The influence of magnetic fields on neutrino-dominated accretion disc

Yi Xie, Zhao-Yan Huang, Xiang-Fu Jia, Shao-Juan Fan and Fang-Fang Liu

School of Physics and Information Engineering, Shanxi Normal University, Linfen 041004, China

Accepted 2009 May 23. Received 2009 May 23; in original form 2008 September 24

ABSTRACT

We consider the influence of magnetic fields on the model of neutrino-dominated accretion flows (NDAFs) for gamma-ray bursts (GRBs) via the assumption that the accretion rate of the disc is totally caused by the torque of the Lorentz force, i.e. the magnetic braking of large-scale magnetic fields and magnetic viscosity of small-scale magnetic fields. We calculate the structure, composition, luminosity of neutrino emission and the Poynting flux, and the rate of mass loss driven by neutrino heating or launched centrifugally by large-scale magnetic fields, based on the physical condition of the magnetized NDAFs. It is shown that the magnetized disc is favourable to interpret the diverse prompt emissions as well as the X-ray flares observed in the early afterglow of GRBs.

Key words: accretion, accretion discs – magnetic fields – methods: numerical – gamma-rays: bursts.

1 INTRODUCTION

The popular theoretical explanations for the central engine of cosmological gamma-ray bursts (GRBs) can generally be divided into two sets of scenarios: the core collapse of massive stars for long bursts (Woosley 1993; Paczynski 1998; Hjorth et al. 2003; Stanek et al. 2003) and the coalescence of neutron stars or black hole (BH) binary systems for short bursts (Eichler et al. 1989; Narayan, Paczynski & Piran 1992; Fryer & Woosley 1998). It is believed that the two processes give rise to a BH of several solar masses with a magnetized disc or a torus around it (Meszaros & Rees 1997b), and many authors have studied the accretion model for GRBs by assuming a steady-state accretion (e.g. Popham, Woosley & Fryer 1999, hereafter PWF; Narayan, Piran & Kumar 2001, hereafter NPK; Di Matteo, Perna & Narayan 2002, hereafter DPN; Kohri & Mineshige 2002; Chen & Beloborodov 2007). This accretion mode is referred to as neutrino-dominated accretion flows (NDAFs) since they are mainly cooled via neutrino emission.

The neutrino mechanism can fuel many GRBs; however, it is somewhat difficult to power the brightest ones due to their relatively low efficiency. Notably, they cannot fuel the X-ray flares that were observed in the afterglow of the prompt bursts in Swift era (Fan, Zhang & Proga 2005). In contrast to the neutrino mechanism, magnetic mechanisms are promising candidates for the central engines of both the bursts and their X-ray flares. For instance, the Blandford & Znajek (1977) process as the central engine for GRBs was suggested by Meszaros & Rees (1997), and was also discussed in detail by Lee, Wijers & Brown (2000). More recently, Dai et al. (2006a) argued that X-ray flares can be produced by differentially rotating, millisecond pulsars after the mergers of binary neutron stars. Proga & Zhang (2006, hereafter PZ06) suggest that the accretion flow may not always be in a steady state. Rather, magnetic fields accumulated near the BH can form a magnetic barrier that temporarily blocks the accretion flow. This makes some dormant epochs at the central engine. The breaking of the barrier would lead to restarting the central engine, which is required to explain the recent observations of X-ray flares. The interrupted accretion mode was also suggested by Igumenshchev, Narayan & Abramowicz (2003), Narayan, Igumenshchev & Abramowicz (2003) and Spruit & Uzdensky (2005).

It is widely known that the accretion in the disc is driven by viscous stress, and the kinematic viscosity caused by ordinary molecular is extremely inefficient; thus, many authors (e.g. Balbus & Hawley 1991) suggested that magnetic fields on the disc could play an effective role in transferring the angular momentum of the disc. If this is also authentic in the NDAFs context, because of its hyperaccretion rate, the magnetic fields on the disc must be extremely strong and may have considerable effects on the disc. Meanwhile, if short GRBs are generated through compact star mergers, as is supported by Swift observations, the jet powering the late X-ray flares must be launched via magnetic processes rather than via neutrino-antineutrino annihilations (Fan et al. 2005). The above given two reasons motivated us to consider the influence of magnetic fields on NDAFs.

