Supercritical Accretion onto a Non-magnetized Neutron Star: Why is it Feasible?

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Abstract

To understand why supercritical accretion is feasible onto a neutron star (NS), we carefully examine the accretion flow dynamics by 2.5-dimensional general relativistic radiation magnetohydrodynamic (RMHD) simulations, comparing the cases of accretion onto a non-magnetized NS and that onto a black hole (BH). Supercritical BH accretion is relatively easy, since BHs can swallow excess radiation energy, so that radiation flux can be inward in its vicinity. This mechanism can never work for an NS, which has a solid surface. In fact, we find that the radiation force is always outward. Instead, we found significant reduction in the mass accretion rate due to strong radiation-pressure-driven outflow. The radiation flux \( F_{\text{rad}} \) is self-regulated such that the radiation force balances with the sum of gravity and centrifugal forces. Even when the radiation energy density greatly exceeds that expected from the Eddington luminosity \( L_{\text{Edd}} \), the radiation flux is always kept below a certain value, which makes it possible not to blow all the gas away from the disk. These effects make supercritical accretion feasible. We also find that a settling region, where accretion is significantly decelerated by a radiation cushion, is formed around the NS surface. In the settling region, the radiation temperature and mass density roughly follow \( T_{\text{rad}} \propto r^{-1} \) and \( \rho \propto r^{-3} \), respectively. No settling region appears around the BH, so matter can be directly swallowed by the BH with supersonic speed.

Key words: accretion, accretion disks – magnetohydrodynamics (MHD) – radiation: dynamics – stars: black holes – stars: neutron

1. Introduction

There is growing evidence of the supercritical (or super-Eddington) accretion objects (hereafter, super-Eddington accretors) in the universe. Super-Eddington accretors are very powerful engines and thus play essential roles in various astrophysical phenomena (e.g., emitting high-energy emission and/or launching relativistic baryon jets). They can also give large impacts on their environments through intense radiation and massive outflow, thereby giving rise to interesting activities (e.g., creating huge ionized nebulae). It is thus worth studying the detailed processes associated with super-Eddington accretors from various viewpoints.

One of the most promising candidates for the super-Eddington accretors is ULXs, compact Ultraluminous X-ray sources, which were successively discovered in nearby active galaxies (Fabbiano et al. 1989; Liu 2011; Walton et al. 2011). The ULXs are off-nuclear point sources producing very large X-ray luminosity, \( L_x > 10^{39} \text{ erg s}^{-1} \), far exceeding the Eddington limit \( (L_{\text{Edd}}) \) of a stellar mass black hole (BH). There are two major scenarios so far proposed and discussed to explain their nature: (1) sub-Eddington accretors harboring an intermediate mass black hole (IMBH) with mass exceeding \( 100 M_\odot \) (e.g., Makishima et al. 2000; Miller et al. 2004), and (2) super-Eddington accretors harboring a stellar mass BH with super-Eddington rates with \( M \gg L_{\text{Edd}}/c^2 \) (e.g., King et al. 2001; Watarai et al. 2001; Poutanen et al. 2007). Quite recently, one very convincing piece of evidence in favor of the latter scenario has been reported; that is, the discovery of pulses in one of the ULXs M82 X-2 (Bachetti et al. 2014). This discovery has established that at least some part of ULXs is super-Eddington accretors (ULX Pulsars, see Fürst et al. 2016; Israel et al. 2017a, 2017b, for the discovery of other cases).

The ULXs are not the only candidates for super-Eddington accretors; however, there are actually plenty of other objects known to date, that are suspected to host supercritical accretion flow. One good example is ULSs, Ultraluminous supersonic sources, which have similarly high X-ray luminosities but exhibit much softer X-ray spectra with typical photon energy of \( \sim 0.1 \text{ keV} \) (e.g., Di Stefano & Kong 2003; Kong et al. 2004). These features can be understood, if one observes super-Eddington accretors from a nearly edge-on direction (Gu et al. 2016; Urquhart & Soria 2016; Ogawa et al. 2017). Other candidates include microquasars, TDE (tidally disrupted events), narrow-line Seyfert 1 galaxies (Wang et al. 1999; Mineshige et al. 2000), and so on. Super-Eddington accretors are unique in the sense that their energy release rate does not depend on their internal properties at all but on the external conditions; i.e., mass supply rate to the compact object vicinity.

