The Magnetic Field of the Ultraluminous X-ray Pulsar M82 X-2

Dimitris M. Christodoulou\textsuperscript{1}, Silas G. T. Laycock\textsuperscript{2}, and Demosthenes Kazanas\textsuperscript{3}

ABSTRACT

Pulsations were recently detected from the ultraluminous X-ray source X-2 in M82. The newly discovered pulsar has been described as a common neutron star with a 1 TG magnetic field that accretes above the Eddington rate and as a magnetar-like pulsar with a 100 TG magnetic field that is above the quantum limit. We show here that this X-ray source is not exotic in any sense. The standard equations of accretion along field lines predict that, for the measured spin period $P_S$ and spinup rate $\dot{P}_S$, the isotropic X-ray luminosity $L_X$ must be near the Eddington limit (i.e., $L_X \approx 3.5 \times 10^{38}$ erg s$^{-1}$); and that the surface magnetic field $B$, that does not depend on $P_S$, must be modest (i.e., $B \approx 1 - 10$ TG). The observed higher luminosity can be explained by a moderate amount of geometric beaming that occurs in our direction. Other ultraluminous X-ray sources may also turn out to be common pulsars with similar physical characteristics, but since the emission must occur at a favorable angle to the observer, we expect that very few such pulsars will be discovered in the future.

Subject headings: accretion, accretion disks—stars: magnetic field—stars: neutron—X-rays: binaries—X-rays: individual (NuSTAR J095551+6940.8)

1. Introduction

Ultraluminous X-ray (ULX) sources are extragalactic compact accreting objects characterized by super-Eddington luminosities ($L_X \sim 10^{39-41}$ erg s$^{-1}$) and unusual soft X-ray

\textsuperscript{1}University of Massachusetts Lowell, Dept. of Mathematical Sciences, Lowell, MA 01854. E-mail: Dimitris\textunderscore Christodoulou@uml.edu

\textsuperscript{2}University of Massachusetts Lowell, Dept. of Physics & Applied Physics, Lowell, MA 01854. E-mail: Silas\textunderscore Laycock@uml.edu

\textsuperscript{3}NASA Goddard Space Flight Center, Code 663, Greenbelt, MD 20771. E-mail: Demos.Kazanas@nasa.gov
spectra with blackbody emission around $\lesssim 0.3$ keV and a downturn above $\sim 5$ keV (Gladstone, Roberts, & Done 2009, Feng & Soria 2011, Motch et al. 2014). The extreme luminosities can be understood either as emission at the Eddington limit ($L_X \lesssim L_{\text{Edd}} = 1.3 \times 10^{38} \ M/\text{M}_\odot \ \text{erg s}^{-1}$) from intermediate-mass ($M \sim 10^{2-4} \text{M}_\odot$) black holes or as anisotropic emission with $L_X \gtrsim L_{\text{Edd}}$ from stellar-mass black holes and neutron stars (Soria 2007, Feng & Soria 2011, Medvedev & Poutanen 2013, Bachetti et al. 2014, Motch et al. 2014, Pasham et al. 2014).

The latter interpretation is gaining favor from recent observations: (a) Luangtip et al. (2014) analyzed a Chandra sample of nearby ULX sources and found a change in the spectral index around $L_X \sim 2 \times 10^{39} \ \text{erg s}^{-1}$ that may indicate a transition to the super-Eddington accretion regime by $10 \text{M}_\odot$ black holes; (b) Motch et al. (2014) determined that the ULX source P13 in NGC7793 harbors a stellar-mass black hole with $M < 15 \text{M}_\odot$; and (c) Bachetti et al. (2014) determined that the ULX source X-2 in M82 harbors a pulsar with spin period $P_S = 1.3725 \ \text{s}$ and spinup rate $\dot{P}_S = -2 \times 10^{-10} \ \text{s s}^{-1}$. In all of these cases, the observations indicate that the X-ray sources radiate above their Eddington limits, a result that makes uncomfortable many researchers.