In Xie, Huang & Lei (2007, hereafter X07) , we improved the model of NDAFs by considering the effects of magnetic braking, and the luminosity and stability properties of the inner region of the disc were generally discussed. However, X07 discussed only the critical case where the growth rate of the azimuthal component of the magnetic fields generated by differential rotation from the radial fields equals its loss rate by buoyancy (thus the magnitude of the fields was possibly overestimated). Moreover, in the above condition, the overpressure fields were inevitably formed at about ten times the Schwarzschild radius. In this paper, we adopted the configuration of magnetic fields in which the toroidal magnetic fields are likely to dominate the accretion disc. It is based on the
recent analysis of the effects of strong toroidal fields on magnetorotational instability (MRI) by Stone et al. (1996), as well as other analyses in the literature. Then, we extended the X07 analysis to a general case via a relation between the magnetic fields and disc pressure, in which the magnetic pressure reaches only about 15 per cent of the total pressure of the accretion flow. This result allows us to determine important properties of the magnetized disc, such as the neutrino luminosity and the output power of the Poynting flux; we are also able to survey the disc’s composition and the mass loss as outflow produced by the neutrino irradiation and centrifugally driven by poloidal magnetic field lines, respectively. This work is organized as follows. We present the physics of magnetized NDAFs and some properties mentioned above in Section 2, and their numerical results are given in Section 3.

2 THE MODEL

2.1 Physics of magnetized NDAFs

The main challenge to the theory of accretion discs is the uncertainty of the nature and magnitude of the viscosity, since the angular momentum transport resulting from the action of ordinary molecular viscosity is extremely inefficient (Pringle 1981). It is widely believed that the magnetic fields on the disc can greatly affect angular momentum transfer and hence the accretion rate via a variety of modes (e.g. Blandford 1976; Blandford & Payne 1982; Balbus & Hawley 1991). Since it is the basic law of the interaction between electromagnetic fields and charged ions, we assumed that the accretion rate is completely caused by the torque of the Lorentz force (the ordinary molecular viscosity was ignored completely). We also suppose that the viscous dissipation was caused by small-scale magnetic fields in the disc (Balbus & Hawley 1991), and, meanwhile, the large-scale magnetic fields threading the disc extract part of the rotational energy and produce magnetic waves that escape the disc as the Poynting flux, i.e. the mechanism of magnetic braking. Then, the magnetic viscosity \( \nu_{\text{mag}} \) (the superscript means magnetic origin) can be parametrized as \( \nu_{\text{mag}} = \alpha_{\text{mag}} c_s H \) (Shakura & Sunyaev 1973; Pringle 1981), where \( c_s \) is the sound velocity of disc \((c_s^2 = \frac{P}{\rho}, \) in which \( P \) is the total pressure of the accretion flow and \( \rho \) is the density of the disc\) and \( H \) is the half-thickness of the disc. Invoking hydrostatic equilibrium perpendicular to the disc plane, \( \nu_{\text{mag}} = \frac{P}{\rho} \frac{m_B}{Z} \frac{r^2}{D} \). For a Keplerian orbit, we have \( \Omega = \sqrt{\frac{GM}{r^3}} \), and magnetic viscosity can be written as

\[
B_s B_0 / 4\pi = \frac{2}{3} \alpha P,
\]

where \( B_s \) and \( B_0 \) are, respectively, the radial and azimuthal component, and for simplicity, we omitted the superscript ‘mag’ of \( \alpha \).

Various local shearing sheet simulations of accretion discs have found that the quantities \( B_s^2, B_0^2, B^2, B_s \) all have somewhat standard ratios. For example, table 2 of Stone et al. (1996) gives (note that their \( x, y, z \) correspond to \( r, \phi, z \))

\[
B_s^2 \simeq \frac{1}{10} B_0^2,
\]

and

\[
B_z^2 \simeq \frac{1}{20} B_0^2.
\]

This gives a simple prescription which can be combined with any assumed value of alpha to solve for the field components. We assume the above equations are always satisfied in magnetized NDAFs.

Incorporating equations (1)–(3), we have

\[
B_z^2 / 8\pi \simeq \frac{10\alpha P}{60}.
\]

Equation (4) consistently describes a magnetized disc with angular momentum transport totally caused by the torque of the Lorentz force. We can naturally get the magnitude of magnetic fields through the relation between the poloidal fields \( B_z \) and disc pressure \( P \) for a given radius \( r \) and parameter \( \alpha \).

