Accretion Power in Astrophysics

Third Edition

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1 Accretion as a source of energy

1.1 Introduction

For the nineteenth century physicists, gravity was the only conceivable source of energy in celestial bodies, but gravity was inadequate to power the Sun for its known lifetime. In contrast, at the beginning of the twenty-first century it is to gravity that we look to power the most luminous objects in the Universe, for which the nuclear sources of the stars are wholly inadequate. The extraction of gravitational potential energy from material which accretes on to a gravitating body is now known to be the principal source of power in several types of close binary systems, and is widely believed to provide the power supply in active galactic nuclei and quasars. This increasing recognition of the importance of accretion has accompanied the dramatic expansion of observational techniques in astronomy, in particular the exploitation of the full range of the electromagnetic spectrum from the radio to X-rays and γ-rays. At the same time, the existence of compact objects has been placed beyond doubt by the discovery of the pulsars, and black holes have been given a sound theoretical status. Thus, the new role for gravity arises because accretion on to compact objects is a natural and powerful mechanism for producing high-energy radiation.

Some simple order-of-magnitude estimates will show how this works. For a body of mass M and radius R_* the gravitational potential energy released by the accretion of a mass m on to its surface is

$$\Delta E_{\rm acc} = GMm/R_* \tag{1.1}$$

where G is the gravitation constant. If the accreting body is a neutron star with radius $R_* \sim 10\,\rm km$, mass $M \sim M_\odot$, the solar mass, then the yield $\Delta E_{\rm acc}$ is about 10^{20} erg per accreted gram. We would expect this energy to be released eventually mainly in the form of electromagnetic radiation. For comparison, consider the energy that could be extracted from the mass m by nuclear fusion reactions. The maximum is obtained if, as is usually the case in astrophysics, the material is initially hydrogen, and the major contribution comes from the conversion, (or 'burning'), of hydrogen to helium. This yields an energy release

$$\Delta E_{\rm nuc} \;=\; 0.007mc^2 \tag{1.2}$$

where c is the speed of light, so we obtain about 6×10^{18} erg g^{-1} or about one twentieth of the accretion yield in this case.

It is clear from the form of equation (1.1) that the efficiency of accretion as an energy release mechanism is strongly dependent on the compactness of the accreting object: the larger the ratio M/R_*, the greater the efficiency. Thus, in treating accretion on to objects of stellar mass we shall certainly want to consider neutron stars ($R_* \sim 10$ km) and black holes with radii $R_* \sim 2GM/c^2 \sim 3(M/M_\odot)$ km (see Section 7.7). For white dwarfs with $M \sim M_\odot$, $R_* \sim 10^9$ cm, nuclear burning is more efficient than accretion by factors 25–50. However, it would be wrong to conclude that accretion on to white dwarfs is of no great importance for observations, since the argument takes no account of the timescale over which the nuclear and accretion processes act. In fact, when nuclear burning does occur on the surface of a white dwarf, it is likely that the reaction tends to 'run away' to produce an event of great brightness but short duration, a nova outburst, in which the available nuclear fuel is very rapidly exhausted. For almost all of its lifetime no nuclear burning occurs, and the white dwarf (may) derive its entire luminosity from accretion. Binary systems in which a white dwarf accretes from a close companion star are known as *cataclysmic variables* and are quite common in the Galaxy. Their importance derives partly from the fact that they provide probably the best opportunity to study the accretion process in isolation, since other sources of luminosity, in particular the companion star, are relatively unimportant.

For accretion on to a 'normal', less compact, star, such as the Sun, the accretion yield is smaller than the potential nuclear yield by a factor of several thousand. Even so, accretion on to such stars may be of observational importance. For example, a binary system containing an accreting main-sequence star has been proposed as a model for the so-called symbiotic stars.

For a fixed value of the compactness, M/R_*, the luminosity of an accreting system depends on the rate $\dot{M}$ at which matter is accreted. At high luminosities, the accretion rate may itself be controlled by the outward momentum transferred from the radiation to the accreting material by scattering and absorption. Under certain circumstances, this can lead to the existence of a maximum luminosity for a given mass, usually referred to as the Eddington luminosity, which we discuss next.