The magnetic field of NGC7793 P13 has not been estimated, but in the case of M82 X-2, Bachetti et al. (2014) used their measurement of the accretion torque to obtain a modest value of the magnetic field $B \approx 1 \ \text{TG}$. In contrast, Ekşi et al. (2014) argued that the accretion torque is very small near spin equilibrium and that leads to an underestimate of the magnetic field which in M82 X-2 must be of magnetar strength ($B \sim 100 \ \text{TG}$) (see also Ho et al. 2014, Kluś et al. 2014, Lyutikov 2014). We, on the other hand, find that the magnetic field implied from the measurements of M82 X-2 is modest ($B \approx 1 - 10 \ \text{TG}$). No unusual physics is necessary to obtain this estimate, the result is derived from standard accretion theory and the measured value of $\dot{P}_S$. More importantly, such modest $B$-values are consistent with isotropic X-ray emission near the Eddington limit for the measured values of $P_S$ and $\dot{P}_S$. This consistency implies that M82 X-2 is a common pulsar that emits X-rays anisotropically toward our direction, an explanation also favored by Bachetti et al. (2014).

These results are already present in the equations of accretion onto compact objects (e.g., Frank et al. 2002), but the physics is obscured by the multiple scalings and some minor inconsistencies in the definitions of the physical quantities involved. We thus derive and discuss the accreting pulsar’s isotropic luminosity and surface magnetic field in § 2 without scaling the various physical variables, and we discuss our conclusions for M82 X-2 and other ULX sources in § 3 below.
2. Equations for Accretion onto a Compact Object

We adopt Gaussian units for electromagnetic quantities (Jackson 1962) and we consider accretion from a disk-like structure that is formed around a compact object (Pringle & Rees 1972, Frank et al. 2002).

2.1. Accretion Torque

We introduce two control parameters, $\xi$ and $n$, in the calculation of the accretion torque from disk parameters:

$$\tau_d = n\dot{M}\sqrt{GM(\xi r_m)},$$  \hspace{1cm} (1)

where $\dot{M}$ is the mass accretion rate, $G$ is the gravitational constant, $M$ is the mass of the compact object, and $r_m$ is the spherical magnetospheric (Alfvén) radius. In particular:

(a) We assume that the accreted matter is forced to follow magnetic field lines at a cylindrical radius $R_m \equiv \xi r_m$ (Frank et al. 2002) and $\xi$ accounts for the effects of varying the location of $R_m$ around $r_m$. For example, Frank et al. (2002, § 6.3) discuss a range of $0.5 < \xi < 2$ and they favor $\xi \approx 0.5$ (Ghosh & Lamb 1979), while others (e.g., Stella et al. 1986; Galache et al. 2008) adopt $\xi = 1$ for simplicity.

(b) Eksi et al. (2014) and Lyutikov (2014) argue for spin equilibrium in M82 X-2 that reduces the accretion torque by a factor of $\approx 10^{-100}$ (although the observed event has certainly taken the system out of equilibrium). We account for such a reduction by inserting the parameter $n$ into eq. (1), where $0 \leq n \lesssim 3.9$ (Ghosh & Lamb 1979), and then $n \approx 0.1 - 0.01$ according to these arguments.

We note however that setting $n \approx 0.1 - 0.01$ is not going to have the desired outcome (a much smaller disk torque). The accretion torque can be measured (Bachetti et al. 2014) as it can also be expressed in terms of the transfer of angular momentum onto the compact object:

$$|\tau_*| = |\dot{\mathcal{L}}_*| = \frac{2\pi I_*}{P_S^2}|\dot{P}_S|,$$  \hspace{1cm} (2)

where $\mathcal{L}_* = 2\pi I_*/P_S$ and $I_*$ are the angular momentum and the moment of inertia, respectively, of the compact object with spin period $P_S$ and derivative $\dot{P}_S$. Thus, assuming that $|\tau_*| = \tau_d$, for any reduction of $n$ in eq. (1), the $\dot{M}$ will have to increase accordingly to produce the required $|\tau_*|$ magnitude. This will artificially push the X-ray luminosity $L_X \propto \dot{M}$ higher by a factor of $1/n$, while we should be interested in the luminosity that the standard disk torque with $n = 1$ is capable of producing. We will return to this issue in § 2.3 below.
2.2. Magnetospheric Radius