Following NPK and DPN, the conservation of mass and the approximate expression for angular momentum balance give the relation for mass accretion rate as \( M = 6\pi \sigma \rho H \Omega_1 \approx 6\pi \nu \rho H \) (for now we ignore the mass of the wind). Taking the magnetic braking effects into account, the relation is modified by

\[
M = 6\pi \sigma \rho H / \Omega_1 + \frac{4(B_s B_0 / 4\pi)^2}{2\pi r}.
\]

The derivation of equation (5) is given in Appendix A.

In the equation of state, we include the contributions from radiation pressure, gas pressure, degeneracy pressure, neutrino pressure, and, different from that of NPK and DPN, the magnetic pressure is included as

\[
P = P_{\text{gas}} + P_{\text{rad}} + P_{\text{deg}} + P_{\nu} + P_B.
\]

The gas pressure from the free nucleons and \( \alpha \)-particles is \( P_{\text{gas}} = \frac{m_n}{m_e} \frac{1}{4} T^4 \), where \( \alpha \) is the radiation constant, \( T \) is the disc temperature and \( X_n \) is the mass fraction of free nucleons (see Woosley & Baron 1992; Qian & Woosley 1996, hereafter QW96). The radiation pressure is \( P_{\text{rad}} = \frac{1}{8\pi} a T^4 \), where the factor \( \frac{1}{8\pi} \) includes the contribution of relativistic electron–positron pairs. In the degeneracy term \( P_{\text{deg}} = \frac{2\mu_e}{3}(\frac{1}{\Omega_1})^2(\frac{1}{\sqrt{\pi}}) \), \( \mu_e \) is the mass per electron, and it is taken as 2 by assuming equal number of protons and neutrons. The neutrino pressure \( P_{\nu} \) is \( P_{\nu} = u_e / 3 \), in which \( u_e \) is the neutrino-energy density which is well defined in DPN. Finally, in terms of magnetic pressure, we consider the three components in the calculation, i.e. \( P_B = B_z^2 / 8\pi = (B_s^2 + B_0^2 + B_z^2) / 8\pi \).

In the energy balance equation, the viscous heating equals neutrino-radiative losses plus advective loss and the fraction of rotational energy extracted by large-scale magnetic field directly:

\[
Q_{\text{vis}} = q_{\text{adv}} + q_{\nu} + Q_{\text{B}}.
\]

In which \( Q_{\text{vis}} \) represents the viscous dissipation, and we take it as the approximation \( Q_{\text{vis}} = \frac{3GM}{4\pi T} \). \( q_{\nu} \) represents the total cooling rate due to both neutrino losses (using the prescription given by DPN, which includes neutrino-optical-depth effects) and photodisintegration. \( q_{\text{adv}} \) is the advective cooling rate; we approximate it by (see e.g. Narayan & Yi 1994; Abramowicz et al. 1995)

\[
q_{\text{adv}} = \Sigma v_{\text{ds}} \frac{dT}{dr} \simeq \alpha T \frac{1}{3} \frac{1}{\Omega_1} \left( \frac{11}{3} a T^3 + \frac{3}{2} m / 4 \right).
\]

where \( s \) is specific entropy, \( \Sigma = -\Delta \ln s / \Delta \ln r \) is assumed to be equal to 1. \( Q_{\text{B}} \) represents the rotational energy of the disc extracted by large-scale magnetic fields anchored at the disc surface and can be calculated as \( Q_{\text{B}} = \Omega_0 \tau_0 T / 4 \) \( \Omega_0 \) (see Lee et al. 2000), in which \( dT \) is the torque exerted by the annular ring with width \( Dr \) of the disc due to the Lorentz force (see equation A2, and \( ds = 2\pi r \) dr). Then, we have

\[
Q_{\text{B}} = 2\pi \Omega_0 (B_s B_0 / 4\pi);
\]

this fraction of energy does not contribute to heating the disc matter, but escapes the disc as magnetic waves. Thus, it should be subtracted from viscous dissipation. This is also consistent with equation (5).

---

© 2009 The Authors. Journal compilation © 2009 RAS, MNRAS 398, 583–590
Equations (4)–(7) contain four independent unknowns, $P$, $\rho$, $T$ and $B$, as functions of $r$ and constitute a complete set of equations which can be numerically solved with given $M$, $\alpha$ and $M$. In the following calculations, we fix $M = 3 M_\odot$ (the corresponding Schwarzschild radius $R_s$ is $2GM/c^2 = 8.85 \times 10^3 \text{cm}$), and $\alpha = 0.1$.

