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Magnetic collimation of relativistic outflows in jets with a high mass flux

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ABSTRACT

Self-collimation of relativistic magnetohydrodynamic (MHD) plasma outflows is inefficient in a *single*-component model consisting of a wind from a stellar object or an accretion disc, in the sense that the collimated portion of the mass and magnetic fluxes is uncomfortably low. The theory of magnetic collimation is applied here to a *two*-component model consisting of a relativistic wind-type outflow from a central source and a non-relativistic wind from the surrounding disc. Through a direct numerical simulation of the MHD flow in the nearest zone by the relaxation method and a solution of the steady-state problem in the far zone, it is shown that in this two-component model it is possible to collimate into cylindrical jets, in principle, up to all the mass and magnetic fluxes that are available from the central source. In such a case the non-relativistic disc-wind plays the role of the jet collimator. It is also shown that this collimation is accompanied by the formation of a sequence of shock waves at the interaction interface of the relativistic and non-relativistic outflows.

Key words: MHD – plasmas – stars: mass-loss – stars: winds, outflows – ISM: jets and outflows – galaxies: jets.

1 INTRODUCTION

Collimated relativistic outflows in the form of jets are observed from several classes of astrophysical sources. For example, active galactic nuclei (AGN) produce extragalactic jets with rather high bulk flow Lorentz factors, $\gamma \sim 5$–10 (Urry & Padovani 1995), while some galactic superluminal sources are also associated with jets with Lorentz factors close to 2 (Mirabel & Rodrigues 1999). It is generally accepted that the leading mechanism capable of collimating an initially uncollimated outflow from such sources is the so-called mechanism of magnetic self-collimation, apparently a fundamental property of magnetized outflows from rotating astrophysical sources (Heyvaerts & Norman 1989). Basically, the toroidal magnetic field generated by the rotation of the source spontaneously collimates part of the outflow around the axis of rotation. This idea of magnetic collimation of astrophysical winds, first suggested by Lovelace (1976), Blandford (1976) and Bisnovatyi-Kogan & Ruzmaikin (1976), has been amply demonstrated analytically for non-relativistic (Vlahakis & Tsinganos 1998) and relativistic velocities (Chiueh, Li & Begelman 1991; Bogovalov 1995). In addition, exact magnetohydrodynamic (MHD) solutions of collimated outflows from central gravitating fields have been constructed in Sauty & Tsinganos (1994), Sauty, Tsinganos & Trussoni (1999), Vlahakis & Tsinganos (1999) and Vlahakis et al. (2000). Nevertheless, the general theory of magnetic self-collimation of magnetized winds, tells us nothing concerning the *fraction* of the collimated fluxes in the wind, or, the *distance* at which the collimated outflow is formed. To answer these questions direct calculations for every specific case should be performed.

In that connection, in recent years several numerical simulations have been performed for *non-relativistic* outflows showing the formation of jets from magnetized and rotating sources (Ouyed & Pudritz 1997; Krasnopolsky, Li & Blandford 1999; Ustyugova et al. 1999; Kudoh, Matsumoto & Shibata 1998, etc.). For example, in Bogovalov & Tsinganos (1999, henceforth BT99) and Tsinganos & Bogovalov (2000, henceforth TB00), full self-consistent solutions of the steady-state problem from small scales up to infinitely large distances from the source are obtained for different laws of rotation of the central source. In these studies it was found that the fraction of the cylindrically collimated part of the wind is of the order of 1 per cent of the total mass and magnetic fluxes of the initially (i.e. before rotation sets in) uncollimated wind from the source, when the source rotates uniformly. With a non-uniform rotation, typical of outflows from Keplerian accretion discs, the situation is even worse. This collimated outflow fraction decreases and at some field lines a decollimation effect could also be observed in some parameter regimes (see, e.g., Ustyugova et al. 1999).