Calculations involving the spherical magnetospheric (Alfvén) radius \( r_m \) in the literature (Elsner & Lamb 1977, Ghosh & Lamb 1979, Frank et al. 2002, Ekşi et al. 2014) suffer from inconsistencies in the definitions of the magnetic moment \( \mu \) and the ram pressure \( P_{\text{ram}} \) of the inflowing matter. Eq. (11) of Ghosh & Lamb (1979) agrees with eq. (6) of Elsner & Lamb (1977) only if \( P_{\text{ram}} = \rho v^2/2 \) in the former calculation, where \( \rho \) is the mass density and \( v \) is the inflow speed of matter at spherical radius \( r \). In both cases, the magnetic moment is defined as \( \mu = Br^3 \) on dimensional grounds, where \( B \) is the magnetic field at \( r \); but other authors use \( \mu = Br^3/2 \) instead (e.g., Ekşi et al. 2014). Thus, various factors of \( 1/2 \) propagate in the calculations but they cause only small differences in the coefficients of the final results because of the steep dependence of \( B^2(r) \) on \( 1/r^6 \).

We adopt here the following definitions: for the ram pressure, \( P_{\text{ram}} = \rho v^2/2 \) (consistent with Bernoulli’s equation); and for the magnetic moment, \( \mu = Br^3/2 \) (Jackson 1962, eqs. [5.41] and [5.42]). Then \( 2\mu_* = Br^3 = B_\ast R^3_\ast \), where asterisks indicate the quantities at the surface of the compact object, and the magnetic pressure at radius \( r \) can be written as \( P_{\text{mag}} = B^2/(8\pi) = \mu_*^2/(2\pi r^6) \). As usual (e.g. Frank et al. 2002), \( \rho|v| = \dot{M}/(4\pi r^2) \) and \( |v| \) is close to the free-fall speed \( v_{\text{ff}} = \sqrt{2GM/r} \), in which case the balance between the two pressures at \( r = r_m \) determines the Alfvén radius:

\[
\frac{2}{3} \mu_*^4 = \frac{8}{3} \frac{GM}{\dot{M}} \left( \frac{2GM\mu_*^4}{L_X^2 R^2_\ast} \right)^{1/7}. \tag{3}
\]

It is convenient to replace \( \dot{M} \) in terms of the X-ray luminosity \( L_X \), which is assumed to be one-half of the total accretion power, i.e.,

\[
L_X = \frac{3}{2} \frac{GM\dot{M}}{R_*}, \tag{4}
\]

in which case eq. (3) takes the form

\[
\frac{2}{3} \mu_*^4 = \frac{8GM\mu_*^4}{L_X^2 R^2_\ast} \left( \frac{2GM\mu_*^4}{L_X^2 R^2_\ast} \right)^{1/7}. \tag{5}
\]

2.3. Isotropic X-Ray Luminosity

Assuming that \( |\tau_*| = \tau_d \) and using the cylindrical radius \( R_m = \xi r_m \), we get from eqs (1), (2), and (4):

\[
L_X = \frac{\pi I_*}{n R_* P_S S} |\dot{P}_S| \sqrt{\frac{GM}{R}}. \tag{6}
\]
If we ask for the inflowing matter that reaches $R_m$ to also corotate with the magnetic field lines, then $R_m$ must be equal to the corotation radius $r_c$:

$$R_m = r_c \equiv \left( \frac{GMP_s^2}{4\pi^2} \right)^{1/3}. \tag{7}$$

Eliminating $R_m$ between eqs. (6) and (7), we find that

$$L_X = \frac{\pi I_*}{n} \frac{\dot{P}_s |\dot{P}_s|}{2\pi G M} \left( \frac{2\pi GM}{P_s^2 R_*^3} \right)^{1/3}. \tag{8}$$

