The nature of SS433 and the ultraluminous X-ray sources

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ABSTRACT

The periodic precession (162-d) and nodding (6.3-d) motions of the jets in SS433 are driven in the outer regions of the disc, whereas the jets themselves, being relativistic, are launched near the black hole at the disc centre. Given that the nutation period is comparable to the dynamical time-scale in the outer regions of the disc, it seems unlikely that these periods can be communicated efficiently to the disc centre. Here, we propose that the massive outflow observed in SS433 is launched at large radii in the disc, about 1/10 of the outer disc edge, and that it is this outflow which responds to the oscillations of the outer disc and determines the direction of the jets. The massive outflow is launched at large radius because the mass transfer rate is hyper-Eddington. This implies not only that the total luminosity of SS433 exceeds \( L_{\text{Edd}} \) by a considerable factor, but also that the radiative output is collimated along the outflow. We thus suggest that SS433 is an ultraluminous X-ray source (ULX) viewed ‘from the side’. We also suggest that the obscured INTEGRAL sources may be SS433-like objects, but with slightly lower mass transfer rates.

Key words: accretion, accretion discs – black hole physics – X-rays: binaries.

1 INTRODUCTION

The extraordinary properties of SS433 have long attracted attention. This 13.1-d binary system shows a pair of well-collimated jets with velocity \( \pm 26c \) which precess on a cone of half-angle \( \theta \approx 20^\circ \) with period \( P_{\text{pr}} = 162 \text{ d} \) [Abell & Margon 1979; Fabian & Rees 1979; Margon 1979; Milgrom 1979; Margin 1984; see Fabrika (2006) for a recent review of observational properties]. In addition to the jets, there is a powerful outflow with velocity \( v_o \approx 1500 \text{ km s}^{-1} \) which may have inflated the surrounding nebula W50 (Begelman et al. 1980). The jets have mass-loss rate \( \dot{M} \lesssim 5 \times 10^{-7} \text{ M}_\odot \text{ yr}^{-1} \) and kinetic luminosity \( L_k \gtrsim 10^{39} \text{ erg s}^{-1} \) (e.g. Begelman et al. 1980; Königl 1983; Brinkmann & Kawai 2000; Marshall, Canizares & Schultz 2002) and interact in a helical pattern with W50 (Hjellming & Johnston 1981; Blundell & Bowler 2004).

The very regular precession of the jets is probably related to other super-orbital periods seen in X-ray binaries. Pringle (1996, 1997) showed that strong self-irradiation causes an accretion disc to warp out of the orbital plane and to precess. (Wijers & Pringle 1999, see also Ogilvie & Dubus 2001) showed that such discs can exhibit precessions with the long periods seen in some X-ray binaries, including SS433. If the disc is sufficiently warped that the mass transfer stream does not strike the outer disc edge, then the disc shape is such that the radiatively driven precession is retrograde (see fig. 7 of Wijers & Pringle 1999). However, such 162-d, radiatively driven, precessional warping occurs predominantly in the outer parts of the disc, whereas the jets, since they are relativistic, are presumably driven from very close to the compact accretor at its centre (assumed to be a black hole). Thus, in this model, it is necessary for the precession of the outer disc to communicate itself to the centre of the disc where the initial jet direction is determined. Indeed, the problem is worse than this for two reasons. First, as seems likely, the black hole is spinning, then the inner disc is aligned with the spin of the hole (Bardeen & Petterson 1975) independent of the spin of the outer disc. Secondly, the precessing jets also show a nodding motion (nutation) with period close to one-half of the 13.1-d orbital period \( P_{\text{orb}} \) (actually the synod period \( 2/P_{\text{orb}} + 1/P_{\text{pr}} \) as the disc precession is retrograde). Disc nodding with this period is a direct effect of the \( m = 2 \) part of the tidal torque (Katz et al. 1982; Bate et al. 2000) acting on the outer part of the disc. The amplitude of such a short period oscillation would be reduced even more severely (Katz 1986) during inwards propagation than that of the 162-d precession, simply because diffusion damps rapid oscillations more strongly than slower ones. Hence, there seems little hope of communicating the nodding motion to the inner disc, where one might expect the jet directions to be fixed (Bate et al. 2000).