Parallel to the accumulation of observational evidence supporting the ubiquitous existence of super-Eddington accretors, semi-analytic and simulation studies have been conducted rather extensively. The possibility of supercritical accretion onto the compact star was first discussed in the pioneering paper by Shakura & Sunyaev (1973; hereafter SS73). Abramowicz et al. (1988) found an equilibrium solution of the supercritical disk and constructed the so-called slim disk model, in which advection of radiation entropy plays a crucial role (see Watarai & Fukue 1999, for a simplified self-similar solution of the slim disk). The general relativistic (GR) version of the slim disk was first constructed by Beloborodov (1998), who claimed that the thermalization timescale could be longer...
than the accretion timescale, so that radiation and matter temperatures may deviate. The supercritical accretion disk has also been discussed in the context of magnetized and/or non-magnetized NSs. In the case of accretion onto a magnetized NS, the accretion mode through the disks quenches due to the strong magnetic pressure. Gas then falls onto the NS surface along the magnetic field lines, thereby forming accretion columns (Basko & Sunyaev 1976; Lyubarskii & Sunyaev 1988). The emission from the accretion columns can reach $10^{40}$ erg s$^{-1}$ (Mushotkov et al. 2015), which is consistent with recent observations of the ULX pulsars.

The pioneering simulation work was made by Ohsuga et al. (2005) using radiation hydrodynamic (RHD) simulations (see also Eggum et al. 1985, 1988). They could for the first time succeed in producing steady-state supercritical accretion flow and revealed various unique features, such as the anisotropic radiation field, wide-angle outflow, large-scale circulation of gas within the flow, and so on. The most up-to-date simulations were performed under the full GR treatments, including the magnetic field for BHs (McKinney et al. 2014; Sadowski et al. 2014, 2017; Sadowski & Narayan 2015a, 2016; Takahashi et al. 2016) and for NSs (Takahashi & Ohsuga 2017), and found the formation of strong outflows (Sadowski & Narayan 2015b; Takahashi & Ohsuga 2015). Takahashi et al. (2016) demonstrate that the hot accretion flow is formed close to the compact object and it can be responsible for hard X-ray emission.

We wish to address one key question here; why is supercritical accretion feasible? Another related question is whether there are no practical limits on mass accretion rates and luminosities, provided that a sufficient amount of mass is supplied externally. Through the number of simulation studies conducted recently, we now have a consensus that it is really feasible to put as much material into a BH as you like. We should be careful, however, since the simulations only give results, while it is our task to specify the mechanisms underlying them. A popular argument made in this context is as follows: supercritical accretion is feasible, since radiation goes out in the perpendicular direction to the disk plane, thus giving little effects on the matter that accretes along the disk plane. This explanation is not complete, however, since it misses the consideration of the force balance on the equatorial plane, although radiation force should also give enormous impacts on the material there. What is needed is to give a clear explanation as to why matter can accrete toward the region full of radiation energy.

It is interesting to note in this respect that Ohsuga & Mineshige (2007) discussed this problem, by using their RHD simulation data. They have found two key ingredients that make it possible to excite supercritical flow: the anisotropic radiation field created by large $\tau$ accretion flow from the equatorial plane and photon trapping effects. Photons created deep inside the thick accretion flow are trapped within the flow and finally swallowed by a BH before escaping from the surface of the flow. The outgoing radiative flux is thus largely attenuated (or sometimes flux becomes inward) so that supercritical accretion is feasible onto BHs.

How about the cases of NS accretions? We should point out that photon trapping cannot be so effective on a long timescale there, since photons should eventually be emitted from the solid surface of an NS. As a result, radiation force should always be outward, thereby decelerating accreting gas. Supercritical accretion is relatively easier, if the NS is strongly magnetized and if accretion occurs through a narrow accretion column (i.e., ULX pulsars). This is because excess radiation energy can then almost freely escape from the side wall of the accretion column (Basko & Sunyaev 1976; Kawashima et al. 2016; Takahashi & Ohsuga 2017). In this paper, we make careful analysis of the GR simulation data to find an answer to the question that asks why the super-Eddington accretion onto a non-magnetized NS is feasible. The paper is organized as follows: we will describe the methods of calculations in Section 2 and then present results in Section 3. The final section is devoted to the discussion of observational implications and other related issues.

2. Basic Equations and Numerical Method

We numerically solve general relativistic radiation magnetohydrodynamic (GR-RMHD) equations in which the radiation equation is based on a moment formalism with applying a M-I closure (Levermore 1984; Kanno et al. 2013; Sadowski et al. 2013). In the following, Greek suffixes indicate spacetime components, and Latin suffixes indicate space components. We take the light speed $c$ as unity unless otherwise stated. Then length and time are normalized by gravitational radius $r_g = GM/c^2$ and its light crossing time $t_g = r_g/c$, where $G$ is the gravitational constant and $M$ is the mass of a central object. We take $M = 1.4 M_\odot$ and $M = 10 M_\odot$ for NS and BH, respectively.

Basic equations consist of mass conservation,

$$\left(\rho u^\nu\right)_\nu = 0,$$

the energy-momentum conservation for magnetofluids,

$$T_{\mu\nu}^\rho = G_{\mu\nu},$$

the energy-momentum tensor for radiation field,

$$R_{\mu\nu} = -G_{\mu\nu},$$

and induction equation,

$$\partial_t (\sqrt{-g} B^\mu) = \left[\sqrt{-g} (B^\nu \partial_\nu B^\mu - B^\mu B^\nu)\right],$$

where $\rho$ is the proper mass density, $u^\mu$ is the four velocity, $v^\nu = u^\nu / u^0$ is the laboratory frame three velocity, $B^\mu$ is the laboratory frame magnetic three field, and $g = \det(g_{\mu\nu})$ is the determinant of the metric, $g_{\mu\nu}$.