It is convenient to rewrite this result in terms of the angular velocity of the compact object $\Omega_S \equiv 2\pi/P_S$ and the Keplerian angular velocity on the surface $\Omega_{K*} \equiv \sqrt{GM/R_*^3}$:

$$L_X = \frac{I_*}{2n} \left( \Omega_S \Omega_{K*}^2 \right)^{1/3} |\dot{\Omega}_S|. \tag{9}$$

The cubic root in this equation is the geometric mean of $\Omega_S$, $\Omega_{K*}$, and $\Omega_{K*}$. Thus, the luminosity $L_X \propto |\dot{\Omega}_S|$ but the proportionality constant is heavily weighted by $\Omega_{K*}$, the larger of the two frequencies in the geometric mean, which also expresses the depth of the gravitational potential well of the compact object. The spin contributes to the geometric mean to a much lesser extent. This implies that a source that is spun up by accretion will increase its power output marginally ($L_X \propto \Omega_S^{1/3}$) and corotation will move inward very slowly. For a compact object spinning near breakup ($\Omega_S = \Omega_{K*}$), corotation is near the surface of the object ($r_c \approx R_*$) and eq. (9) with $n = 1$ reduces to $L_X = I_* \Omega_{K*} |\dot{\Omega}_K|/2$ which is simply the rate of radiative loss of one-half of the rotational kinetic energy $E_{\text{rot}} = I_* \Omega_{K*}^2/2$ that reaches the surface.

In eqs. (8) and (9), the effect of parameter $n < 1$ is to increase the X-ray luminosity by a factor of $1/n$. This is understood as an increase to the mass inflow rate $\dot{M}$: since $L_X \propto \dot{M}$ (eq. [4]), then $\dot{M} \propto 1/n$ as well. This occurs because in the torque equation $|\tau| = \tau_d$ we ask for the compact object to be spun up to the observed level of $|\dot{\Omega}_s|$ by a disk torque $\tau_d \propto n\dot{M}$, and so the mass inflow rate must increase in proportion to compensate for the reduced $n < 1$. Such a transient high $\dot{M}$ rate can certainly match the power output of M82 X-2, but it is not a good measure of the strength $B_*$ of the stellar magnetic field that is certainly overestimated when the peak $L_X$ emission from the system is used (Ekşin et al. 2014, Lyutikov 2014). More robust estimates can be obtained for both $L_X$ and $B_*$ from the measured values of $P_S$ and $\dot{P}_S$, as in § 2.4 and 2.5 below.
2.4. Stellar Magnetic Field

We set again $|\tau_s| = \tau_d$ and we combine eqs (1), (2), (4), (5), and (8) to get a relation that involves the stellar magnetic moment. We find that

$$\mu_* = \frac{1}{2} \left( \frac{2\pi^2 n^2 \xi^7}{2} \right)^{-1/4} \sqrt{GMI_* |\dot{P}_S|},$$

or, using $\mu_* = B_* R_*^3 / 2$, that the surface magnetic field is

$$B_* = \left( \frac{2\pi^2 n^2 \xi^7}{2} \right)^{-1/4} \sqrt{\frac{GMI_*}{R_*^6} |\dot{P}_S|}.$$  \hspace{1cm} (11)

Unlike in the case of the X-ray luminosity (eq. [8]), the results in eqs. (10) and (11) do not at all depend on the precise amount of the inflowing accretion power that is converted to X-rays. Therefore, the efficiency of the conversion process (assumed to be 50\% in eq. [4]) does not enter in the calculation of $B_*$. This result makes physical sense; it can be verified easily by changing the ”efficiency” factor of 1/2 in the definition of $L_X$ in eq.(4) and repeating the calculation.