Recent calculations of *relativistic* plasma outflows in the same model and with the same method as those used in BT99 show that a cylindrically collimated flow at large distances is formed in relativistic winds as well, in full agreement with the general theory of magnetic collimation. However, the fraction of the collimated mass and magnetic flux is even smaller than the corresponding values

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found in the case of non-relativistic winds (Bogovalov 2001). This fraction is of the order of 0.1 per cent of the total mass and magnetic flux outflowing from a rapidly and uniformly rotating source. The decrease of the efficiency of magnetic self-collimation in comparison with the non-relativistic case seems to be an intrinsic property of the relativistic plasma. There seem to be two effects accounting for such a result. First, the electric field generated by rotation becomes of a similar magnitude to the magnetic field at relativistic velocities and affects the plasma in a direction opposite to the collimating force arising from the toroidal magnetic field. Secondly, the electromagnetic energy flux increases the total inertia of the plasma, thus decreasing the response of the plasma to the collimating force (for more details see the discussion in Bogovalov 2001). A similar conclusion regarding the low efficiency of collimation has also been reached in Beskin, Kuznetsova & Rafikov (1998).

It is generally accepted that in astrophysical outflows associated with accretion discs the source of the energy in the outflowing plasma is the gravitational energy that is released by the accreting matter. Hence, the wind luminosity should be less than the luminosity of the accretion disc, where approximately half of the gravitational energy is released (Shakura & Sunayev 1973). Then, the bolometric luminosity of AGN, L_b, should be comparable with the energy released by accretion. According to our results (BT99) for a single-component model the bolometric luminosity in the jet, $\dot{E}_{jet}$, should be less than the accretion luminosity by a factor of the order of 10^3, $\dot{E}_{jet} \lesssim 10^{-3} L_b$. This theoretical result conflicts with observations where, for example, X-ray and gamma-ray emission is attributed to the emission from jets (Urry & Padovani 1995). And, sometimes the luminosity of the jets is less than the total luminosity of AGN at other wavelengths by only a factor of the order of 10 (Elvis et al. 1994). This energy flux of the hard emission can be considered as a lower limit for the energy flux in the jet (i.e. without beaming effects). If so, we have in AGN $\dot{E}_{jet} \gtrsim 0.1 L_b$. Also, analysis of observational data shows that at least in some non-relativistic jets from young stellar objects (YSO) the ratio $\dot{M}_{jet}/\dot{M}_{wind}$ could be even larger than unity (see the discussion in Bogovalov & Tsinganos 2001, henceforth BT01). Altogether, then, the calculated fraction of cylindrically collimated outflows seems to be uncomfortably low to fit observations of non-relativistic and relativistic jets, and the calculated ratio $\dot{M}_{jet}/\dot{M}_{wind} \lesssim 10^{-3}$ for relativistic outflows evidently contradicts observations.

This evident contradiction of theoretical predictions and observations can be considered as a challenge to the mechanism of magnetic self-collimation of astrophysical jets. One of the simplest resolutions of this contradiction for the case of *non-relativistic* outflows has been proposed in BT01. There it has been argued that a decrease of the toroidal component of the magnetic field in the central part of the outflow results in a more efficient collimation of this part of the wind caused by compression by the outer part of the outflow. This reduction of the inner toroidal magnetic field could be achieved in several ways: by a lower inner angular velocity, a reduction of the inner poloidal magnetic field, or, an increase of the velocity of the plasma in the central region. In BT01 we considered the collimation of the non-relativistic wind in a two-component model consisting of a central source and an accretion disc. Note that a similar idea has been discussed in the context of AGN (Sol, Pelletier & Asseo 1989; Pelletier et al. 1996) or YSO (Ferreira, Pelletier & Apple 2000).

In the particular case studied in BT01, the reduction of the toroidal magnetic field in the wind from the central source is achieved by assuming that the angular velocity of the central source is zero. Under this condition the disc-wind becomes the collimator of all the non-relativistic outflow from the central source. The main objective of the present paper is to demonstrate by direct numerical calculations that the mechanism of magnetic collimation can, in principle, provide collimation of a relatively large fraction of the outflow from a central source into a relativistic jet. To that end, in the following two sections we briefly outline the model and the method used, while in Section 4 the main results of our calculations are presented and discussed.