The energy-momentum tensor for magnetofluid and radiation are given by

$$T_{\mu
ue} = (\rho + e + p_{\text{gas}} + 2p_{\text{mag}}) u^\mu u^\nu + (p_{\text{gas}} + p_{\text{mag}}) g_{\mu\nu} - b^\mu b^\nu,$$

$$R_{\mu\nu} = p_{\text{rad}} (4 u^\rho u^\nu_{\text{rad}} + g^{\rho\nu}),$$

where $p_{\text{gas}}$, $e$, $p_{\text{mag}}$, $p_{\text{rad}}$, and $u^\rho_{\text{rad}}$ are the gas pressure, gas internal energy, magnetic pressure, radiation pressure, and radiation frame’s four velocity. The gas internal energy is related to the gas pressure by $e = (\Gamma - 1) p_{\text{gas}}$ with $\Gamma = 5/3$ being the specific heat ratio. The magnetic four vector $b^\mu$ is related to its three vectors through $b^\mu = B^i h^i_{\mu}/u^0$, where $h^i_{\mu} = \delta^i_{\mu} + u^i u^\nu$ is the projection tensor and $\delta^i_{\mu}$ is the Kronecker delta. The magnetic pressure is represented by $p_{\text{mag}} = b_{\mu} b^\mu / 2$. 

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The gas and radiation field interact with each other through a radiation four force \( G^\mu \), which is represented by

\[
G^\mu = -\rho \kappa_{\text{abs}} (R^\mu_{\nu\sigma} u^\nu + 4\pi B u^\mu) - \rho \kappa_{\text{esc}} (R^\mu_{\nu\rho} u^\nu + R^\rho_{\nu\nu} u^\nu u^\rho) + G^\mu_{\text{comp}},
\]

where \( \kappa_{\text{abs}} = 6.4 \times 10^{22} \rho T_{\text{gas}}^{-3.5} \text{ cm}^{-1} \) and \( \kappa_{\text{esc}} = 0.4 \text{ cm}^{-1} \) are free–free absorption and Thomson-scattering opacities. The gas temperature is calculated by \( T_{\text{gas}} = \mu \rho p_{\text{gas}} / \rho k_B \), where \( \mu \) is the proton mass, \( k_B \) is the Boltzmann constant, and \( \mu = 0.5 \) is the mean molecular weight. The blackbody intensity is given by \( B = a_{\text{rad}} T_{\text{gas}}^4 \) with \( a_{\text{rad}} \) being the radiation constant. We included the thermal Comptonization as follows:

\[
G^\mu_{\text{comp}} = -\rho \kappa_{\text{esc}} \dot{E}_{\text{rad}} \frac{4k_B(T_e - T_{\text{rad}})}{m_e} \times \left[ 1 + 3.683 \left( \frac{k_B T_e}{m_e} \right) + 4 \left( \frac{k_B T_e}{m_e} \right)^2 \right]^{-1} \frac{k_B T_e}{m_e} u^\mu,
\]

where \( T_e \) is the electron temperature, \( \dot{E}_{\text{rad}} \) is the comoving frame radiation energy density, \( T_{\text{rad}} = (\dot{E}_{\text{rad}}/a_{\text{rad}})^{1/4} \) is the radiation temperature, and \( m_e \) is the electron rest mass (Sadowski et al. 2015). We take \( T_e = T_{\text{gas}} \) for simplicity.

We solve these equations in polar coordinate \((r, \theta, \phi)\) with Kerr–Schild metric by assuming axisymmetry with respect to the rotation axis, \( \theta = 0 \) and \( \pi \). The computational domain consists of \( r = r_{\text{in}} = 245 r_g, \theta = [0, \pi] \). Here we set the inner radius \( r_{\text{in}} \) to be 10 km for the NS and 0.98r_{\text{fl}} for the BH, where \( r_{\text{fl}} = M + (M^2 + a^2)^{1/2} \) is a horizon radius with \( a \) being the spin parameter. We take \( a = 0 \) in this paper. Numerical grid points are \((N_r, N_\theta, N_\phi) = (264, 264, 1) \). A radial grid size exponentially increases with radius, and a polar grid is given by \( \theta = \pi x_2 + (1 - h) \sin(2\pi x_2)/2 \), where \( h = 0.5 \) and \( x_2 \) spans uniformly from 0 and 1. We adopted the outgoing boundary at outer radius, and the reflective boundary is adopted at \( \theta = 0 \) and \( \pi \). At the inner boundary \( r = r_{\text{in}} \), a mirror symmetric boundary condition is employed for the case of the NS, while an outgoing boundary condition is used for the case of the BH. That is, the matter as well as the energy is not swallowed by the NS.