2.2 Composition of disc matter

The neutrino-optical depth is dependent on the composition of the disc, such as the electron fraction, the free nucleon fraction and the electron degeneracy (e.g. Liu et al. 2007), and particularly, the neutron content of the disc is important for determining the properties of the disc’s outflows. We use a simplified treatment (see Metzger, Piro & Quataert 2008, hereafter M08) to calculate the proton mass fraction $X_p$, neutron fraction $X_n$ and $\alpha$-particle fraction $X_\alpha$ along the radius of the disc, in which the values of electron fraction $Y_e[Y_e = X_p/(X_n + X_p)]$ were fixed at 0.1 and 0.3. If we assume that the disc is sufficiently cold that the position and electron capture rates are negligible (it is valid at relatively large radii), the difference between the free neutron and proton mass fraction can be written as

$$X_n - X_p = 1 - 2Y_e.$$  

(10)

Nuclear statistical equilibrium (NSE) between protons, neutrons and $\alpha$-particles is established in the disc when the disc temperature is $\geq 0.5 \text{ MeV}$, a condition that is likely to be fulfilled for the hyper-accretion disc. Thus, we can determine the composition using the NSE, which is expressed by the Saha relation (Shapiro & Teukolsky 1983)

$$X_n^2 X_p^2 = 1.54 \times 10^4 X_n \rho_{10}^{-3} T_{10}^{9/2} \exp\left(-\frac{32.81}{T_{10}}\right).$$  

(11)

Mass conservation requires

$$X_n + X_p + X_\alpha = 1.$$  

(12)

By combining equations (10)–(12), we solve for all of the mass fractions at given results of $\rho$ and $T$ from Section 2.1.

2.3 Magnetic fields accumulating in inner region

PZ06 noted that the flows of collapsar and radiatively inefficient accretion flow (RIAF) simulations are similar during the early phase of their evolution. The late evolution of RIAFs shows that the torus accretion can be interrupted for a short time by a strong poloidal magnetic field in the vicinity of a BH. This mainly motivated their suggestion that the extended GRB activity may be a result of an accretion disc. Thus, we would like to check whether the ‘magnetic barrier’ described in PZ06 could occur in the model of magnetized NDAFs by comparing the time-scale of the magnetic fields diffusing and the time-scale of the fields accumulating via the process of accretion. During the accretion, the inner edge of the disc would, in effect, act as a one-way membrane for collecting magnetic fields and, starting with arbitrarily small values of the disc magnetic field (Ghosh & Abramowicz 1997), would be able to build up the magnetic fields threading the horizon, $B_H$, to the value that in turn is sufficient to stop the accretion processes through its strong radial pressure $B_H^2/8\pi$. The azimuthal component of the magnetic fields is buoyant, and will rise to the surface of the disc on a time-scale:

$$t_B \approx \frac{H}{v_A} = \frac{H(4\pi\rho)^{1/2}}{B}.$$  

(13)

where $v_A$ is the Alfvén velocity, $v_A = B/(4\pi\rho)^{1/2}$. Since the azimuthal component of the fields $B_\phi$ is related to the poloidal component $B_z$ through equations (1) and (2), the quantity $t_B$ can be used to describe the time-scale of the magnetic fields diffusing. Meanwhile, it is assumed that the fields are frozen in the disc gas; thus the viscous time $\tau_v$ can present the time-scale of the magnetic flux accumulating in the vicinity of the BH, and we write

$$t_v = \int_{3R}^R \frac{1}{\Omega r} dr,$$  

(14)

in which $v_r$ is the radial velocity of the disc, $v_r = \frac{1}{2} \frac{\Omega r^2}{\tau_v}$.  

2.4 Neutrino luminosity and BP power

For a comparison of our model with standard NDAFs, we calculate the neutrino luminosity of NDAFs with and without magnetic fields in this section. Following DPN, the neutrino transfer and neutrino opacities are all self-consistently considered in both cases (for detailed expression, see DPN). The neutrino luminosity is given by

$$L_\nu = \int_{R_{\text{min}}}^{R_{\text{max}}} \frac{Q^*_\nu}{4\pi} dr,$$  

(15)

where $R_{\text{min}}$ and $R_{\text{max}}$ are taken as $3R_s$ and $100R_s$, respectively.