Written in terms of $\Omega_S$ and $\Omega_{K*}$, eq. (11) becomes

$$B_* = \left( \frac{n^2 \xi^7}{2} \right)^{-1/4} \left( \frac{\Omega_{K*}}{\Omega_S} \right) \sqrt{\frac{I_*}{R_*^3} |\dot{\Omega}_S|}.$$ \hspace{1cm} (12)

This equation relates the surface magnetic pressure to the $R\phi$ component of the stress tensor $[\rho v_{RvK}]_{R_m}$ at $R_m$, where $\rho(R_m) v_{R(R_m)} = \dot{M}/(4\pi R_m^2)$ and $v_K(R_m) = (GM/R_m)^{1/2}$: squaring both sides of eq. (12), we find that

$$\frac{B_*^2}{8\pi} = \left( \frac{2\xi^7}{2} \right)^{-1/2} \left( \frac{\Omega_{K*}}{\Omega_S} \right)^2 \left( \frac{R_m}{R_*} \right)^3 |\rho v_{RvK}|_{R_m},$$ \hspace{1cm} (13)

independent of the control parameter $n$. This is because the torque balance $|\tau_s| = \tau_d$ implies that

$$\frac{I_*}{R_*^3} |\dot{\Omega}_S| = 4\pi n \ |\rho v_{RvK}|_{R_m} \left( \frac{R_m}{R_*} \right)^3,$$ \hspace{1cm} (14)

and $n$ cancels out when eqs. (12) and (14) are combined. Therefore, one cannot regulate the surface magnetic pressure by adjusting the torque at corotation, a result that is physically correct.

Eq. (13) shows that the stress tensor at $R_m$ is strongly rescaled by both the spin frequency of the compact object (ratio $(\Omega_{K*}/\Omega_S)^2 >> 1$) and by the scale of the accretion
flow (ratio \((R_m/R_∗)^3 \gg 1\)) in order to match the magnetic pressure on the surface of the compact object. Furthermore, eq. (13) reveals the following two notable accretion properties:

(a) If the corotation radius \(R_m\) is smaller, then the magnetic field is weaker and vice versa, precisely as expected. (b) If the compact object is spun up (\(\Omega_S\) increases), then the magnetic field appears to be weaker and vice versa, again as expected. The two effects do not work against one another because the direction in which the two ratios move is the same.

Let us now imagine that the compact object is spun up from accretion at a certain moment in time. Then \(P_S\) decreases, \(\Omega_S\) increases, and corotation pushes inward (eq. [7]). This results in a decrease of both ratios in eq. (13) and then, since \(B_∗\) remains unchanged on the surface, the \(R_\phi\) component of the stress tensor at \(R_m\) in eq. (13) must necessarily increase. Since the condition \(|\tau_*| = \tau_d\) and eqs. (1) and (2) imply that \(\dot{M} \propto 1/(P_S^2 R_m^{1/2})\), then \(\dot{M}\) also increases and this leads to an increase in the X-ray luminosity. So accretion runs away to higher luminosities; as long as there is available matter of density \(\rho(R_m)\), the inflow will continue and the compact object will continue spinning up. The accretion event will begin powering down only when the density \(\rho(R_m)\) or the \(\dot{M}\) at corotation will drop and the loading of matter onto magnetic field lines will taper off.