In this paper, we put forward a solution to this problem. We propose that the powerful 1500 km s\(^{-1}\) outflow comes about because the mass transfer rate is hyper-Eddington (\( \sim 5000 M_{\odot} \text{ yr}^{-1} \)), and that the outflow is therefore radiatively driven from well outside the Schwarzschild radius, i.e. \( R_{\text{outflow}} \gg R_s \) (Shakura & Sunyaev 1973).
(By $M_{\text{Edd}}$ we mean the rate at which spherical accretion on to a black hole with radiative efficiency $\eta \approx 0.1$ produces the Eddington luminosity.) The direction of the outflow is determined by the tilt of the disc at this large radius and is subject to both the precessional and nutational motions of the outer disc. The jets, however, are still launched from close to the accretor with initial direction fixed by the spin of the hole, but with final direction aligned with the powerful (precessing and nutating) outflow.

Such a high mass transfer rate is to be expected if SS433 is a direct descendant of a high-mass X-ray binary (HMXB) in which the more massive companion star now fills its Roche lobe and is transferring mass on a thermal time-scale (King, Taam & Begelman 2000). A consequence of this model is that SS433 would look far brighter if viewed along the axis of the outflow, and we suggest that in this case it might resemble an ultraluminous X-ray source (ULX).

In Section 2, we consider the structure of the disc in SS433. We show that at these high accretion rates, disc warp is communicated by wave-like motions rather than by diffusion (viscous torques). In Section 3, we discuss the nature and magnitude of the powerful outflow, and also argue that at such high accretion rates the bolometric luminosity of SS433 can considerably exceed $L_{\text{Edd}}$. In Section 4, we show that the outflow is massive enough to deflect the jets and align them with the axis of the outflow. In Section 5, we argue that SS433 would resemble a ULX, if observed along the axis of the outflow. We thus arrive at a picture of the accretion disc in SS433 as tilted and aligned with the binary angular momentum, and an outer disc warped significantly out of the orbital plane by radiation torques. In this situation, much of the mass transferred by the companion joins the disc near the circularization radius (Flannery 1975), $R_0 \simeq 0.07 a = 5.2 \times 10^{10}$ cm, rather than near its edge (Paczynski 1977) $R_{\text{edge}} \simeq 0.22 a \simeq 3 R_0 \simeq 1.6 \times 10^{12}$ cm. This then leads to a central disc aligned with the binary angular momentum, and an outer disc tilting to large angles (see fig. 7 of Wijers & Pringle 1999). The latter is still roughly as massive as a disc fed at its outer edge (Wijers & Pringle 1999) because it acts as the reservoir for the angular momentum lost by the matter accreting through the inner disc, before passing this back to the companion star via tides (as in all discs in binary systems). Thus, its surface density $\Sigma$ must match on to that of the accreting inner disc at the interface. Since there is little accretion through the outer disc, the usual disc diffusion equation implies that $\nu \Sigma \propto R^{-1/2}$ rather than $\nu \Sigma \simeq \text{constant}$. Hence, the total mass and angular momentum in this outer non-accreting disc is only slightly less than in a normal accretion disc.

The large expected tilt angle of the outer disc may be the reason for the otherwise puzzling deduction (Stewart et al. 1987) that the disc partially blocks the observer’s view of the central X-ray jet at certain precession phases. Stewart et al. interpreted their result in terms of a thick disc with aspect ratio $H/R \sim 0.4$. In order to cause the same obscuration, a thin tilted disc would have to lie at an angle $\tan^{-1} 0.4 \simeq 20^\circ$ to the orbital plane. We note that this angle is remarkably close to the half angle $\theta$ of the jet precession cone, suggesting that the jet somehow becomes aligned with the axis of the tilted disc. This again brings up the question of communicating the tilt inwards.