The BP power output from a disc is equal to the power of disc magnetic braking and can be calculated as (Livio, Ogilvie & Pringle 1999; Lee, Wijers & Brown 2000)

$$L_{\text{BP}} = \int_{R_{\text{min}}}^{R_{\text{max}}} 2\Omega(r)(B_\phi B_z/4\pi)2\pi r dr.$$  

(16)

2.5 Disc winds

2.5.1 Neutrino-heated winds

In this subsection, we discuss the general conditions of a neutrino-driven wind escaping from the disc. In order to escape to large radii, a nucleon has to gain enough energy from the neutrino flux to overcome its gravitational potential at the disc surface. A nucleon with rest mass $m_n$ driven from a thin disc at radius $r$ must absorb a net energy provided by neutrino heating of at least $GMm_n/r$. Since the initial velocity of the nucleon is small, the nucleon is incapable of carrying the amount of energy obtained from the neutrino flux. Almost all of this energy has to go into photon radiation and relativistic electron–positron pairs via electron–neutrino absorption on baryons, i.e. $P \rightarrow \bar{\nu}_e \rightarrow n + e^+$ and $n + e^- \rightarrow P + e^-$, which dominates other forms of neutrino heating (e.g. neutrino-antineutrino annihilation and neutrino-electron scattering which could also produce electron–positron pairs and transfer neutrino energy directly to electrons and positrons). The energy initially stored in photon radiation and electron–positron pairs is converted into the mechanical energy of the nucleons at much larger radii, where the temperatures are low (see QW96; M08). Assuming that the $\nu_e$ and $\bar{\nu}_e$ luminosities and spectra are approximately equal, the neutrino-heating rate through a surface density $\Sigma$ on a narrow annulus $\Delta r$ is

$$\Delta L_{\nu_e} = \frac{\Sigma E_{\nu_e} \Delta L_{\nu_e}}{2\pi r \Delta r^2}.$$  

(17)

where the neutrino-absorption cross-section is $\sigma_{\nu N} \simeq 5 \times 10^{-44} (\bar{\nu}_e)^2 \text{MeV} \text{cm}^2$. $E_{\nu_e}$ is the neutrino energy and $\Delta L_{\nu_e}$ is the neutrino luminosity above the annulus $\Delta r$. A wind that emerges from the disc in the $z$-direction converts its energy to the matter of disc atmosphere at the so-called ‘gain region’ because neutrino
heating overwhelms neutrino cooling at a low temperature. This net heating drives an outflow. The thermal power deposited in the upper disc atmosphere is

\[ E_v = \frac{q^\alpha}{\Sigma} M_{\text{gain}}, \]  

(18)

in which \( M_{\text{gain}} \) is the mass of the atmosphere in the gain region (see M08 for a detailed description). By equating \( E_v \) with bounding energy \( E_b = GM_\alpha M_\alpha/2r(\Delta M_\alpha^s\rho\Omega_1) \) is the mass-loss rate of a wind that originates from the annulus \( r - r + \Delta r \), we get

\[ \Delta M_\alpha^s = \frac{11}{6} \frac{\sigma_\alpha a_\alpha T_\alpha^4}{m_\Sigma(GM)^2} \frac{\Delta L_v}{\Delta r}. \]  

(19)

This expression assumes that the absorbing layer is optically thin for neutrino emission.

2.5.2 Centrifugally driven winds

The high ratio of total energy of outflow to the rest energy of baryons required of ultrarelativistic GRB outflows can be achieved if a strong poloidal magnetic field supplies an energetically dominant Poynting flux. Blandford & Znajek (1977) first proposed that the spin energy of a BH could be extracted by large magnetic fields threading its horizon. Later, Blandford & Payne (1982) showed the possibility that an outflow of matter can be driven centrifugally by large-scale magnetic fields anchored at the surface of the disc. These two mechanisms are usually referred to as BZ and BP processes, respectively. The BZ mechanism, where the outflow follows magnetic field lines anchored to the BH, is probably the best way to prevent baryonic pollution, and can even initially produce a purely leptonic wind. However, for the BP process, the risk of baryonic pollution is much larger since the wind originates from high-density regions (Daigne & Mochkovitch 2002). These winds of different Lorentz factors launched by BZ and BP processes compose a spine/sheath jet structure, which was first pointed out by Meier (2003) from observational reasons and presented by Wang et al. (2008, hereafter W08) to explain the jets for active galactic nuclei and BH binaries. Based on the physical conditions described in Section 2.1, we now adopt the prescription of W08 to calculate the rate of matter outflow driven centrifugally by large-scale magnetic fields anchored at the surface of the disc. The outflow rate from a ring with width \( r - r + dr \) can be expressed as

\[ dM_\text{BP} = \frac{3\alpha P}{r \Sigma_{\Omega_\phi c}} B_\phi^2 R_\phi^2 \left( 1 + \sqrt{1 - a_\alpha^2} \right)^2 S \left( \frac{r}{r_m} \right)^{S-1} \frac{d}{d} \left( \frac{r}{r_m} \right). \]  