2 THE MODEL

In the model we adopt in this study the flow has two components. A radially expanding outflow having its origin at the central source and a non-relativistic outflow at the accretion disc that initially expands radially as well. We assume in addition that the wind from the disc is non-relativistic since relativistic winds are poorly collimated by magnetic stresses. The total poloidal magnetic field initially has a monopole-like structure. The central source is assumed to be spherical, while the accretion disc is attached to the spherical central source as a slab. The thickness of this slab defines the magnetic and mass flux from the disc. For simplicity, the radius of the disc is assumed to be twice as large as the radius of the central source. At this stage of our investigation we do not care about a close correspondence of the parameters of the model to some specific astrophysical object since our purpose here is to clarify the principal possibility to collimate relativistic winds and define the conditions under which this may take place. For the same reason, we omit the thermal pressure of the plasma (which is assumed to be cold) and the gravitational field. We are mainly interested on the electromagnetic stresses acting in the wind.

The only parameter that we specify close to reality is the Lorentz factor of the relativistic plasma, $\gamma \sim 5$. The velocity of the disc-wind was taken to be equal to $0.3c$. This velocity is already small enough to reduce the decollimating effect of the electric field and, on the other hand, it is still not too small compared with the velocity of the wind from the central source to provide strong gradients in the flow that might easily destroy the solution. To avoid strong gradients in the velocity and density, the initial values of these variables are smoothed with a function of the form

$$U(\psi) = \frac{(U_j - U_0)}{1 + \exp\{[\psi - (1 - \psi_{disc})]/0.05\}} + U_0, \qquad (1)$$

where U_j is the four-velocity of the plasma ejected from the central source, while U_0 is the four-velocity of the plasma from the disc. The density is taken to keep the mass flux ρU independent of ψ. The angular velocity changes at the interface between the central object and the disc from 0 to a specified value. The total magnetic flux from the source is normalized to unity at the equator.

3 METHOD OF SOLUTION

To obtain the steady-state solution of the problem over a wide range of scales, from distances comparable with the dimension of the central source up to much larger distances, we used a combination of two methods. A more detailed discussion of these methods is given in BT99 and Bogovalov (2001).

In our approach, the steady-state solution in the nearest zone that contains all critical surfaces and the governing partial differential equation (PDE) is of mixed type (elliptic/hyperbolic) is obtained by using a relaxation method, as in several other studies (cf. Ouyed & Pudritz 1997; Krasnopolsky et al. 1999; Ustyugova et al. 1999). We use the same software as used in Bogovalov (2001), except for the simple linear interpolation of the variables in the cells, which is replaced here by the van Leer (1977) interpolation scheme in the code for the time dependent simulation. This modification allows us to reduce strongly the usual Lax–Wendroff artificial oscillations present in our previous works.

A detailed discussion of the parametrization of the flow in the initially monopole-like solution is given in Bogovalov (2001). For the convenience of the reader here we recall that all geometrical parameters are expressed in units of the radius of the fast mode surface at the equator. The flow is described by the parameters α, σ and $\psi_{\rm disc}$. The parameter σ is defined as $\sigma = (R_{\rm f}/R_{\rm l})^2$, where $R_{\rm f}$ is the initial radius of the fast mode surface (i.e. before rotation of the disc starts), $R_{\rm l}$ is the radius of the light cylinder and $\alpha = \sqrt{\sigma}/U_0$, where U_0 is the initial four-velocity of the plasma. In the case of uniform rotation and uniform flow these parameters describe the total flow. In our case of non-uniform rotation and non-uniform flow we attribute these parameters to the field line leaving the inner edge of the accretion disc. The angular velocity of the disc rotation is taken such that $\alpha \sim r^{-\delta}$. For $\delta = \frac{3}{2}$ we have the familiar Keplerian law. The fraction of the magnetic field flux leaving the disc is defined by the parameter $\psi_{\rm disc}$. All the rest of the magnetic and mass flux originates at the central source.

In the second step, the solution in the far zone is obtained by extending to large distances the solution obtained in the nearest zone. This ability to extend the inner zone solution is based on the fact that the outflow in the far zone is already superfast magnetosonic. Therefore, the problem can be treated as an initial-value problem (a Cauchy-type problem) with the initial values being taken on an arbitrary surface located at the base of the far zone. The initial values on this surface are taken from the solution of the problem in the nearest zone. Since later we shall focus more on the solution in the far zone, the method of solution shall be briefly outlined here for the convenience of the reader.