We start simulations from an equilibrium torus given by Fishbone & Moncrief (1976), but the gas pressure in this solution is replaced by a gas + radiation pressure by assuming a local thermodynamic equilibrium. The inner edge of the initial torus is situated at \( r = 20 r_g \), while its pressure maximum is situated at \( 33 r_g \). Weak poloidal magnetic fields are initially embedded in the torus. The magnetic flux vector \( A_\phi \) is proportional to \( \rho \), and a ratio of maximum \( p_{\text{mag}} \) and \( P_{\text{gas}} + P_{\text{rad}} \) is set to be 100. Outside the torus, the gas is not magnetized and the density and the pressure are given by \( \rho = 10^{-4} \rho_0 (r/r_g)^{-1.5} \) and \( p_{\text{gas}} = 10^{-6} \rho_0 (r/r_g)^{-2.5} \), where \( \rho_0 \) is the maximum mass density inside the torus. We also set \( p_{\text{rad}} = 10^{-10} \rho_0 \) and \( u_{\text{rad}} = (1, 0, 0, 0) \) outside the torus.

In this paper, we take \( \rho_0 = 0.1 \text{ g cm}^{-3} \) for the NS. On the other hand, the relatively small maximum mass density is employed for the BH \( (\rho_0 = 1.4 \times 10^{-2} \text{ g cm}^{-3}) \). By such adjustment, we can compare the models of NS and BH under almost equal conditions, since the mass of the NS is about one order of magnitude smaller than that of the BH. In the present work, we ignore the rotation of a central object \((a = 0)\). We also consider an unmagnetized NS. Thus we can directly study effects of the physical boundary at a surface of central objects by comparing results between the BH and NS.

### 3. Results

#### 3.1. Overview of the Two Cases

In the following, we show time averaged data between \( t = 3000 r_g - 5000 r_g \) at which the mass accretion continuously occurs onto a central object. We first give in Figure 1 global supercritical accretion flow patterns, comparing the two cases of NS accretion (left) and BH accretion (right). The color contours in Figure 1(a) represent gas density distribution with the same color scales (but note that the density normalizations \( \rho_0 \) is a factor of \( \sim 7 \) greater in the left panel), and arrows show fluid stream lines. White lines indicate photosphere measured from outer boundary at \( r = 245 r_g \) along fixed \( \theta \). The size of the NS (=10 km) corresponds to 4.8 \( r_g \) for a mass of \( 1.4 M_\odot \). Figure 1(b) shows stream lines around the NS (left) and BH (right). Red and blue lines indicate that the radial velocity is in positive and negative directions, respectively.

The flow patterns displayed in these figures are distinct in many respects. First of all, the flow lines are roughly conical (i.e., the line directions are more or less radial) in the innermost region (at \( r \lesssim 15 r_g \) in the BH case (see the right panel), while they are chaotic in the innermost region in the NS case (the left panel). Second, the high-density regions (indicated by the red color) are tightly collimated near the BH and thus have a conical structure in the BH accretion, while it is rather broadened and covers the large surface area of the NS. Third, we see more significant outflow motion in the NS case. In particular, the strong outflow is ejected even below the photosphere (indicated by the thick white line). The outflow has a large opening angle from \( \pm 60^\circ \) and its four velocity in the orthonormal frame is 0.2 around \( r \approx 60 r_g \) and \( \theta = 60^\circ \), while it is only 0.005 for the BH case. The mass flux is an order of magnitude larger for NS than that of BH. As a consequence, some of the inwardly flowing material in the NS accretion flow does not reach the NS surface but is reflected and turns its direction to outward. No such reflection motion is significant in the BH accretion flow (see Figure 1(b)). These differences should be understood in terms of the different mechanisms of absorbing radiation effects.

Figure 2 shows radial profiles of mass inflow rate \( \dot{M}_\text{in} \) (dashed), outflow rate \( \dot{M}_\text{out} \) (dotted), and net inflow rate \( \dot{M}_\text{net} = \dot{M}_\text{in} - \dot{M}_\text{out} \) (solid), for NS (red) and BH (black). For the NS, the mass inflow rate is about \( \dot{M}_\text{in} \approx 300 L_{\text{Edd}} \) around 10\( r_g \). It steeply decreases with a decrease in radius near the NS surface. Also the mass outflow rate has a similar trend with that of the inflow rate, but it is slightly smaller than the inflow rate. This indicates that substantial mass is blown away from the disk. We note that the mass supply (inflow) rate around \( r \approx 20 r_g \) is about 10\( 3 L_{\text{Edd}} \) in both cases, since we start from a similar initial torus. Even though the mass outflow rate is much higher for the NS than that for BH. Thus, it indicates that the NS can drive more massive outflows than the BH. We also note that the net inflow rate is approximately constant inside \( r \lesssim 15 r_g \) for the BH case. Thus, the inflow–outflow equilibrium is realized inside this radius. For the NS case, the net inflow rate is not constant but it slightly increases with increasing...
radius, even though the computational time is the same 
\( t = 3000-5000 \) \( rg \) in both simulations. This would be due to the mass accumulations on the NS as shown above (see also Figure 1). To summarize, a fraction of about a few tens of percent of the input mass can accrete onto a BH, whereas only 10% of less of the input mass can accrete onto an NS. The other fraction of mass is lost as outflow.