If the 6.3-d wobble were propagated inwards by viscosity, then it would be damped out over a distance such that the viscous time-scale from the outer disc is of the order of 6.3 d. Since the dynamical time-scale at the outer disc edge is only a few days, and since the viscous time-scale is longer than the orbital time-scale by a factor of $\sim(R/H)^2 \alpha^{-1} \gg 1$, propagation of the nutation period could not proceed very far. However, in a nearly Keplerian disc, tilt is communicated by bending waves, rather than viscosity, provided that

$$\frac{H}{R} \gtrsim \max(\alpha, |\Omega - \kappa|/\Omega),$$

where $\alpha$ is the dimensionless viscosity parameter, $\Omega$ the angular frequency of the disc and $\kappa$ the local epicyclic frequency (see the discussion in Wijers & Pringle 1999). For the parameters of SS433, we find, using the formulae given by Wijers & Pringle (1999), that at the outer disc edge $|\Omega - \kappa|/\Omega \simeq 0.03$. In addition, for an accretion rate of $M \sim 10^{-5}$ to $10^{-3}$ $M_{\odot}$ yr$^{-1}$, we find that in the outer disc $H/R \simeq 0.1$ (Shakura & Sunyaev 1973). Thus, for typical values of $\alpha$ in the range 0.01–0.1, we see that warp propagation can occur via bending waves in SS433.1 How far inwards the nutational wobble and the precessional warp can propagate is not easy to determine. The warp waves are damped by viscosity, and also perhaps by internal hydrodynamical instabilities (see the discussion in Bate et al. 2000). But the wave amplitude can also grow, by conservation of wave action, depending on the density profile of the disc.

It is unlikely, however, that the tilt can propagate inwards as far as regions close to the central black hole. If the hole has non-zero spin, then the disc near the black hole must be aligned with its spin axis out to a warp radius of a few tens of Schwarzschild radii $R_{s} = 3 \times 10^6$ cm (cf. King et al. 2005). If we adopt the usual assumption that the black hole has its spin aligned with the orbital plane, then since the inner disc is the region from which the jets must be launched, the jets initially move along the black hole (and orbital) spin axis. We thus arrive at a picture of the accretion disc in SS433 as tilted at angles $\theta \simeq 20^\circ$ at large radii, and aligned with the orbital plane and black hole spin close in. As we have remarked, bending waves are able to propagate the disc tilt and its nodding motion inwards to parts of the disc gaining mass from the companion to some extent, but not to the innermost parts where the jets are launched. We thus need some other agency to bend the jets to align them with the axis of the outer disc, and thus to communicate the disc precession and nodding to them.

3 THE OUTFLOW FROM SS433

The ultimate cause of the unique features of SS433 is probably its extremely high mass transfer rate, which far exceeds the Eddington rate for a $10 M_{\odot}$ black hole accretor. Only a small part of the transferred mass is lost in the jets ($M_j \simeq 5 \times 10^{-7} M_{\odot}$ yr$^{-1}$). Evidently,

1Since the analysis of Wijers & Pringle (1999) assumed that disc tilt is communicated only through viscosity, strictly speaking it is necessary for that analysis to be reworked for the case of SS433. We will assume here, however, that the basic conclusions still stand.

2This appears not to be the case for some of the microquasars, but for SS433 the accretion rate is so high that alignment may have had time to occur (Maccarone 2002).

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most of it is expelled as the massive high-speed outflow \( \left( v_w \simeq 1500 \text{ km s}^{-1} \right) \) seen in the form of the 'stationary' H\( \alpha \) line and broad absorption lines. This velocity suggests that the outflow is expelled from a radius (van den Heuvel 1981; Seifina et al. 1991)

\[
R_{\text{outflow}} \simeq \frac{2GM}{v_w^2} \simeq \left( \frac{c^2}{v_w^2} \right) R_s \simeq 1.2 \times 10^{11} \text{ cm}. \tag{2}
\]

Support is given to these ideas by the numerical simulations of Okuda et al. (2005). Because accretion is highly super-Eddington, we expect \( R_{\text{outflow}} \) to correspond to the 'spherization radius' given by (Shakura & Sunyaev 1973; from their equation 7.1) as

\[
R_{\text{out}} = \frac{27M_{\text{out}}}{4M_{\text{Edd}}} R_s \tag{3}
\]