The problem in the far zone is solved in an orthogonal curvilinear system of coordinates denoted by ψ and η. This system of coordinates has a rather simple physical meaning. The variable ψ in the axisymmetric flow denotes the flux function and gives the poloidal magnetic field $B_{\rm p}$ as,

$$\boldsymbol{B}_{\rm p} = \frac{\boldsymbol{\nabla}\psi \times \hat{\boldsymbol{\varphi}}}{r}, \tag{2}$$

where $\hat{\boldsymbol{\varphi}}$ is the unit vector in the azimuthal direction. A geometrical interval in these coordinates can be expressed as

$$(\mathrm{d}\boldsymbol{r})^2 = g_\psi^2\,\mathrm{d}\psi^2 + g_\eta^2\,\mathrm{d}\eta^2 + r^2\,\mathrm{d}\varphi^2, \tag{3}$$

where g_ψ and g_η are the corresponding line elements (components of the metric tensor).

Here the unknown variables are $z(\eta, \psi)$ and $r(\eta, \psi)$. The metric coefficient g_η can be obtained from the transfield equation (BT99)

$$g_\eta = \exp\left[\int_0^\psi G(\eta, \psi)\,\mathrm{d}\psi\right], \tag{4}$$

where

$$G(\eta, \psi) = \left[\frac{\partial}{\partial\psi}\left(\frac{B^2 - E^2}{8\pi}\right) - \frac{1}{r}\frac{\partial r}{\partial\psi}\left(U_\varphi v_\varphi c\rho - \frac{B_\varphi^2 - E^2}{4\pi}\right)\right] \times \left(U_{\rm p} v_{\rm p} c\rho - \frac{B_{\rm p}^2 - E^2}{4\pi}\right)^{-1}. \tag{5}$$

The lower limit of the integration in equation (4) is chosen to be 0, such that the coordinate η is uniquely defined. In this way η coincides with the coordinate z where a surface of constant η crosses the axis of rotation.

The metric coefficient g_ψ can be obtained from equation (2) in terms of the magnitude of the poloidal magnetic field,

$$g_\psi = \frac{1}{r B_{\rm p}}. \tag{6}$$

The equations for r and z are then,

$$r_\eta = -\frac{z_\psi g_\eta}{g_\psi}, \quad z_\eta = \frac{r_\psi g_\eta}{g_\psi}, \tag{7}$$

with g_η calculated using equation (4). Here $r_\eta = \partial r/\partial\eta$, $z_\eta = \partial z/\partial\eta$, $r_\psi = \partial r/\partial\psi$, $z_\psi = \partial z/\partial\psi$. For the numerical solution of the system of equation (7) the two-step Lax–Wendroff method on the lattice with a dimension equal to 1000 has been is used.

4 RESULTS

For a system consisting of a spherical central source to which is attached a rectangular disc with a radius twice as large as the radius of the central source, Fig. 1 shows the steady-state solution in the nearest zone. Since only the disc rotates, the poloidal electric currents are generated along the field lines that are rooted on this disc. A notable feature from this figure is the collimation of the initially radial outflow from the disc towards the axis of rotation and subsequent compression of the wind from the central source. It may also be seen that the Alfvén and fast MHD surfaces split in the lines that are rooted in the rotating disc.

In Fig. 2 we show the solution in the far zone. For this flow we were able to reach radial distances $R \approx 53 R_{\rm f}$, wherein the solution was terminated because of the formation of a shock wave. The formation of shocks during the process of collimation of a wind from a central source is apparently a rather general phenomenon and it was also found in our previous paper BT01 dealing with a non-relativistic