3.2. Various Energy Density Distributions

Next, we consider energy composition in the accretion disks with different central objects. The kinetic, gas, magnetic, and radiation energy densities are expressed as

\[
E_{\text{kin}} = \rho (\gamma - 1) \gamma, \tag{9}
\]

\[
E_{\text{gas}} = (e + p_{\text{gas}}) \gamma^2 - p_{\text{gas}}, \tag{10}
\]

\[
E_{\text{mag}} = b^2 \gamma^2 - (n_s b^n)^2, \tag{11}
\]

and

\[
E_{\text{rad}} = n_s n_g R^{\alpha \beta}, \tag{12}
\]

where \( n_s = (-\alpha, 0) \) is the normal observer’s four velocity, \( \alpha = (-g^{00})^{-1/2} \) is the lapse function, and \( \gamma = -n_s u^\alpha \) is the Lorentz factor. The energy density is normalized by \( \rho_0 \).

The left three panels in Figure 3 show spatial distributions of \( E_{\text{kin}}, E_{\text{mag}}, \) and \( E_{\text{rad}} \). Again, the conical flow structure around the BH is clearly shown in the lower panels of Figure 3, except for the magnetic energy distribution that shows a more spherically symmetric shape (see the second panel from the left). By contrast, the NS accretion case displayed in the upper panels shows a somewhat distinct pattern. The upper, third panel from the left, for example, shows that the large \( E_{\text{rad}} \) region is found more widely around the NS than around the BH. This indicates that there is intense radiation emitted from the NS surface and from the innermost flow region. Kinetic energy distribution displayed in the upper left panel shows a similar structure, implying the launch of outflow occurring widely from the surface of the accretion flow. Such enhanced energy regions around the central object are not found in the lower panels, since excess energy can be absorbed by the BH.

Right panel in Figure 3 shows the comoving frame radiation energy density distributions normalized by \( L_{\text{Ed}}/(4 \pi r^2 c) \), where we recover the light speed \( c \) for the sake of clarity. We found that this quantity largely exceeds unity, typically \( \sim 10^3 \) or even greater, in the entire inflow region. This is true in both of NS and BH cases, though the photon accumulation region is much wider in the former. This fact indicates that there exists a region full of radiation energy and that its radiation energy density is so high that it would be able to blow away the large amount of gas by counteracting the gravity force. Nevertheless, we find that the inflow region stably
persists around the compact objects. This is because the inflow exists deeply inside the photosphere (see Figure 1) so that the radiation flux can be much attenuated to become \( F_\text{rad} \ll E_\text{rad} c \). As a result, the gas is never prevented from accretion (Ohsuga & Mineshige 2007). This issue will be discussed again later.

Figure 4 shows the density weighted, angle-averaged energy densities in various forms along \( r \). We take an average of a physical quantity, \( f \), over the entire solid angle (\( \Omega \)) according to

\[
\langle f \rangle = \frac{\int d\Omega f \rho \sqrt{-g}}{\int d\Omega \rho \sqrt{-g}},
\]

where \( g = \det g_{\mu\nu} \).

Comparing these panels, we understand that the kinetic energy \( E_\text{kin} \) dominates over all other energy forms inside the accretion disks in both cases. While the radiation energy density \( E_\text{rad} \), the second largest one, increases with decreasing radius in both cases, there exists an interesting distinction between the two: the ratio of \( E_\text{rad} / E_\text{kin} \) increases with a decreasing radius near the central object in the NS accretion, while the opposite is the case in the BH accretion. In the proximity of the NS, especially, the radiation energy density is comparable to the kinetic energy density (see also Figure 3).

(Note that the kinetic energy is due mostly to the rotation, not the accretion.) These facts indicate that the radiation-pressure force makes a significant contribution in force balance near the NS (this point will be discussed in the next subsection). Around the BH, in contrast, the ratio of \( E_\text{rad} / E_\text{kin} \) stays nearly constant on the order of \( \sim 10\% \) but rather decreases in the innermost part. This is the direct consequence of photons being swallowed by the BH. We should note, however, that the difference between \( E_\text{kin} \) and \( E_\text{rad} \) may depend on the mass accretion rate.