(this is close to the trapping radius). Hence for this picture to be consistent, we require

\[
\dot{M}_{\text{out}} = \frac{4c^2}{27v_w^4} M_{\text{Edd}} = 5000 M_{\text{Edd}} \simeq 5 \times 10^{-4} M_\odot \text{ yr}^{-1}. \tag{4}
\]

This is consistent with evolutionary calculations of the mass transfer rate \( \dot{M}_2 \simeq \dot{M}_{\text{out}} \) (King et al. 2000), and would suggest a mass transfer lifetime \( t_{\text{ov}} \simeq 4 \times 10^6 \text{ yr} \) for a companion of mass \( M_2 \simeq 20 M_\odot \). It is also consistent with estimates of the gas flow rate from SS433 (van den Heuvel 1981; Shklovskii 1981; Blundell et al. 2001), and agrees with the estimate of \( \dot{M}_{\text{out}} \) by King et al. (2000), which used the emission measure of the stationary H\( \alpha \) and associated infrared free–free radiation. Things would change very little with other assumed masses (e.g. \( M_1 = 2.9 M_\odot \), \( M_2 = 10.9 M_\odot \)) Hillwig et al. 2004, who combine absorption-line velocities from the companion with antiphased emission-line velocities from the vicinity of the accretor). The main change is that the smaller accretor mass would give a smaller Eddington limit and outflow rate (both by a factor of ~3). These are still compatible with the evolutionary calculations of King et al. (2000), and our conclusions would remain qualitatively unchanged.

### 3.1 The Limiting Luminosity of SS433

While most of the outflow from SS433 is expelled from \( R_{\text{outflow}} \), we expect further outflow from the disc inside this radius. If each disc radius is close to its local Eddington limit, the outflow must be arranged so that the accretion rate at disc radius is close to its local Eddington limit, the outflow must be arranged so that the accretion rate at disc radius is close to its local Eddington limit, the outflow must be arranged so that the accretion rate at disc radius is close to its local Eddington limit, the outflow must be arranged so that the accretion rate at disc radius is close to its local Eddington limit.

\[
\frac{\dot{M}_{\text{out}}}{2\pi R_{\text{outflow}} v_w} \simeq M_\odot \text{ yr}^{-1} \tag{5}
\]

The jets hit this wall at angle \( \theta \). We now check if this causes significant outwards deviation (i.e. along a cylindrical radius) in the outflow. In order for the jet to be deflected sufficiently, the jet impact must cause the flow in the wall to deflect through an angle \( \leq \text{sin} \theta \). Thus, we require the wall flow to acquire a radial component of velocity \( \Delta v \ll v_w \sin \theta \). The jet momentum flux density normal to the wall is

\[
\rho_j v_j^2 \sin \theta \frac{1}{\text{sin} \theta}. \tag{6}
\]

where \( \sin \theta \) in the numerator gives the radial component (in cylindrical coordinates), and \( l/\text{sin} \theta \) in the denominator comes from the fact that this is spread over a larger wall area because of the projection. An element of the outflow passing through the region where the jet hits feels the jet’s momentum flux for a time

\[
\Delta t = \frac{l}{v_w l/\text{sin} \theta}. \tag{7}
\]

where \( l \) is the diameter of the jet as it hits the wall. The jet collision thus imparts sideways velocity

\[
\Delta v = \frac{\rho_j v_j^2 \sin^2 \theta l}{\text{sin} \theta \Sigma_{\text{outflow}}} \frac{1}{v_w} \tag{8}
\]

If the full opening angle of the jet is \( \theta_j \), we have \( l = R \sin \theta_j \), where \( R = R_{\text{outflow}}/\sin \theta_j \). Mass conservation requires that

\[
\dot{M}_j = \frac{\pi}{2} \rho_j v_j l^2 \tag{9}
\]

where we have assumed that \( \frac{1}{2} \dot{M}_j \) is the mass flux in each jet. We thus find

\[
\Delta v \frac{v_w}{v_j} = \frac{2}{\pi} \frac{\dot{M}_j}{M_{\text{out}}} \frac{l}{v_j} \frac{\sin \theta}{\text{sin} \theta_j} \frac{2\pi R_{\text{outflow}} v_w}{\dot{M}_{\text{out}}} \tag{10}
\]

which reduces to

\[
\Delta v \frac{v_w}{v_j} = 4 \frac{\dot{M}_j}{M_{\text{out}} v_w \sin \theta_j} \sin \theta \frac{l}{R_{\text{outflow}}} \tag{11}
\]