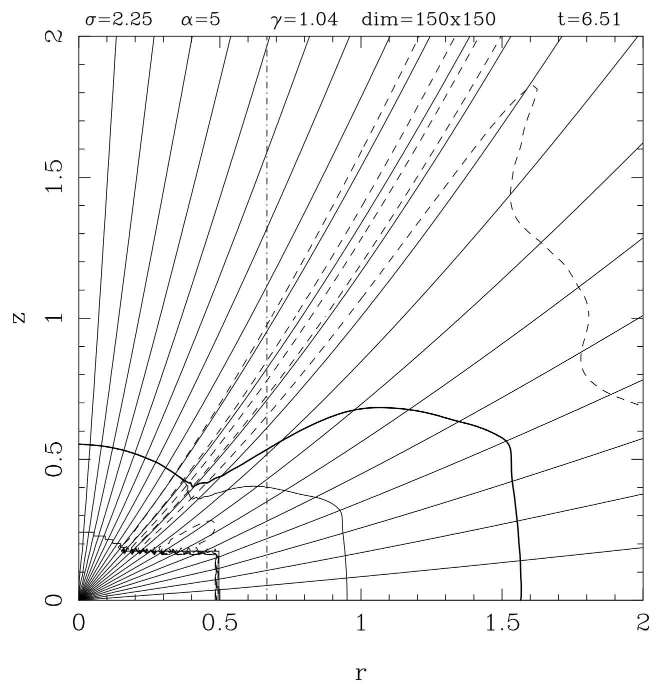

Figure 1. Solution of the problem in the nearest zone for the two-component model with $\psi_{\rm disc} = 0.7$. Thin solid lines indicate lines of poloidal magnetic field and dashed lines indicate poloidal electric currents. The thick solid line indicates the fast mode surface and the thin solid line indicates the Alfvén surface. The values of the parameters $\sigma = 2.25$ and $\alpha = 5$ refer to the inner edge of the disc.

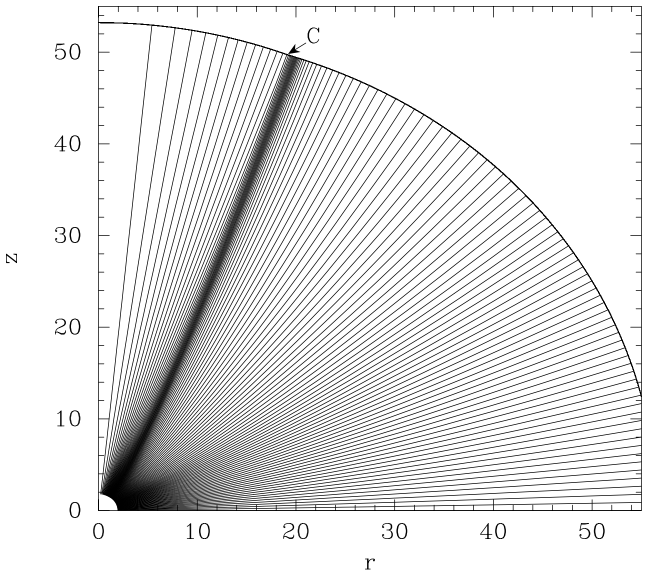

Figure 2. Magnetic field lines of the outflow shown in Fig. 1 but in the far zone. A notable feature in this case is the compression of the part of the inner relativistic flow from the central source into the thin layer of enhanced poloidal magnetic field and finally the formation on the inner part of this layer of a shock wave.

outflow. In the following, we consider this phenomenon in some more detail.

First, we note that the formation of shock discontinuities in collimated outflows from central sources has many similarities with shock formation in the flow of a supersonic gas past a concave wall discussed in Landau & Lifshitz (1959, p. 429). In our case this 'wall' is formed by the thin layer of the collimated plasma with enhanced magnitude of the poloidal magnetic field as can be seen in Fig. 2. Actually, here we deal with the situation of the collision of two components of the flowing plasma. One of them has already turned at some angle towards the system axis in relation to the initial direction of the flow, while the other flow component from the central source is still moving in the initial radial direction. This may be seen in Fig. 2. A surface of constant η at the termination of the calculations is shown there where this surface has a break at the point marked by C. The cusp is not evident in Fig. 2. The poloidal velocity of the plasma is perpendicular to this surface. This means that at this point the velocity of the plasma changes direction sharply. Fig. 3 shows the poloidal and toroidal magnetic fields along this surface of constant η. Two features are clearly seen in Fig. 3. First, the formation of a shock front in the poloidal magnetic field (at $\psi \approx 0.14$). And secondly, that the compression of the flow from the central source is caused by the toroidal magnetic field of the wind from the accretion disc.

Unfortunately, with our available numerical code we cannot pass through the shock wave. Our method of solution of the problem at large distances is based on the assumption of an adiabatic flow, which is violated at shocks. However, there is one way to go around this problem. In Fig. 4 we show a sketch illustrating the expected structure of the shock waves that could be formed during the collimation of the flow from the central source by the external collimator.