1.2 The Eddington limit

Consider a steady spherically symmetrical accretion; the limit so derived will be generally applicable as an order-of-magnitude estimate. We assume the accreting material to be mainly hydrogen and to be fully ionized. Under these circumstances, the radiation exerts a force mainly on the free electrons through Thomson scattering, since the scattering cross-section for protons is a factor $(m_\mathrm{e}/m_\mathrm{p})^2$ smaller, where $m_\mathrm{e}/m_\mathrm{p} \cong 5 \times 10^{-4}$ is the ratio of the electron and proton masses. If S is the radiant energy flux (erg s^{-1}cm^{-2}) and $\sigma_\mathrm{T} = 6.7 \times 10^{-25}$ cm^2 is the Thomson cross-section, then the outward radial force on each electron equals the rate at which it absorbs momentum, $\sigma_\mathrm{T} S/c$. If there is a substantial population of elements other than hydrogen, which have retained some bound electrons, the effective cross-section, resulting from the absorption of photons in spectral lines, can exceed σ_T considerably. The attractive electrostatic Coulomb force between the electrons and protons means that as they move out the electrons drag the protons with them. In effect, the radiation pushes out electron–proton pairs against the total gravitational force $GM(m_p + m_e)/r^2 \cong GMm_p/r^2$ acting on each pair at a radial distance r from the centre. If the luminosity of the accreting source is L(erg s^{-1}), we have $S = L/4\pi r^2$ by spherical symmetry, so the net inward force on an electron–proton pair is

$$\left(GMm_p - \frac{L\sigma_T}{4\pi c}\right)\frac{1}{r^2}.$$

There is a limiting luminosity for which this expression vanishes, the Eddington limit,

$$L_{\rm Edd} = 4\pi GMm_p c/\sigma_T \tag{1.3}$$

$$\cong 1.3 \times 10^{38}(M/M_\odot)\,\text{erg s}^{-1}. \tag{1.4}$$

At greater luminosities the outward pressure of radiation would exceed the inward gravitational attraction and accretion would be halted. If all the luminosity of the source were derived from accretion this would switch off the source; if some, or all, of it were produced by other means, for example nuclear burning, then the outer layers of material would begin to be blown off and the source would not be steady. For stars with a given mass–luminosity relation this argument yields a maximum stable mass.

Since $L_{\rm Edd}$ will figure prominently later, it is worth recalling the assumptions made in deriving expressions (1.3,1.4). We assumed that the accretion flow was *steady* and *spherically symmetric*. A slight extension can be made here without difficulty: if the accretion occurs only over a fraction f of the surface of a star, but is otherwise dependent only on radial distance r, the corresponding limit on the accretion luminosity is $fL_{\rm Edd}$. For a more complicated geometry, however, we cannot expect (1.3,1.4) to provide more than a crude estimate. Even more crucial was the restriction to *steady* flow. A dramatic illustration of this is provided by supernovae, in which $L_{\rm Edd}$ is exceeded by many orders of magnitude. Our other main assumptions were that the accreting material was largely hydrogen and that it was fully ionized. The former is almost always a good approximation, but even a small admixture of heavy elements can invalidate the latter. Almost complete ionization is likely to be justified however in the very common case where the accreting object produces much of its luminosity in the form of X-rays, because the abundant ions can usually be kept fully stripped of electrons by a very small fraction of the X-ray luminosity. Despite these caveats, the Eddington limit is of great practical importance, in particular because certain types of system show a tendency to behave as 'standard candles' in the sense that their typical luminosities are close to their Eddington limits.

For accretion powered objects the Eddington limit implies a limit on the steady accretion rate, $\dot{M}$(g s^{-1}). If all the kinetic energy of infalling matter is given up to radiation at the stellar surface, R_*, then from (1.1) the *accretion luminosity* is

$$L_{\rm acc} = GM\dot{M}/R_*. \tag{1.5}$$

It is useful to re-express (1.5) in terms of typical orders of magnitude: writing the accretion rate as $\dot{M} = 10^{16}\dot{M}_{16}\,\mathrm{g\ s^{-1}}$ we have

$$\begin{aligned} L_{\mathrm{acc}} &= 1.3 \times 10^{33}\dot{M}_{16}(M/M_{\odot})(10^9\,\mathrm{cm}/R_*)\,\mathrm{erg\ s^{-1}} && (1.6)\\ &= 1.3 \times 10^{36}\dot{M}_{16}(M/M_{\odot})(10\,\mathrm{km}/R_*)\,\mathrm{erg\ s^{-1}}. && (1.7)\end{aligned}$$