With \( \dot{M}_j \sim M_{\text{Edd}} v_j \sim c \) this is

\[
\Delta v \frac{v_w}{v_j} \leq 4 \left( \frac{R_s}{R_{\text{outflow}}} \right)^{1/2} \frac{\sin \theta}{\sin \theta_j} \tag{12}
\]

With \( (R_{\text{outflow}}/R_s)^{1/2} = 200, \theta \sim 20^\circ, \theta_j = 3^\circ \) (Fabrika 2006) gives

\[
\Delta v / v_w < 0.045 \ll \sin \theta \sim 0.34.
\]
We conclude that the outflow is well capable of bending the jet direction until it is parallel to the outflow axis, and thus that the outflow can communicate the precession and nodding motions to the jet direction. If the jet deflection takes place through one oblique shock, then in this simple picture the flow is deflected through an angle $\chi = \theta = 20^\circ$, and the post-shock velocity is the jet velocity of $v_\| \approx 0.26c$. We can then make use of (the special-relativistic generalization of) the standard shock theorem (cf. Landau & Lifschitz 1959, equation 86.5) to compute the properties of the shock. The main result of this, assuming that the shock is very oblique so that the post-shock velocity is supersonic, is that only about $\sim 1$ per cent of the pre-shock kinetic energy, i.e. $\sim 10^{37}$ erg s$^{-1}$, is dissipated in the shock.\(^3\) It is more realistic to suppose that the jets are deflected through a series of oblique shocks. In this case, the amount of dissipation of jet energy is lower, and it is now distributed along the jet. Thus, the amount of energy dissipated in deflecting the jets makes little contribution to the overall energy output of the system, and does not slow the jets significantly. These bending shocks, and concomitant radiative losses, are hidden from the observer within the outflow (see below). The observed jet X-ray emission (Watson et al. 1986; Stewart et al. 1987) comes from further out, presumably from internal shocks, and is observable since the outflow has spread geometrically at such distances. The only way in which the energy dissipated in the jet-deflecting shocks might be observable is if the post-shock temperatures ($\sim 10^{11}$ K) give rise to high-energy (e.g. MeV) emission, and if the system were viewed down the axis of the outflow.

Evidently, the pre-shock jet is much slower than those with Lorentz factors $\gamma \gtrsim 10$ inferred in sub-Eddington, or near-Eddington ($M \lesssim M_{\text{Edd}}$), accreting systems (Miller-Jones, Fender & Nakar 2006). It is not clear why this is so, but it may indicate that because in SS433 the accretion is strongly super-Eddington ($M \gtrsim M_{\text{Edd}}$), the jets in this case are accelerated by radiation pressure, rather than by magnetic processes [cf. the outflows computed by Okuda et al. (2005)].

Our picture makes a further prediction. We would not expect the tilt angle $\theta$ of the outer disc and outflow to be precisely constant in time, and indeed small variations are seen. If the bending shocks are sufficiently radiative – which seems likely given the high post-shock temperatures – the jets conserve only the velocity component $v_\| \propto \cos \theta$ parallel to the outflow axis. We would therefore expect a cosinusoidal anticorrelation, i.e.

$$v_\| \propto 1/\cos \theta,$$

(13)

between the observed jet velocity and the precession cone angle $\theta$. Such an anticorrelation is seen (Blundell & Bowler 2005).