In general, a system of two shock waves and a weak discontinuity may be formed. The first shock indicated by '1' is the result of the collision of the two different flow components, as happens in the flow of gas past a concave wall discussed in Landau & Lifshitz (1959).

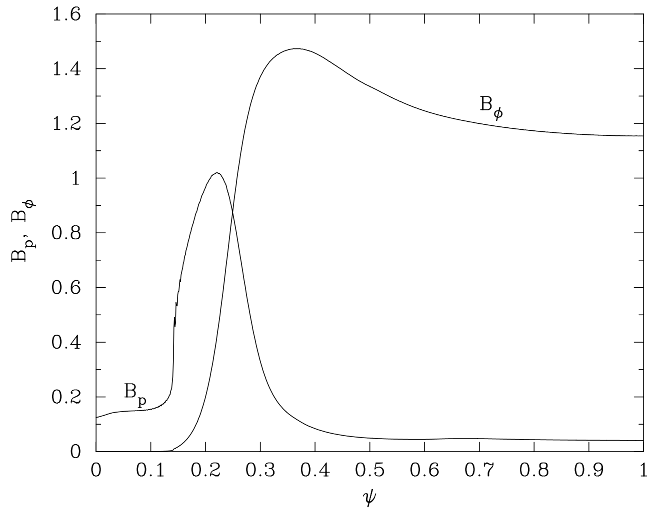

Figure 3. Distribution of the poloidal and toroidal magnetic fields along a surface of constant η crossing the z-axis at $53R_f$. The formation of a shock in the poloidal magnetic field at $\psi \approx 0.14$ may be seen.

This shock is accompanied by the formation of a weak discontinuity, labelled by '2' in Fig. 4(a). The final shock front indicated by '3' in this figure is formed by the reflection of the collimated flow from the rotational axis. Of course this reflection occurs not by the axis but from the conically converging flow itself.

It is important to note that the formation of shock '1' together with its associated weak discontinuity '2' is not always formed. Shock '1' is formed by the turning of the uncollimated flow from the central source towards the rotational axis by the wind from the disc. Before the formation of this shock, a layer with an increased poloidal magnetic field is formed in the relativistic flow. This layer

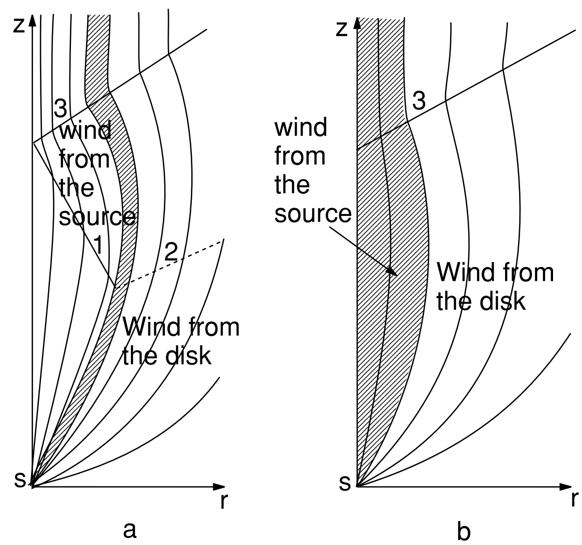

Figure 4. In panel (a) a sketch of the shock waves and singular surfaces that are expected to be formed in the general case of a two-component outflow is presented. The oblique shock front marked by '1' is formed at the collision of the two parts of the collimated and still uncollimated flows. An outgoing weak discontinuity from one end of this shock is marked by '2'. The shock front marked by '3' is formed at the self-reflection of the collimated flow at the axis of rotation. Under special conditions the collision shock may not be formed. In this case, the structure of a shock such as that shown in panel (b) is expected.