The reason for rewriting (1.5) in this way is that the quantities $(M/M_{\odot})$, $(10^9\,\mathrm{cm}/R_*)$ and $(M/M_{\odot})$, $(10\ \mathrm{km}/R_*)$ are of order unity for white dwarfs and neutron stars respectively. Since $10^{16}\ \mathrm{g\ s^{-1}}(\sim 1.5 \times 10^{-10} M_{\odot}\ \mathrm{yr^{-1}})$ is a typical order of magnitude for accretion rates in close binary systems involving these types of star, we have $\dot{M}_{16} \sim 1$ in (1.6,1.7), and the luminosities $10^{33}\ \mathrm{erg\ s^{-1}}$, $10^{36}\ \mathrm{erg\ s^{-1}}$ represent values commonly found in such systems. Further, by comparison with (1.4) it is immediately seen that for steady accretion $\dot{M}_{16}$ is limited by the values $\sim 10^5$ and 10^2 respectively. Thus, accretion rates must be less than about $10^{21}\ \mathrm{g\ s^{-1}}$ and $10^{18}\ \mathrm{g\ s^{-1}}$ in the two types of system if the assumptions involved in deriving the Eddington limit are valid.

For the case of accretion on to a black hole it is far from clear that (1.5) holds. Since the radius does not refer to a hard surface but only to a region into which matter can fall and from which it cannot escape, much of the accretion energy could disappear into the hole and simply add to its mass, rather than be radiated. The uncertainty in this case can be parametrized by the introduction of a dimensionless quantity η, the *efficiency*, on the right hand side of (1.5):

$$\begin{aligned} L_{\mathrm{acc}} &= 2\eta GM\dot{M}/R_* && (1.8)\\ &= \eta \dot{M}c^2 && (1.9)\end{aligned}$$

where we have used $R_* = 2GM/c^2$ for the black hole radius. Equation (1.9) shows that η measures how efficiently the rest mass energy, c^2 per unit mass, of the accreted material is converted into radiation. Comparing (1.9) with (1.2) we see that $\eta = 0.007$ for the burning of hydrogen to helium. If the material accreting on to a black hole could be lowered into the hole infinitesimally slowly - scarcely a practical proposition - all of the rest mass energy could, in principle, be extracted and we should have $\eta = 1$. As we shall see in Chapter 7 the estimation of realistic values for η is an important problem. A reasonable guess would appear to be $\eta \sim 0.1$, comparable to the value $\eta \sim 0.15$ obtained from (1.8) for a solar mass neutron star. Thus, despite its extra compactness, a stellar mass black hole may be no more efficient in the conversion of gravitational potential energy to radiation than a neutron star of similar mass.

As a final illustration here of the use of the Eddington limit we consider the nuclei of active galaxies and the closely related quasars. These are probably the least understood class of object for which accretion is thought to be the ultimate source of energy. The main reason for this belief comes from the large luminosities involved: these systems may reach $10^{47}\ \mathrm{erg\ s^{-1}}$, or more, varying by factors of order 2 on timescales of weeks, or less. With the nuclear burning efficiency of only $\eta = 0.007$, the rate at which mass is processed in the source could exceed $250\ M_{\odot}\ \mathrm{yr^{-1}}$. This is a rather severe requirement and it is clearly greatly reduced if accretion with an efficiency $\eta \sim 0.1$ is postulated instead. The accretion rate required is of order 20 $M_\odot$ yr^{-1}, or less, and rates approaching this might plausibly be provided by a number of the mechanisms considered in Chapter 7. If these systems are assumed to radiate at less than the Eddington limit, then accreting masses exceeding about $10^9 M_\odot$ are required. White dwarfs are subject to upper limits on their masses of 1.4 $M_\odot$ and neutron stars cannot exceed about 3 $M_\odot$ thus, only massive black holes are plausible candidates for accreting objects in active galactic nuclei.

1.3 The emitted spectrum

We can now make some order-of-magnitude estimates of the spectral range of the emission from compact accreting objects, and, conversely, suggest what type of compact object may be responsible for various observed behaviour. We can characterize the continuum spectrum of the emitted radiation by a temperature T_{rad} defined such that the energy of a typical photon, $h\bar{\nu}$, is of order kT_{rad}, $T_{\mathrm{rad}} = h\bar{\nu}/k$, where we do not need to make the choice of $\bar{\nu}$ precise. For an accretion luminosity L_{acc} from a source of radius R, we define a blackbody temperature T_{b} as the temperature the source would have if it were to radiate the given power as a blackbody spectrum:

$$T_{\mathrm{b}} = (L_{\mathrm{acc}}/4\pi R_*^2 \sigma)^{1/4}. \tag{1.10}$$

Finally, we define a temperature T_{th} that the accreted material would reach if its gravitational potential energy were turned entirely into thermal energy. For each proton–electron pair accreted, the potential energy released is $GM(m_{\mathrm{p}} + m_{\mathrm{e}})/R_* \cong GMm_{\mathrm{p}}/R_*$, and the thermal energy is $2 \times \frac{3}{2}kT$; therefore