5 SS433 AS A ULX

We can now see why SS433 appears so underluminous compared with its mechanical energy output. Most of its radiative luminosity must be produced within $R_{\text{outflow}}$ of the accretor, i.e. within the outflow. The Thomson optical depth radially across this outflow is $\tau_\perp = \kappa_{\text{e}} \Sigma_{\text{outflow}} \simeq 110$, whereas its central regions are essentially transparent along the flow axis, i.e. $\tau_\parallel \lesssim 1$. Most of the luminosity must therefore escape in directions close to this axis, collimated into a solid angle $\Omega$ set by the geometric spreading of the outflow away from the axis. An observer within this cone would measure a radiation flux density given by $\sim 10 L_{\text{Edd}}/(4\Omega)$ rather than $L_{\text{Edd}}$ and thus identify the source as ultraluminous.

This suggests that one basic cause of ULX behaviour is the collimation of the outgoing radiation (King et al. 2001), and that, as suspected for some time (King et al. 2001; King 2002; Charles et al. 2004), SS433 is a ULX viewed ‘from the side’. The total luminosity escaping sideways (i.e. from the outside of the ‘container’) is $L_{\text{Edd}}$, but thermalized over an area $\sim 4\pi R_{\text{outflow}}^2$. This gives an effective temperature $T_{\text{eff}} \sim 10^5$ K. While this radiation is enough to ensure that the outer disc is efficiently warped, it is far too soft to be observable through the interstellar photoelectric absorption column ($7 \times 10^{21}$ cm$^{-2}$; Kotani et al. 1996) towards SS433. The picture of ULXs presented here differs from the scaled-down BL Lac object (Körding, Falcke & Markoff 2002) sometimes invoked, in that, in our picture the radiation is collimated by scattering, rather than relativistically beamed. Moreover, the jets produced near the accretor have to make their way out through the outflow, and also appear to be definitely sub-relativistic.

We note that some ULXs are surrounded by nebulae which are too large to be supernova remnants and appear to be collisionally ionized (Grisé, Pakull & Motch 2006a,b; Pakull, Grisé & Motch 2006). These are reminiscent of the W50 nebula around SS433, reinforcing our suggestion that ULXs are super-Eddington accreting binaries like SS433.

It is not clear how to estimate a priori the solid angle $\Omega$, the inverse of which specifies just how ultraluminous a ULX can be, without undertaking a numerical simulation involving radiation hydrodynamics (cf. Okuda et al. 2005). Below, we argue that to achieve consistency between the number of progenitor HMXBs in our Galaxy ($\sim 10$) and the number of ULXs ($\lesssim 1$) we require $\Omega/4\pi \lesssim 1/10$ for the latter. This would give a peak apparent luminosity for a hyper-accreting $10 M_\odot$ black hole of $\gtrsim 10^{40}$ erg s$^{-1}$.

The radiation spectrum predicted for ULXs by this picture is a combination of the usual complex medium-energy X-ray binary spectrum, plus a very strong soft component resulting from thermalization of a large fraction of the accretion luminosity $\sim 10L_{\text{Edd}}$ over the walls of the outflow, i.e. over an area $\sim \pi R_{\text{outflow}}^2$, but beamed into a solid angle $\Omega$. This gives a beamed flux with colour temperature $T_{\text{c}} \sim 3 \times 10^6$ K, assuming $\Omega/4\pi \sim \sim 0.1$. Strong soft components of this type are indeed seen in ULXs with sufficiently low absorption columns (e.g. Miller et al. 2003). In our picture, the large blackbody radius $R_b$ deduced from this component gives an estimate of the super-Eddington factor as

$$M_{\text{out}}/M_{\text{Edd}} \simeq \frac{4R_b}{27R_{\odot}},$$

(14)

rather than an estimate of the black hole mass from the assumption that $R_b \sim 3R_{\odot}$. Thus, the IMBH masses $\gtrsim 10^3 M_\odot$ deduced by Miller et al. (2003) for the ULXs NGC 1313 X-1, X-2 become in our picture estimates of the super-Eddington factor $M_{\text{out}}/M_{\text{Edd}} \gtrsim 450$ for a $10 M_\odot$ black hole.