plays the role of the concave obstacle in Landau & Lifshitz (1959). The position of the starting point of the shock front in the flow past a concave profile is defined by the first intersection of one family of characteristics of the flow. The position of this point depends on the curvature radius of the obstacle and the ratio of the speed of sound to the local speed of the gas. It is worth stressing that this point moves away from the obstacle as the ratio C_s/v_p increases, where C_s is the speed of sound and v_p is the poloidal flow speed. Apparently, in our case the role of the speed of sound is played by the fast mode magnetosonic speed C_f, since the plasma is cold. The idea is to increase the ratio C_f/v_p and simultaneously to reduce the flux from the central source, in this way reducing the interval between the concave obstacle and the axis of rotation. The increase of the ratio C_f/v_p can be achieved by increasing the collimating force. Owing to this the crossing of the characteristics occurs close to the central source where the magnetic field is higher and therefore the ratio C_f/v_p is higher. We changed the index of variation of the angular velocity of the disc rotation from $\delta = 1.5$ used before to $\delta = 0.5$. Thus, the disc now rotates more uniformly. To reduce the space between the concave profile and the rotational axis we simply reduced the flux from the central source and increased the flux from the disc, so that $\psi_{disc} = 0.85$. Only 15 per cent of the total flux leaves the central source now. These two changes resulted in turning all of the flow from the central source to the rotational axis without the formation of shock front '1'. In some sense we moved this shock 'beyond' the axis of rotation. The resulting flow in the far zone is shown in Fig. 5.

A well-collimated jet can be seen in this case, confirming our expectations that in the two-component model the collimation of the relativistic plasma into a tightly collimated jet is indeed possible. Although there is no colliding shock wave '1' in this flow, the reflecting shock '3' is still present here. To demonstrate the presence of this shock, a portion of the flow pattern near the axis of rotation with an expanded scale on the *r*-axis is shown in Fig. 6. The position of the reflecting shock that traces the location of a break of the slope of the field lines is marked in this figure by a grey line.

It may appear unexpected that we were unable to cross the colliding shock front '1' while we were able to cross the reflecting shock front '3'. The explanation is the following. The amplitude of the reflecting shock is very small and in our case it appears as a weak discontinuity. In this case the jump in the thermal pressure is not significant and therefore the cold plasma approximation is valid through this shock front.

In Fig. 7 we show the final distribution of the toroidal and poloidal magnetic fields as a function of ψ, at $\eta = 550$. The poloidal magnetic field is uniform in the jet and then goes to zero sharply for $\psi > 0.15$. It is seen that all the magnetic flux from the central source with $\psi < 0.15$ is accumulated into the jet. The toroidal magnetic field compressing the jet is approximately equal to the poloidal magnetic field at the interface of the fluxes from the central source and from the disc. Note that in Fig. 7 some oscillations appear in the toroidal magnetic field that extend up to $\psi = 0.41$. This oscillating structure is typical of collisionless shock waves. In our case this is formed

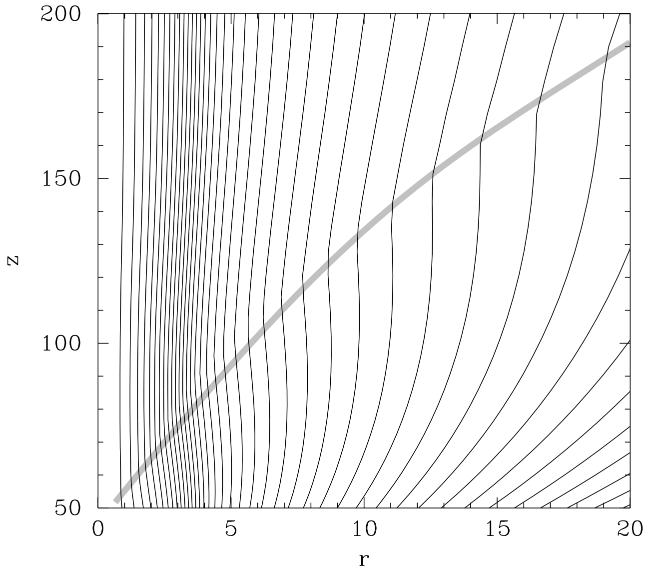

Figure 6. Magnification near the axis of rotation of the flow pattern shown in Fig. 5. The scale on the *r*-axis is increased. The reflecting shock '3' is shown by a thick grey line.