The magnetic energy is unimportant in both cases; the ratio of \( E_\text{mag} / E_\text{kin} \) is always around a few percent. Likewise, the gas energy \( E_\text{gas} \) is everywhere negligible because the gas temperature is low enough. An interesting distinction between the BH and NS cases is found regarding the magnetic energy distribution; that is, it is nearly isotropic in the BH accretion, while it is concentrated on the polar and equatorial regions in the NS accretion (see Figure 3). In our simulations, we start from the poloidal magnetic field. The magnetic flux is swept up according to the gas accretion and it is accumulated near the central object. Since we assume ideal MHD and axisymmetry, the magnetic field is dissipated by a small numerical resistivity and most of the flux remains around the pole.

Figure 3. First three panels from the left show kinetic, magnetic, and radiation energy density. The right panel shows the comoving frame radiation energy density normalized by \( L_\text{Edd} / (4\pi c^2) \). Top and bottom panels correspond to the cases for the neutron star and black hole, respectively.
3.3. Force Balance on the Equatorial Plane

Next, we show the radial profile of forces acting on the fluid elements. We consider a steady-state equation of motion

\[
f_r^{\text{adv}} + f_r^{\text{grav}} + f_r^{\text{cent}} + f_r^{\text{rad}} + f_r^{\text{gas}} + f_r^{\text{mag}} + f_r^{\text{cor}} = 0,
\]

Here \( f_r^{\text{adv}}, f_r^{\text{grav}}, f_r^{\text{cent}}, f_r^{\text{rad}}, f_r^{\text{gas}}, f_r^{\text{mag}}, f_r^{\text{cor}} \) are defined according to Moller & Sadowski (2015) as

\[
f_r^{\text{adv}} = -u^i \partial_i u_r,
\]

\[
f_r^{\text{grav}} = \gamma_r \frac{T_r^i}{w} \Gamma_{ir},
\]

\[
f_r^{\text{cent}} = \gamma_r \frac{T_r^i}{w} \Gamma_{ir}',
\]

\[
f_r^{\text{rad}} = \frac{G_r}{w},
\]

\[
f_r^{\text{gas}} = -\frac{\partial_r p_{\text{gas}}}{w},
\]

\[
f_r^{\text{mag}} = -\frac{\partial_r (b^2/2) + \partial_i (b^i b_r)}{w},
\]

\[
f_r^{\text{cor}} = f_r^{\text{metric}} - f_r^{\text{grav}} - f_r^{\text{cent}} + f_r^{\text{cent}},
\]

where

\[
f_r^{\text{metric}} = \frac{1}{w} T_r^{\Gamma} \gamma_{ir} - \frac{T_r^i}{w} \rho u^j \partial_i \sqrt{-g},
\]

\[
f_r^{\text{cent}} = -\frac{u_j}{w} \partial_j [(w - \rho) u^i],
\]

where \( w = \rho + e + p_{\text{gas}} + 2p_{\text{mag}} \) denotes the relativistic enthalpy. Here Equations (15)–(20) correspond to advection term, gravity force, centrifugal force, radiation force, gas pressure gradient force, and Lorentz force. \( f_r^{\text{cor}} \) is the relativistic correction term.

Figure 4 shows various forms of density weighted, angle-averaged radial force along \( r \) normalized by the gravity force. Here \( f_r^{\text{tot}} \) is the total force without gravity force, so that steady accretion would be realized, where \( |f_r^{\text{tot}}|/|f_r^{\text{grav}}| = 1 \). Let us first examine the NS case displayed in the upper panel. We immediately notice that the centrifugal force balances almost completely with the gravity force at large radii far from the central object. Hence, the rotation profile is nearly Keplerian and radiation force is negligible there. With a decrease in radius, however, the outward radiation force grows, since the NS surface cannot swallow the radiation so that the radiation energy is accumulated there. The radiation energy density profile, hence, has a negative gradient along \( r \), which gives rise to outward radiation-pressure force. The centrifugal force decreases with a decreasing radius so that the radiation force and centrifugal force can be comparable close to the NS surface. This occurs because the gravitational attraction force by the NS is weakened by the outward radiation-pressure force. As a result, the disk rotation becomes highly sub-Keplerian, although the flow is still in a quasi-steady state. The important fact is that radiation force does never exceed the gravitational force, which makes it feasible to induce supercritical accretion flow.

It is then of great importance to pay attention to the behavior of the centrifugal force. We find a clear tendency for it to decline inward very close to the NS. This is caused by the accumulation of low angular momentum above the NS surface and never happens in the BH case, since matter should be immediately swallowed. However, the gradient of the radiation energy density is not large enough to totally compensate for the gravitational attraction force toward the NS. Finally, the advection term is very small, compared to the gravity force, but it does not vanish. That is, the matter is slowly accreting onto the NS surface with accretion velocity being much less than the free-fall velocity. We may call this slowly accreting zone (at \( r < 10g \)) the settling region. As a result, the supercritical accretion is feasible for the NS.
Next, let us examine the force balance in the BH accretion case in comparison with the NS case. A big distinction is found in the behavior of the radiation force, which is negative in the BH case, while it was positive in the NS case. This is because not only the gas but also the radiation energy is swallowed by the BH. The negative radiation flux pushes the gas toward the BH. This explains why supercritical accretion onto a BH is feasible (see Ohsuga & Mineshige 2007 for the discussion based on the pseudo-Newtonian dynamics).