(20)

in which we simply assume that \( B_\phi \) is the same in the disc and at the base of the wind. The total outflow rate can be expressed by equation (18) in W08, which could be modified to fit the model as

\[ \dot{M}_{\text{BP}} = \frac{3\alpha P}{r \Sigma_{\Omega_\phi c}} B_\phi^2 R_\phi^2 \left( 1 + \sqrt{1 - a_\alpha^2} \right)^2 \left[ \left( \frac{r_{\text{out}}}{r_m} \right)^S - 1 \right], \]  

(21)

in which \( a_\alpha \) is the BH spin that is moderately taken as 0.5, and \( S \) is an outflow parameter that is approximately taken as 1.01 in calculation. We analytically express the parameter \( a_\alpha \), which is an adjustable parameter due to the uncertainty of equation (13) in W08, as

\[ a_\alpha = \frac{3\alpha P}{\Sigma_{\Omega_\phi c}}. \]

We expect that an average value of the terminal Lorentz factor of the outflow will be given by \( \Gamma = \frac{\dot{\mathcal{E}}}{\dot{\mathcal{E}}_{\text{kin}}}, \) where \( \mathcal{E}_{\text{kin}} \) is the total power injected into the wind and \( M_\alpha \) is the total mass loss from the disc.

3 NUMERICAL RESULT

This section presents our results of the structure of a magnetized neutrino-cooled accretion disc, obtained by numerically solving the full set of equations in Section 2.1. We first show the physical quantities of disc matter as functions of \( r \) in Fig. 1, and in the calculations we fix \( M = 3 M_\odot \), \( M = M_\odot s^{-1} \) and \( \alpha = 0.1 \). Then, based on those numerical results of physical quantities of the disc matter, i.e. the disc pressure \( P \), the density \( \rho \), the temperature \( T \) and the magnitude of magnetic fields \( B \), we calculate the composition of the disc matter, the neutrino luminosity and BP power, and discuss the main properties of the disc wind driven by neutrino emission and launched centrifugally by BP process, respectively. We show each component of pressure in Fig. 1(a). The magnetic pressure component is shown by the red line; the gas pressure, degeneracy pressure, radiation pressure and neutrino pressure, and total pressure are shown by the solid, dotted, dashed, long-dashed and blue lines, respectively. It is seen that gas pressure dominates from \( 5R_\phi \) outwards to \( ~50R_\phi \) and the magnetic pressure is about 15 per cent of total pressure of the disc. However, beyond this region (i.e. \( r \geq 50R_\phi \)), radiation pressure begins to dominate, and the disc may become thermally unstable, since \( T^4 \propto P \propto \Sigma H \) and \( Q \propto \Sigma T^{3/2} \), and the criterion for thermal stability (see e.g. DPN) is not satisfied. However, the recent work of Hirose, Krolik & Blaes (2008) shows that stress fluctuations precede pressure fluctuations, contrary to the usual supposition that the pressure controls the saturation level of the magnetic energy; thus the radiation-dominated regions can also be stable and no thermal runaway occurs over a time-span of \( ~40 \) cooling times. In Figs 1(b) and (c), we can see that from \( 1000R_\phi \) inwards to \( 5R_\phi \), \( \rho \) increases by about four orders of magnitude and reaches \( \sim 10^{10} \text{ g cm}^{-3} \) in the innermost region of the disc. \( T \) increases by about two orders of magnitude and reaches \( \sim 8 \times 10^{10} \) K in the innermost region. In Fig. 1(d), we show \( B_\phi \) and \( B_\phi \) with dashed, solid and dotted lines, respectively. We can see that \( B_\phi \) increases by about two orders and reaches \( 3.5 \times 10^{15} G \) at \( 5R_\phi \), which agrees well with the magnetohydrodynamic simulations of merger process (Price & Rosswog 2006) that the fields are amplified by Kelvin–Helmholz instabilities and \( B_\phi \sim 0.71 B_\phi \sim 0.22 B_\phi \).

