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Magnetically confined wind shocks in X-rays – A review

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Abstract

A subset (∼10%) of massive stars present strong, globally ordered (mostly dipolar) magnetic fields. The trapping and channeling of their stellar winds in closed magnetic loops leads to *magnetically confined wind shocks* (MCWS), with pre-shock flow speeds that are some fraction of the wind terminal speed. These shocks generate hot plasma, a source of X-rays. In the last decade, several developments took place, notably the determination of the hot plasma properties for a large sample of objects using *XMM* and *Chandra*, as well as fully self-consistent MHD modeling and the identification of shock retreat effects in weak winds. Despite a few exceptions, the combination of magnetic confinement, shock retreat and rotation effects seems to be able to account for X-ray emission in massive OB stars. Here we review these new observational and theoretical aspects of this X-ray emission and envisage some perspectives for the next generation of X-ray observatories.

Keywords: X-rays; MHD; Massive stars

1. Introduction

Hot luminous, massive stars of spectral type O and B are prominent sources of X-rays which can originate from three distinct sources: shocks in their high-speed radiatively driven stellar winds, wind-wind collisions in binary systems and magnetically confined wind shocks.

In single, non-magnetic O stars, the intrinsic instability of wind driving by line-scattering leads to embedded wind shocks that are thought to be the source of their relatively soft (∼0.5 keV) X-ray spectrum, with a total X-ray luminosity that scales with stellar bolometric luminosity, $L_x \sim 10^{-7} \times L_{bol}$ (Chlebowski, 1989; Berghoefer et al., 1997; Nazé et al., 2011). In massive binary systems the collision of the two stellar winds at up to the wind terminal speeds can lead to even higher L_x, generally with a significantly harder (up to 10 keV) spectrum (for a review, see Rauw & Nazé, 2016).

Here we discuss a third source of X-rays from OB winds, namely those observed from the subset (∼10%) of massive stars with strong, globally ordered (often significantly dipolar) magnetic fields (Petit et al., 2013); in this case, the trapping and channeling of the stellar wind in closed magnetic loops leads to *magnetically confined wind shocks* (MCWS) (Babel and Montmerle, 1997a,b, hereafter BM97a,b), with pre-shock flow speeds that are some fraction of the wind terminal speed, resulting in intermediate energies for the shocks and associated X-rays (∼2 keV). A prototypical example is provided by the magnetic O-type star θ^1 Ori C, which shows moderately hard X-ray emission with a rotational phase variation that matches well the expectations of the MCWS paradigm (Gagné et al., 2005).

Here, we discuss theoretical aspects of magnetic confinement that determine the extent of the influence of the field over the wind. We, then, describe an effect called 'shock

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retreat', which can moderate the strength of the X-rays, or even quench it altogether in extremely low mass loss rate stars, and the effects of rotation. In Section 3, we describe current results in the X-ray observations of early-type magnetic stars, while Section 4 briefly discusses future outlooks of MCWS in X-ray astronomy.

2. Theoretical perspective

To explain X-ray emission from the Ap-Bp star IQ Aur (Babel and Montmerle, 1997a) introduced MCWS model. In their approach, they effectively prescribed a fixed magnetic field geometry to channel the wind outflow (see also Shore and Brown, 1990). For large magnetic loops, wind material from opposite footpoints is accelerated to a substantial fraction of the wind terminal speed (i.e., $\geqslant 1000$ km s^{-1}) before the channeling toward the loop tops forces a collision with very strong shocks, thereby heating the gas to temperatures (10^7–10^8 K) that are high enough to emit hard (few keV) X-rays. This star has a quite strong field ($\sim$4 kG) and a rather weak wind, with an estimated mass loss rate of about $\sim 10^{-10} M_\odot$ yr^{-1}, and thus indeed could be reasonably modeled within the framework of prescribed magnetic field geometry. Later, BM97b applied this model to explain the periodic variation of X-ray emission of the O7 star θ^1 Ori C, which has a lower magnetic field ($\sim$1100 G) and significantly stronger wind (mass-loss rate $\sim 10^{-7} M_\odot$ yr^{-1}), raising now the possibility that the wind itself could influence the field geometry in a way that is not considered in the simple fixed-field approach.