Another distinction is that there is no force balance near the BH in the sense that the total force no longer balances with the gravity force near the BH. This means that mass continuously falls onto the BH with finite velocity. In particular, the accretion motion is supersonic and is close to the speed of light in the BH vicinity.

We note that the centrifugal force exceeds the gravity force inside \( r < 10r_g \) for BH, but the total force balance holds if we consider the relativistic correction factor \( f_{\text{cor}} \), i.e., a quasi-steady state is actually realized. There is an issue regarding how we decompose each force term in Equation (14). The centrifugal force \( f_{\text{cent}} \) approaches the non-relativistic one far from the BH, but this force does not balance with gravity force everywhere. It deviates from the gravity force close to the central object. The relativistic correction term \( f_{\text{cor}} \) is important in this region. For example, the innermost stable circular orbit is never obtained without \( f_{\text{cor}} \). The gravity force almost balances with the centrifugal force and the correction force in this region, but the advection and radiation forces are also important and thus the total force balances with the gravity force.

4. Discussion

In the present paper, we have carefully examine the gas dynamics of supercritical flow around the NS, in comparison with that around the BH, through the GR-RMHD simulations. Supercritical accretion is feasible in both NS and BH cases but for distinct reasons. While it is photon trapping that works in the BH case, the removal of mass and energy in the form of intense outflow is a key to realizing supercritical accretion onto the NS. The flow dynamics is also distinct: sub-sonic, settling flow occurs around the NS surface, whereas matter nearly free falls onto the BH. In the following, we will discuss some related issues more or less related to supercritical NS accretion.

4.1. Outflow from Inside the Spherization Radius

It is widely known that SS73 have proposed the standard disk model, but in the same paper they also made pioneering discussions regarding the gas dynamics of the supercritical accretion flow onto the BH. In their section IV, SS73 introduced the notion of the spherization radius, inside which gas flows toward the central BH in a spherically symmetric fashion. They also pointed out that outflow emerges from inside this radius. They evaluated the spherization radius to be on the order of \( r_{\text{ph}} \sim 10(Mc^2/L_{\text{edd}})r_g \), corresponding to the trapping radius, inside which photon trapping is significant (see also Begelman & Meier 1982). In the present case, we estimate \( r_{\text{ph}} \sim 10^2r_g \) for \( Mc^2/L_{\text{edd}} \gtrsim 300 \), see Figure 2) thus being far outside the picture box of Figure 1.

The right panel of Figure 1 clearly shows that the inflow and outflow stream lines are separated all the way down to the BH event horizon. In other words, there are no stream lines, which turns its direction from inward to outward. By contrast, the left panel of Figure 1 shows somewhat similar stream lines as those illustrated in Figure 8 of SS73; that is, some stream lines change their directions from inward to outward. Rather, we see that the change of the direction occurs even in the very vicinity of the NS surface. In fact, the inflow and outflow rates nearly coincide in the innermost region (inside \( \sim 10r_g \), see, Figure 2) so that the net accretion rate is kept around the critical rate. This is exactly a situation as that postulated by SS73.

4.2. Bernoulli Parameter

To visualize the relative importance of the outflow in the NS accretion, we calculate the local Bernouilli parameter according to Sadowski & Narayan (2015b);

\[
B_e \equiv -\frac{T_i - R_i + \rho u'_r}{\rho u'}
\]

where \( T_i \) and \( R_i \) are the \( t - r \) components of the MHD and radiation energy–momentum tensors (representing the energy flux of MHD and radiation processes), respectively, and \( \rho u' \) stands for the rest-mass energy flux.

The results are shown in Figure 6 for the NS and BH cases in the left and right panels, respectively. The locations of the photospheres are also indicated by the thick white lines there. It is obvious that the blue regions, in which \( B_e < 0 \), are wider in the BH case. Especially, we find that the Bernoulli parameter is negative mostly below the photosphere close to the BH, while it is positive in the NS case (except near the equatorial plane).