Fig. 2 shows the composition of disc matter. We fix the electron fraction at 0.1 (upper panel, note that Chen et al. 2007 show that all \( \nu \)-cooled discs are very neutron rich in the inner region with \( Y_e \sim 0.1 \) and 0.3 (bottom panel). In the outer region between \( 1000R_\phi \) and \( \sim 150R_\phi \), almost no particles have disintegrated, so the \( \alpha \)-particle fraction \( X_\alpha \) remains \( \sim 0.2 \) and \( \sim 0.6 \), respectively. From \( \sim 150R_\phi \) inwards, the disintegration of \( \alpha \)-particles causes \( X_\alpha \) to decrease and \( X_\beta \) to increase dramatically.

We show our results for the Alfven velocity \( v_A \) and radial velocity of disc \( v_r \) (solid and dashed lines, respectively) in the upper panel of Fig. 3, and diffusing time \( t_d \) and viscous time \( t_v \) (solid and dashed lines, respectively) in the bottom panel of Fig. 3. One can see that though \( v_A > v_r \), the diffusing time \( t_d \) is still longer than the viscous time \( t_v \), since the half-thickness \( H \) is larger than radius \( r \) after the magnetic fields were included (see Fig. 6). As the process of magnetic fields diffusing is slower than the process of accumulating in the inner region of the disc, we get the conclusion that the magnetic barrier may be formed via the magnetic flux accumulating in the vicinity of the BH. If the barrier is formed between the BH horizon and \( ~3R_\phi \), the balance between the pressure of magnetic fields and the ram pressure of the innermost parts of an accretion flow should be fulfilled, i.e. \( B_\phi^2/8\pi = p_{\text{ram}} \sim \rho c^2 \), which implies \( B_\phi \sim 2 \times 10^{20} G \). The magnetic flux of the barrier is \( \Phi_\theta \sim (3R_\phi)^2 B_\phi \approx 2.8 \times 10^{29} \text{ cm}^{-2} \) (these results agree with
Figure 1. (a) Pressure components for $\dot{m} = 1$. The total pressure is shown by the blue line, magnetic pressure by the red line, gas pressure by the solid line, degeneracy pressure by the dotted line, radiation pressure by the dashed line and the neutrino pressure by the long-dashed line; (b) the density of disc matter; (c) the disc temperature; (d) the strength of magnetic fields in the disc. $B_\phi$, $B_Z$ and $B_r$ are given by dashed, solid and dotted lines, respectively.

Figure 2. Proton fraction $X_P$ (solid line), neutron fraction $X_n$ (dashed line) and $\alpha$-particles fraction $X_\alpha$ (dotted line) as functions of $r$, with the fixed electron fraction $Y_e$ at 0.1 (upper panel) and 0.3 (bottom panel), respectively.

Figure 3. The radial velocity (dotted line) and Alfvén velocity (solid line) are shown in the upper panel, and the viscous time (dotted line) and buoyant time (solid line) are shown in the bottom panel.
The mass-loss rate driven by neutrino emission (upper panel) and large-scale magnetic fields (bottom panel) as functions of r, for $m = 1$ (dashed line), $m = 1$ (solid line) and $m = 10$ (dotted line), respectively.

Table 1. The total mass-loss rate $M_w$ and terminal Lorentz factors $\Gamma$ of outflows for different accretion rates $m$.

<table>
<thead>
<tr>
<th>$m$</th>
<th>$M_w (M_\odot \text{s}^{-1})$</th>
<th>$\Gamma$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.1</td>
<td>$1.88 \times 10^{-4}$</td>
<td>2.83</td>
</tr>
<tr>
<td>1</td>
<td>$5.92 \times 10^{-5}$</td>
<td>760.6</td>
</tr>
<tr>
<td>10</td>
<td>$1.03 \times 10^{-3}$</td>
<td>105.8</td>
</tr>
<tr>
<td>0.1</td>
<td>$2.17 \times 10^{-4}$</td>
<td>7.29</td>
</tr>
<tr>
<td>1</td>
<td>$2.29 \times 10^{-3}$</td>
<td>7.65</td>
</tr>
<tr>
<td>10</td>
<td>0.016</td>
<td>1.49</td>
</tr>
</tbody>
</table>

Figure 4. Neutrino luminosity and BP power. The solid line represents the neutrino luminosity of the model including magnetic fields, and the dashed line represents the BP power; the dotted line representing the neutrino luminosity do not include the influence of magnetic fields.