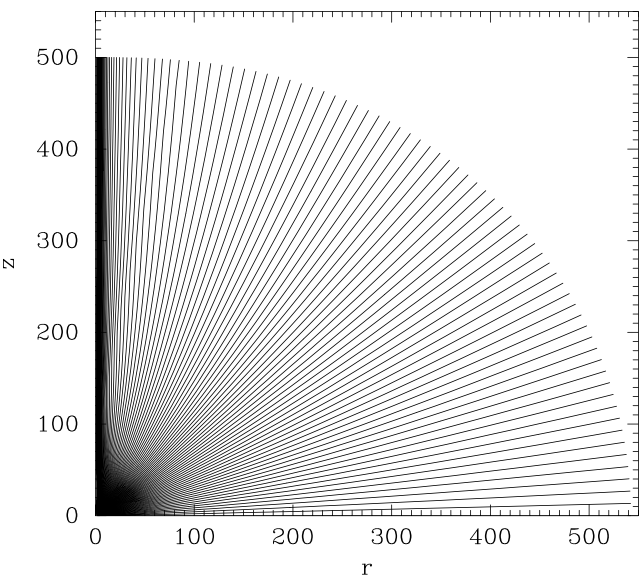

Figure 5. The flow in the far zone with the relativistic jet carrying 100 per cent of the mass flux from the central source. The Lorentz factor of the jet is $\gamma = 5$.

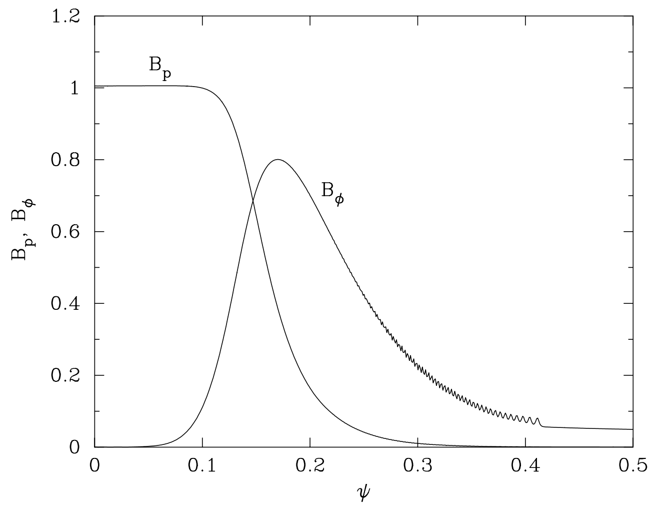

Figure 7. Distribution of the magnetic fields along the last surface of constant $\eta = 550$. The poloidal magnetic field in the jet is uniform. The boundaries of the jet correspond to the magnetic flux ψ where the poloidal magnetic field falls by a factor of 2.

after the reflecting shock wave '3'. The damping of such oscillations is caused by the usual numerical viscosity of the code.

5 CONCLUSION

In this paper we have used a simplified model to demonstrate that the mechanism of magnetic collimation of outflows may provide collimation of a remarkable fraction of the total magnetic and relativistic mass flux from a source provided that the system consists of two components: an initially uncollimated relativistic plasma from the central source and a non-relativistic wind from the surrounding disc. For the relativistic jet we were able to obtain a steady-state solution having a Lorentz factor of $\gamma = 5$. The total magnetic and mass flux from the central source of the relativistic outflow was approximately 15 per cent of the total flux from our system. It is important that all (100 per cent) the mass from the central source should be collimated into a relativistic jet. There seems no doubt that similar results could be obtained if we were able to extend our solution to large distances from the central source. The main difficulty in that respect is to pass through the collision shock wave.

A detailed calculation of the collimation process in conditions appropriate to specific astrophysical objects has not been performed in this paper. Here our purpose was restricted to resolving the difficulty of the theory of magnetic collimation in forming jets with a large fraction of the total mass flux of the outflow from the central source. We succeeded in obtaining this rather interesting result, by using an admittedly simplified and crude model. Another byproduct of this work is that some important processes accompanying the collimation of astrophysical plasmas, such as the formation of shocks, occur in the supersonic or MHD superfast region. A shortcoming of the present study, as far as a direct application of the results to jet formation in specific astrophysical objects is concerned, is that some additional physical ingredients (thermal pressure, non-zero angular velocity of rotation of the central source, etc.), have not been included in the present modelling. However, we feel that this will be worthwhile only if we solve the problem of passing through all shock waves formed in the far zone of the flow. This undertaking remains a challenge for the future.

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