4.3. Radiation Cushion

The next question we wish to address is if there exists a settling regime covering the NS surface. The accretion column created on the magnetized NS surface is composed of the upper free-fall region and the lower settling region (e.g., Basko & Sunyaev 1976; Kawashima et al. 2016). In the latter, accretion velocity is greatly reduced by the decelerating force asserted by the radiation cushion.
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The direct consequence of the existence of the settling region is that the matter density is \( \rho \propto r^{-3} \), radiation pressure is \( P_{\text{rad}} \propto r^{-4} \), and radiation temperature is \( T_{\text{rad}} \propto r^{-3} \). These relations are derived from the hydrostatic balance in the radiation-pressure-dominated atmosphere, which leads to

\[
\frac{GM\rho}{r^2} = -\frac{dP_{\text{rad}}}{dr}
\]

where \( r \) is the radial coordinate. Here, we assume that accretion motion is very slow (accretion velocity is much less than free-fall velocity). Let us further assume little entropy production is significant during the accretion. Then, the adiabatic relation holds between \( P_{\text{rad}} \) and matter density \( \rho \); that is, \( P_{\text{rad}} \propto \rho^{4/3} \). We then find \( dP_{\text{rad}}/P_{\text{rad}}^{4/3} \propto dr/r^2 \), which reads \( P_{\text{rad}} \propto r^{-4} \) and \( \rho \propto P_{\text{rad}}^{3/4} \propto r^{-3} \) (see also Burger & Katz 1983).

To see if such dependences appear in the simulation data of the NS case, we plot matter density and \( T_{\text{rad}} \) as functions of radii in Figure 7. We find that radiation entropy crudely obeys the expected relationship; \( T_{\text{rad}} \propto r^{-3} \) in the innermost region, \( r < 10r_\text{s} \), although the density profile is steeper than \( r^{-3} \). These results indicate that an almost adiabatic settling region is formed close to the NS. The mass density and radiation entropy on the surface of NS increase with time due to the accumulation. Nevertheless, their radial profiles do not change. This indicates that the force balance given in Equation (25) holds during simulation interval. Thus, we can expect that supercritical accretion onto the NS continues to accompany the forming settling region, until the gas in the disk is exhausted and the mass accretion rate decreases.

4.4. Validity of Our Numerical Model

We simply compute opacities assuming fully ionized hydrogen gas. The free–free opacity is, however, much larger by assuming the solar opacity. We expect results would not be affected so much by the metallicity since the local thermodynamic equilibrium \( (T_{\text{gas}} = T_{\text{rad}}) \) is attained mainly by to the Comptonization, for which the cooling timescale is much shorter for the supercritical accretion disks. For the scattering opacity, it decreases about 15%, assuming the solar abundance. The reduction of opacity might reduce the outflow power. But the outflow velocity is determined by the balance between the radiation force \( (\propto k_{\text{sc}c}F_{\text{rad}}) \) and its drag force \( (\propto k_{\text{sc}c}F_{\text{rad}} \), see Takahashi & Ohsuga 2015). The resulting terminal velocity would not be affected by the opacity. Also Ohsuga et al. (2005) shows that the luminosity weakly depends, or is almost independent from, the metallicity. Thus, our conclusion would hold even if we adopted the solar metallicity.

Another concern in our numerical model is the boundary condition on the NS. We simply applied a mirror boundary condition where the gas never flows across the boundary. This boundary condition might be plausible to mimic the NS’s solid surface, while other boundary conditions have been adopted in the past study; e.g., the free boundary condition (Romanova et al. 2012) or the accretion-energy-injection boundary condition (Ohsuga 2007). Also the boundary condition adopted in our simulation does not take into account the interaction between the gas and NS. The magnetic activity in this boundary layer can transport the angular momentum (Armitage 2002). The dissipation of rotation energy of the disk would increase the radiation energy close to the NS. Although recent high-resolution MHD simulations show that the stresses worked in the boundary layer oscillate around zero (Pessah & Chan 2012; Belyaev & Quataert 2017), it is under debate what boundary condition is appropriate to describe the NS surface. We have to perform a comprehensive study around the NS surface with different boundary condition models to investigate the plausible boundary conditions. We leave this problem for important future work.

5. Conclusions

We performed a two-dimensional axisymmetric GR-RMHD simulation of supercritical accretion onto a nonrotating unmagnetized NS, and compared results with nonrotating BHs. Our findings can be summarized as follows:

1. In contrast with the BH case, a significant fraction of mass is blown away by the radiation-pressure-driven outflow and thus the net mass inflow rate reduces for the NS case. Also the anisotropic radiation arising from the anisotropic density distribution helps photons escape from the disk.
2. Inside the accretion disks, the radiation flux is largely attenuated so that the radiation force balances with the sum of centrifugal and gravity forces. Due to the large optical depth in the supercritical disks, the radiation energy density much exceeds that expected from the Eddington luminosity, \( E_{\text{rad}} \propto F_{\text{rad}}c^2/\tau > 100L_{\text{Edd}}/(4\pi r^2c) \).
3. We found that the gas and radiation is accumulated on the NS surface. The settling region, where accretion motion is significantly decelerated by the radiation cushion, is formed. The radiation cushion would be approximately adiabatic, i.e., the radiation energy roughly follows \( \dot{E} \propto r^{-4} \) and the gas and radiation temperature obeys \( \propto r^{-1} \). Such a radiation cushion never appears around the BH so that matter can be directly swallowed by the BH. Also, these mass density and radiation energy density profiles follow radiation-pressure-supported hydrostatic balance.

These facts make supercritical accretion feasible for the NS.

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