### 2.5. Application to M82 X-2

In order to apply eqs. (8) and (11) to the case of M82 X-2, we use \(G = 6.67 \times 10^{-8} \text{ cm}^3 \text{ g}^{-1} \text{ s}^{-2}\) and the following typical pulsar parameters: \(I_* = 2MR_0^2/5\), \(M = 1.4M_\odot\), \(R_0 = 10 \text{ km}\), and \(\xi = 0.5\) (Ghosh & Lamb 1979). We find that

\[
L_X = \frac{3.7 \times 10^{38}}{n} \left(\frac{P_S}{1 \text{ s}}\right)^{-7/3} \left(\frac{|\dot{P}_S|}{10^{-10} \text{ s s}^{-1}}\right) \text{ erg s}^{-1},
\]

and

\[
B_\ast = \frac{7.3 \times 10^{12}}{\sqrt{n}} \sqrt{\frac{|\dot{P}_S|}{10^{-10} \text{ s s}^{-1}}} \text{ G}.
\]

In the case of M82 X-2 with \(P_S = 1.3725 \text{ s}\) and \(|\dot{P}_S| = 2 \times 10^{-10} \text{ s s}^{-1}\), these values become

\[
L_X = \frac{3.5 \times 10^{38}}{n} \text{ erg s}^{-1},
\]

and

\[
B_\ast = \frac{1.0 \times 10^{13}}{\sqrt{n}} \text{ G}.
\]

These typical values show a weak dependence of the parameter \(n\). For \(n \approx 0.1\) (Ekşi et al. 2014), eqs. (17) and (18) give modestly higher values for M82 X-2: \(L_X = 3.5 \times 10^{39} \text{ erg s}^{-1}\). 

and $B_\ast = 3 \times 10^{13}$ G, respectively. The dependence of $B_\ast$ on parameter $\xi$ in eq. (11) is more important: for $\xi = 2$ (i.e., corotation is now outside the Alfvén radius) and $n = 1$, eq. (11) gives a significantly lower value of $B_\ast = 9 \times 10^{11}$ G for M82 X-2 that is comparable to the estimate by Bachetti et al. (2014); and for $\xi = 1$ (spin equilibrium) and $n = 1$, we find that $B_\ast = 3 \times 10^{12}$ G.

3. Discussion and Conclusions

We have used standard accretion theory to estimate the typical values of the isotropic luminosity and the surface magnetic field for the ULX pulsar M82 X-2 for which Bachetti et al. (2014) have recently measured pulsations with period $P_S = 1.3725$ s and spinup rate $\dot{P}_S = -2 \times 10^{-10}$ s s$^{-1}$. For these values of $P_S$ and $\dot{P}_S$, our results show that the luminosity is $L_X = 3.5 \times 10^{38}$ erg s$^{-1}$, a value that is only $\sim 2$ times larger than the Eddington limit. This is a sign of supercritical accretion onto the polar regions of the pulsar as a result of the observed episodic event. This result is not entirely surprising: super-Eddington X-ray emission has been previously observed from A0538-66 in the Large Magellanic Cloud (Skinner et al. 1982) and the spin period of the pulsar was then measured. All other X-ray observations of A0538-66 since then, summarized by Kretschmar et al. (2004), have measured much smaller luminosities that fall below the "propeller line" (Stella et al. 1986, Christodoulou & Laycock 2014). This led Campana et al. (1995) and Campana (1997) to propose that the low-power X-rays are emitted from the magnetosphere where accretion is halted by the centrifugal force and a luminous shock forms. M82 X-2, for which Kong et al. (2007) have reported two cases of no detection, may also be caught in such nonpulsating sub-Eddington states by future observations after the current event has died out, and spectral differences (e.g., softer spectra) ought to provide an indication that the propeller line has been crossed. In fact, another ULX source hosting a stellar-mass black hole (NGC7793 P13) has been observed in such a "faint" state ($L_X = 5 \times 10^{37}$ erg s$^{-1}$; Motch et al. 2014).