2.1. Magnetic confinement

In an interplay between magnetic field and stellar wind, the dominance of the field is determined how strong it is relative to the wind. To understand the competition between these two, ud-Doula and Owocki (2002) defined a characteristic parameter for the relative effectiveness of the magnetic fields in confining and/or channeling the wind outflow. Specifically, consider the ratio between the energy densities of field vs. flow,

$$\eta(r,\theta) \equiv \frac{B^2/8\pi}{\rho v^2/2} \approx \frac{B^2 r^2}{\dot{M} v(r)} = \left[\frac{B_*^2(\theta) R_*^2}{\dot{M} v_\infty}\right] \left[\frac{(r/R_*)^{-2n}}{1 - R_*/r}\right], \quad (1)$$

where the latitudinal variation of the surface field has the dipole form given by $B_*^2(\theta) = B_o^2(\cos^2\theta + \sin^2\theta/4)$. In general, a magnetically channeled outflow will have a complex flow geometry, but for convenience, the second equality in Eq. (1) simply characterizes the wind strength in terms of a spherically symmetric mass loss rate $\dot{M} = 4\pi r^2 \rho v$. The third equality likewise characterizes the radial variation of outflow velocity in terms of the phenomenological velocity law $v(r) = v_\infty(1 - R_*/r)^\beta$, with v_∞ the wind terminal speed and $\beta = 1$; this equation furthermore models the magnetic field strength decline as a power-law in radius, $B(r) = B_*(R_*/r)^{(n+1)}$, where, e.g., for a dipole $n = 2$.

With the spatial variations of this energy ratio thus isolated within the right square bracket, we see that the left square bracket represents a dimensionless constant that characterizes the overall relative strength of field vs. wind. Evaluating this in the region of the magnetic equator ($\theta = 90^o$), where the tendency toward a radial wind outflow is in most direct competition with the tendency for a horizontal orientation of the field, one can thus define an equatorial 'wind magnetic confinement parameter',

$$\eta_* \equiv \frac{B_*^2(90^\circ) R_*^2}{\dot{M} v_\infty} = 0.4 \, \frac{B_{100}^2 \, R_{12}^2}{\dot{M}_{-6} \, v_8}, \quad (2)$$

where $\dot{M}_{-6} \equiv \dot{M}/(10^{-6}\, M_\odot/\text{yr})$, $B_{100} \equiv B_o/(100 \text{ G})$, $R_{12} \equiv R_*/(10^{12} \text{ cm})$, and $v_8 \equiv v_\infty/(10^8 \text{ cm/s})$. In order to have any confinement, $\eta_* \geqslant 1$. As these stellar and wind parameters are scaled to typical values for an OB supergiant, e.g. ζ Pup, the last equality in Eq. (2) immediately suggests that for such winds, significant magnetic confinement or channeling should require fields of order few hundred G. By contrast, in the case of the sun, the much weaker mass loss ($\dot{M}_\odot \sim 10^{-14} M_\odot/\text{yr}$) means that even a much weaker global field ($B_o \sim 1$ G) is sufficient to yield $\eta_* \simeq 40$, implying a substantial magnetic confinement of the solar coronal expansion. But in Bp stars the magnetic field strength can be of order kG with $\dot{M}_\odot \sim 10^{-10} M_\odot/\text{yr}$ leading $\eta_* \leqslant 10^6$. Thus, the confinement in Bp stars is very extreme.

It should be emphasized that $\dot{M}$ used in the above formalism is value obtained for a spherically symmetric non-magnetic wind as the magnetic field may significantly influence the predicted circumstellar density and velocity structure.

2.1.1. Alfvén radius

The extent of the effectiveness of magnetic confinement is set by the Alfvén radius, R_A, where flow and Alfvén velocities are equal. This will also determine the extent of the largest loops and thus the highest shock velocities affecting the hardness of X-ray emission. This radius can be derived from Eq. (1) where the second square bracket factor shows the overall radial variation; n is the power-law exponent for radial decline of the assumed stellar field, e.g. $n = 2$ for a pure dipole, and with $v(r) = v_\infty(1 - R_*/r)^\beta$ where β is the velocity-law index, with typically $\beta \approx 1$. For a star with a non-zero field, we have $\eta_* > 0$, and so given the vanishing of the flow speed at the atmospheric wind base, this energy ratio always starts as a large number near the stellar surface, $\eta(r \to R_*) \to \infty$. But from there outward it declines quite steeply, asymptotically as r^{-4} for a dipole, crossing unity at the Alfvén radius defined implicitly by $\eta(R_A) \equiv 1$.