Strohmayer & Mushotzky (2003) use the occasional presence of a 54 mHz (period $\sim 18$ s) QPO seen in the 2–10 keV flux of the ULX M82 X-1 as evidence against this source being highly beamed, and in favour of it being an intermediate mass black hole (IMBH). There are two main points to the argument. First, the high rms amplitude of the QPO (7–8 per cent) would be washed out if the beaming process involves a lot of scattering. In our picture of SS433, most of the luminosity emerges over a collimating region of size $\sim 4$ light seconds. The fact that the QPO in M82 X-1 is observed in the 2–10 keV energy range suggests that it is observed

\[^3\]This energy would be recovered as the post-shock jet fluid reaccelerates unless the shocks are sufficiently radiative – which is likely given the high post-shock temperatures.
in photons which come from close to the black hole and which have not undergone a large amount of scattering. Thus, we are most likely looking down the central regions of the collimating outflow, which are essentially transparent ($\tau_\perp \lesssim 1$, see above). These regions would not significantly wash out a QPO with period $\sim 18$ s unless there was significant optical depth (Kylafis & Klimis 1987). Secondly, if one assumes that the QPO frequency scales with mass of the compact object, then comparison with typical QPO frequencies of $\sim 0.8 - 3$ Hz in stellar mass objects suggests a mass of $\sim 100 - 300 M_\odot$ for this object. Strohmayer & Mushotzky admit that the second point is a weak one and that ‘the broadband variability in the M82 ULX and the lack of a non-thermal component argue against this identification’. We note that in any case there is as yet no believable physical picture of how QPOs are made, and also that, because of the lack of flux, detecting a QPO at around 1 Hz in M82 X-1 would be problematic.

6 THE OBSCURED INTEGRAL SOURCES

In SS433, the walls of the outflow ‘container’ are optically thick ($\tau_\perp \gg 1$). Thus, most of the X-ray radiative output emitted from the outside of this flow towards a terrestrial observer of SS433 is degraded to photon energies too low to be observable through the ISM. This may help explain the difficulty in finding similar ‘sideways-on’ systems as X-ray sources, even though there should be roughly as many of these as there are HMXBs. However, if the mass transfer rate were rather less super-Eddington than in SS433, the total scattering optical depth $\tau_\perp$ across the outflow would be lower. For $\tau_\perp \lesssim 7$, the combined effects of scattering and absorption do not completely thermalize the emerging radiation, which might thus appear as a very hard X-ray continuum with a high intrinsic absorption column $N_H \lesssim \times 10^{20} \, \text{cm}^{-2}$, presumably accompanied by powerful fluorescent emission lines. This is just what is seen in the new class of obscured HMXB systems found by INTEGRAL (Revnivtsev et al. 2003; Walter et al. 2003; Dean et al. 2005). Note that our picture of hyper-Eddington accreting sources applies also to neutron-star systems, as many of the INTEGRAL sources appear to be, with luminosity limits lower by the ratio of neutron-star to black-hole mass. Some of these obscured sources may be modulated (either in scattering or in absorption) at one-half of the orbital period because of the $m = 2$ tide from the companion star (cf. the 6.3-d nodding motion of the jets in SS433).

7 DISCUSSION

We have shown that the effects of the powerful outflow expelling the super-Eddington mass transfer in SS433 can explain several of aspects of this system. The outflow is able to bend the jets from the accretor and redirect them along the axis of the outer disc, which is tilted by radiation torques. As a result, the jets show both the 162-d precession and the 6.3-d nodding motion driven by the companion star’s tide. Our picture also explains the observed anticorrelation of precession cone angle and jet velocity (Blundell & Bowler 2005).

This picture provides an explanation of why SS433 is faint in observed electromagnetic radiation (especially X-rays) compared with the kinetic luminosity of the jets ($L_\perp \gtrsim 10^{39} \, \text{erg} \, \text{s}^{-1}$). Indeed, we have shown that SS433 gives a plausible picture of ULX behaviour if we imagine viewing it along the outflow axis. The radiative luminosity exceeds the formal Eddington value $L_{\text{Edd}}$ by a factor of $\sim 10$, and most of this luminosity is radiated in a cone around the outflow axis. This picture can account for apparent ULX luminosities $\gg L_{\text{Edd}}$ from stellar-mass binaries. HMXB systems undergoing thermal-time-scale mass transfer like SS433 are inevitable in galaxies with vigorous star formation, accounting for the observed correlation of ULXs with star-forming regions in such galaxies, and their birthrate is known to be of the right order to explain the incidence of ULXs (King et al. 2001), independently of the collimation solid angle $\Omega$. Very bright outbursting soft X-ray transients may reproduce some of these features, as the accretion rates driven by disc instabilities can be highly super-Eddington for some time (cf. Cornellis, Charles & Robertson 2006). These may produce transient ULXs in early-type galaxies (King 2002).