Figure 5. The mass-loss rate driven by neutrino emission (upper panel) and large-scale magnetic fields (bottom panel) as functions of r, for $m = 1$ (dashed line), $m = 1$ (solid line) and $m = 10$ (dotted line), respectively.
Moreover when \( v_{\alpha\phi} = (2c_{s}v_{K})^{1/2} \), the growth rate MRI vanishes. Based on the above discussion, if we assume that the growth of the azimuthal field, \( B_{\phi} \), stops being amplified and saturates when

\[
v_{\alpha\phi} \sim (2c_{s}v_{K})^{1/2}
\]

under this prescription, the results we get are almost the same as the present paper except that \( B_{i} < B_{z} \).

### 4 CONCLUSIONS

In this paper, we present a magnetized neutrino-dominated accretion disc as the central engine of GRBs. We found that the pressure of magnetic fields contributes about 15 per cent of the total pressure of the accretion flows, and the strength of magnetic fields can reach \( 3.5 \times 10^{15} \) \( G \) in the innermost region of the disc for \( \dot{m} = 1 \). The disintegration of \( \alpha \)-particles mainly occurs in the region from \( \sim 150 R_{g} \) inwards to \( \sim 80 R_{g} \). Since the diffusing time-scale is longer than viscous time-scale, a magnetic barrier between the BH and the accretion flow would likely be formed. We also show that after magnetic fields were considered, the neutrino luminosity drops to \( 1.0 \times 10^{48} \) erg s\(^{-1} \) for \( \dot{m} = 0.01 \), which is insufficient to fuel some of the early X-ray flares; however, the BP power is about \( 1.4 \times 10^{49} \) erg s\(^{-1} \), robust enough for the most powerful X-ray flare observed. Thus, the prompt emissions may be powered by both neutrino processes and magnetic processes; however, X-ray flares (or the late activities of the central engine) are mainly caused by the magnetic mechanism. It is also shown that the distribution of neutrino-driven mass-loss rate has a bimodal shape along the disc radius, because both neutrino-emission flux and gravitational potential drop along the radius; however, the magnetic fields driven concentrate in the inner region where it holds the greatest strength of magnetic fields. The outflows of different origins have different Lorentz factors that is convenient for some internal shock models to fit the light curves of GRBs.

### ACKNOWLEDGMENTS

This work has been supported by the National Natural Science Foundation of China under Grant Number 10703002 and the National Science Foundation of Shanxi Normal University under Grant Number YZ08021. The anonymous referee is thanked for his/her helpful comments.

### REFERENCES


Cambridge Univ. Press, Cambridge


Hjorth J. et al., 2003, Nat, 423, 847


Price D. J., Rosswog S., 2006, Sci, 312, 719


### APPENDIX A: DERIVATION OF EQUATION (5)

The surface current density of disc at radius \( r \) is defined as \( J_{r} \), which is proportional to the azimuthal component of the magnetic fields \( B_{\phi} \),

\[
4\pi J_{r}/2 = -B_{\phi}, \tag{A1}
\]

where 2 in the denominator comes from the fact that the radial current density \( J_{r} \) includes the currents into the disc both from above and below. The torque exerted by the annular ring with width \( dr \) of the disc due to the Lorentz force is given by

\[
dT = -2\pi r J_{r} B_{\phi} dr. \tag{A2}
\]
In the case without the process of magnetic braking (i.e. BP process), accretion is produced by the internal viscous torque $T_{vis}$ of the disc. Then, the conservation law of angular momentum is given by (Frank, King & Raine 1992):

$$-\dot{M}_D \frac{\partial (r^2 \Omega_D)}{\partial r} = \frac{\partial T_{vis}}{\partial r},$$

(A3)

in which $\frac{\partial T_{vis}}{\partial r}$ is the contribution due to internal viscous torque, always transporting angular momentum outwards in the disc. Taking magnetic braking effects into account, we think equation (A3) should be modified by

$$-\dot{M}_D \frac{\partial (r^2 \Omega_D)}{\partial r} = \frac{\partial T_{vis}}{\partial r} + \frac{\partial T_{BP}}{\partial r},$$

(A4)

where $\frac{\partial T_{BP}}{\partial r}$ is the contribution due to BP torque. Comparing equations (A3) and (A4), we express $(\dot{M}_D)_{BP}$, the contribution of magnetic braking effect to the accretion rate, as follows:

$$(\dot{M}_D)_{BP} = \frac{\partial T_{BP}}{\partial r} \frac{\partial (r^2 \Omega_D)}{\partial r}.$$  

(A5)

Substituting equations (A1), (A2) into equation (A5), we have

$$(\dot{M}_D)_{BP} = \frac{4(B_B B_z/4\pi)2\pi}{\Omega_D}.$$  

(A6)

This is the second term of equation (5).