Despite the apparent super-Eddington X-ray emission from M82 X-2, its surface magnetic field is estimated to be $B_\ast = 1 - 10$ TG depending on where the corotation radius $r_c$ is located relative to the Alfvén radius $r_m$ (for $\xi \equiv r_c/r_m = 2 - 0.5$, respectively). From these estimates, we conclude that there is no need to invoke powerful, magnetar-type magnetic fields to explain the pulsar’s X-ray power output, standard accretion mechanisms suffice. Eksi et al. (2014) and Lyutikov (2014) advocated for such very strong magnetic fields ($B \sim 10 - 100$ TG) because they thought that M82 X-2 is in spin equilibrium characterized by a substantial reduction in accretion torque. Our results in §§ 2.3–2.5 and the measured spinup rate $\dot{P}_S$ argue against this claim and against using the peak X-ray luminosity in or-
order to estimate the magnetic field (this is also a problem in the calculation of Tong [2014], although $L_X$ was reduced from its peak by an anisotropic factor of 5, causing a reduction of $B_\text{x}$ by $5^3$ and making it agree with our estimate).

A modest amount of anisotropy (geometric collimation $b \sim 0.1$) appears to be sufficient to explain the observed super-Eddington luminosities if $\dot{M}$ exceeds the Eddington rate $\dot{M}_\text{Edd}$ by $\sim 100$ during outbursts (King 2008, 2009, Feng & Soria 2011, Tong 2014). Then, the luminosity is boosted by an overall factor of at least $[1 + \ln(\dot{M}/\dot{M}_\text{Edd})]/b \sim 60$. Furthermore, if the emission from M82 X-2 is collimated, some of the energy must emerge at much longer wavelengths. This is the case according to the radio observations of Kronberg et al. (1985), McDonald et al. (2002), and Fenech et al. (2008); and the infrared observations of Kong et al. (2007) and Gandhi et al. (2011). The radio maps show a core-dominated source which is expected if the pulsar is a modestly aligned rotator and a collimated jet is coming out in our direction. Kong et al. (2007) also produced Chandra X-ray spectra that are hard (photon indices 1.3 – 1.7 from an absorbed power-law model) and show no soft excess. This, combined with the strong X-ray variability on timescales of $\sim 2$ months and the recurring outbursts, indicates that M82 X-2 is not at all dissimilar to Galactic and Magellanic X-ray binaries harboring neutron stars.

There are some more discoveries about ULX sources indicating that they may be supercritical accretors but otherwise not at all exotic: (a) Gilfanov et al. (2004) found that the ULX sources are just the high end of a luminosity function that cuts off at $L_X \sim 3 \times 10^{40}$ erg s$^{-1}$ and in which the known high-mass X-ray binaries make up the low end of a single power-law with slope $\sim 1.6$. (b) Luangtip et al. (2014) found that the spectral index of ULX spectra changes around $L_X \sim 2 \times 10^{39}$ erg s$^{-1}$, a value that may indicate a transition from critical to supercritical accretion by common $10M_\odot$ black holes. (c) Laycock et al. (2014) determined from the radial velocity curve of IC10 X-1 that the the compact object is likely a neutron star, although a low-stellar-mass black hole cannot be ruled out. (d) Liu et al. (2013) determined from optical spectroscopy that the mass of the compact object in M101 ULX-1 is no more than $30M_\odot$. As in the case of NGC7793 P13 (Motch et al. 2014), it is unlikely that this is an intermediate-mass black hole. For this system, Shen et al. (2014) developed a model that explains how reprocessed soft X-rays can be finally emitted from very large radii ($\sim 100$ times beyond the inner radius of the accretion disk). This model may also be applicable to other ULX sources harboring stellar compact objects (see also King 2008).

Soon after the discovery of pulsations in M82 X-2, Doroshenko et al. (2014) tried to find more pulsating ULX sources in archival XMM-Newton observations, but they did not succeed. We believe that it is a matter of time until another pulsating ULX source is found.
But because the X-ray emission from neutron stars must occur at a favorable angle to the observer, we expect very few such pulsars to be discovered in the future (see also King 2009).

This work was supported in part by NASA grant NNX14-AF77G.
REFERENCES

Jackson, J. D. 1962, Classical Electrodynamics, New York: John Wiley & Sons

This preprint was prepared with the AAS LaTeX macros v5.2.