) days for a 180° rotation. According to our model, this time interval corresponds to the shock having traversed a distance \( (\pi/1.25)R_j \) in the frame of the jet. In view of the remarks made in the preceding paragraph, it is plausible to associate the shock in this outburst with an accelerating clump; in this case, the most rapid polarization P.A. swing would have occurred when the clump was still nonrelativistic, so

\[
R_j \lesssim 0.4\gamma_R c t_{180}^{-1} \tag{5}
\]

(cf. eq. [4]). The jet Lorentz factor \( \gamma_j \) can be estimated from the apparent superluminal expansion with \( \beta_{os} \approx 5 \) (for a cosmological redshift of 0.0695 and \( H_0 = 55 \text{ km s}^{-1} \text{ Mpc}^{-1} \) that was measured in this source in 1981 (Phillips and Mutel 1982). In fact, if the expansion occurred at the jet velocity and the source was observed at the Doppler-favored angle \( \approx \cos^{-1} \beta \) (see Blandford and Königl 1979). (Note that for \( \beta < 1 \) the Doppler-favored angle would be much smaller than \( \cos^{-1} \beta_{os} \), which is consistent with the nearly uniform apparent rate of the 1980 swing.) With this value of \( \gamma_j \), we get from equation (5) \( R_j \lesssim 8 \times 10^{16} \text{ cm} \). We may then estimate the distance \( z_{em} \) of the emission region from the origin by assuming that the jet has a conical geometry of half-angle \( \approx 2^\circ 4 \) (which is the apparent angle in the VLBI map of Phillips and Mutel [1982], multiplied by the estimated projection factor \( \sin \delta \approx 0.2 \)). This gives \( z_{em} \lesssim 2 \times 10^{18} \text{ cm} \). This result does not necessarily imply that the putative large-scale jet associated with this object is nonaxisymmetric (or even magnetically dominated) near the origin, since that jet need not be a direct extension of the VLBI jet considered here (see Henriksen, Bridle, and Chan [1982] and § IIIb of Paper I). The inferred limit on \( R_j \) is consistent with the upper limit of \( 8 \times 10^{17} \text{ cm} \) (for the adopted cosmological parameters) on the radius of the unresolved core in the 5 GHz VLBI map of Bââth et al. [1981]. This core probably represents the region which is optically thick at that frequency, and is presumably where the source of the observed polarization P.A. swing is located. In fact, since the swing was detected at frequencies of 8.0 and 14.5 GHz, but not at 4.8 GHz (see Aller, Hodge, and Aller 1981), it is likely that it originated fairly close to the boundary of the 5 GHz optically thick core. If future observations succeed in resolving the region where the polarization P.A. swings occur, then one could obtain from equation (5) a direct estimate of \( \gamma_j \). (This
equation does not apply, however, to swings induced by freely propagating shocks; in that case, $t_{180}$ depends on $\gamma_j$, $\gamma_s$, and $\delta$, so at least one other parameter [e.g., $\beta_{\text{obs}}$] would be required in order to estimate $\gamma_j$.

IV. SUMMARY

In this paper the force-free model of magnetized jets (Paper I) was applied to the interpretation of the large polarization P.A. swings that have been measured in certain BL Lac objects and highly variable quasars. Such swings could result from shock "illumination" of successive cross sections in a relativistic jet which has a nonaxisymmetric magnetic field configuration and which is observed at a small angle to the axis. This model accounts in a simple way for the occurrence of swings in excess of 180° and for the variations in the swing characteristics from source to source. In the case of 0727—115, the apparent jumps in the polarization P.A. and the accompanying oscillations of the degree of polarization could be attributed to relativistic aberration; this interpretation also implies that the P.A. of the putative jet in this source is orthogonal to the polarization P.A. in midswing. Although similar polarization effects could, in principle, arise also in other nonaxisymmetric field geometries or in jets undergoing a helical Kelvin-Helmholtz instability, the force-free equilibrium model offers what seems to be the most natural interpretation, one which leads to a unified picture of nonaxisymmetric phenomena in compact and extended jets.

We thank H. Aller for a discussion of the observations, W. Collins for assistance with the plots, and R. Henriksen and E. Parker for helpful suggestions. This work was supported in part by NASA grant NGL 14-001-001.

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