In addition, this picture helps to explain the uniqueness of SS433 within our own galaxy. We can get some idea of the collimation solid angle $\Omega$ by considering the progenitors of systems like SS433. These are HMXBs just beginning Roche lobe overflow (e.g. Verbunt & van den Heuvel 1995). After $\sim 10^4$ yr, the mass transfer rate becomes super-Eddington and the system becomes an SS433-like object. Other HMXBs, such as the Be-X-ray binaries and supergiant wind-fed systems, appear less likely to evolve into systems like SS433, so we estimate the current number of SS433 progenitors in the Galaxy as $\sim 10$. The thermal-time-scale mass-transfer stage characterizing SS433 should last at least as long as the preceding HMXB stage, and should exhibit a comparable accretion luminosity. So, were SS433-like systems to emit isotropically, we might have expected to find as many SS433-like systems as relevant HMXBs (i.e. $\sim 10$) in the Galaxy. The fact that we have only one SS433-like object, which we suggest is a ULX seen from the side, probably tells us that such objects are hard to detect. However, the fact that we see no ULXs, which we identify as SS433-like objects collimated towards us and which should be easily detectable X-ray sources, suggests that the mean solid angle into which most of the radiation is emitted is $\Omega / 4\pi = \lesssim 1/10$. It may be possible to get better estimates of the mean value of $\Omega$ by comparing the numbers of ULXs and HMXBs in galaxies where both are detected. We note that it is to be expected that $\Omega$ varies among ULXs, possibly as a function of the Eddington ratio.

As a further check on our adopted value of $M_{\text{out}}$, we note that the outflow is evidently quasi-spherical at large distances from the binary, and drives a wind bubble into the interstellar medium. It is easy to check that this must be in the energy-driven phase, giving a radius

$$R_{\text{bubble}} \approx 0.9 \left( \frac{M_{\text{out}} t_w^2}{2 \rho} \right)^{1/5} \rho^{3/5} = 33 \rho^{1/5} \left( \frac{t}{t_{\text{Edd}}} \right)^{3/5} \text{pc}$$

where $\rho = 10^{-20} \rho_{26} \, \text{g} \, \text{cm}^{-3}$ is the mass density of the ISM. We see that the outflow may well have inflated the almost perfectly spherical ‘head’ of the W50 nebula (whose observed radius is $\approx 42$ pc; see Fabrika 2006 and references therein) within its mass transfer lifetime. (Note that the precessing jets inflate the ‘ears’ of W50 still more.)

The jet precession, together with the tidally induced nutation, is a useful diagnostic in SS433. However, it is important to realize that regular precession is not generic to the ULX picture we have set out (Wijers & Pringle 1999; Ogilvie & Dubus 2001). For example, in other similar objects, and in SS433 itself at other epochs, the black hole spin may not be aligned with the binary orbit, the disc may not be tilted by radiation torques, and even if tilted may not precess, let alone regularly. If we were lucky enough to lie in the ULX visibility cone of a system undergoing regular precession like

4The Be-X-ray binaries are generally too wide for full Roche lobe overflow to occur, and the supergiant systems are likely to merge.
Our picture suggests that some ULXs might show mildly blueshifted spectral features. Velocities such as the 0.26 km/s in SS433 are probably at the limit of current observational capabilities, but would offer major insight if detected. Similarly, it is probably optimistic to hope to see 0.5 MeV pair emission from a jet bending shock in ULXs (the oblique shock produces near-relativistic ion temperatures \( \sim 10^7 \) K), particularly if this proves undetectable in edge-on Galactic systems.

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