$$T_{\mathrm{th}} = GMm_{\mathrm{p}}/3kR_*. \tag{1.11}$$

Note that some authors use the related concept of the virial temperature, $T_{\mathrm{vir}} = T_{\mathrm{th}}/2$, for a system in mechanical and thermal equilibrium. If the accretion flow is optically thick, the radiation reaches thermal equilibnum with the accreted material before leaking out to the observer and $T_{\mathrm{rad}} \sim T_{\mathrm{b}}$. On the other hand, if the accretion energy is converted directly into radiation which escapes without further interaction (i.e. the intervening material is optically thin), we have $T_{\mathrm{rad}} \sim T_{\mathrm{th}}$. This occurs in certain types of shock wave that may be produced in some accretion flows and we shall see in Chapter 3 that (1.11) provides an estimate of the shock temperature for such flows. In general, the radiation temperature may be expected to lie between the thermal and blackbody temperatures, and, since the system cannot radiate a given flux at less than the blackbody temperature, we have

$$T_{\mathrm{b}} \lesssim T_{\mathrm{rad}} \lesssim T_{\mathrm{th}}.$$

Of course, these estimates assume that the radiating material can be characterized by a single temperature. They need not apply, for example, to a non-Maxwellian distribution of electrons radiating in a fixed magnetic field, such as we shall meet in Chapter 9.

Let us apply the limits (1.10), (1.11) to the case of a solar mass neutron star. The upper limit (1.11) gives $T_{\rm th} \sim 5.5 \times 10^{11}$ K, or, in terms of energies, $kT_{\rm th} \sim$50 MeV. To evaluate the lower limit, $T_{\rm b}$, from (1.10), we need an idea of the accretion luminosity, $L_{\rm acc}$; but $T_{\rm b}$ is, in fact, very insensitive to the assumed value of $L_{\rm acc}$, since it is proportional to the fourth root. Thus we can take $L_{\rm acc} \sim L_{\rm Edd} \sim 10^{38}$ erg s^{-1} for a rough estimate; if, instead, we were to take a typical value $\sim 10^{36}$ erg s^{-1} (equation (1.10)) this would change $T_{\rm b}$ only by a factor of ~ 3. We obtain $T_{\rm b} \sim 10^7$K or $kT_{\rm b} \sim$1 keV, and so we expect photon energies in the range

$$1 \text{ keV} \lesssim h\bar{\nu} \lesssim 50 \text{ MeV}$$

as a result of accretion on to neutron stars. Similar results would hold for stellar mass black holes. Thus we can expect the most luminous accreting neutron star and black hole binary systems to appear as medium to hard X-ray emitters and possibly as γ-ray sources. There is no difficulty in identifying this class of object with the luminous galactic X-ray sources discovered by the first satellite X-ray experiments, and added to by subsequent investigations.

For accreting white dwarfs it is probably more realistic to take $L_{\rm acc} \sim 10^{33}$ erg s^{-1} in estimating $T_{\rm b}$ (cf. (1.6)). With $M = M_\odot$, $R_* = 5 \times 10^8$ cm, we obtain

$$6 \text{ eV} \lesssim h\bar{\nu} \lesssim 100 \text{ keV}.$$

Consequently, accreting white dwarfs should be optical, ultraviolet and possibly X-ray sources. This fits in neatly with our knowledge of cataclysmic variable stars, which have been found to have strong ultraviolet continua by the Copernicus and IUE satellite experiments. In addition, some of them are now known to emit a small fraction of their luminosity as thermal X-ray sources. We shall see that in many ways cataclysmic variables are particularly useful in providing observational tests of theories of accretion.

1.4 Accretion theory and observation

So far we have discussed the amount of energy that might be expected by the accretion process, but we have made no attempt to describe in detail the flow of accreting matter. A hint that the dynamics of this flow may not be straightforward is provided by the existence of the Eddington limit, which shows that, at least for high accretion rates, forces other than gravity can be important. In addition, it will emerge later that, certainly in many cases and probably in most, the accreting matter possesses considerable angular momentum per unit mass which, in realistic models, it has to lose in order to be accreted at all. Furthermore, we need a detailed description of the accretion flow if we are to explain the observed spectral distribution of the radiation produced: crudely speaking, in the language of Section 1.3, we want to know whether $T_{\rm rad}$ is closer to $T_{\rm b}$ or $T_{\rm th}$.

The two main tools we shall use in this study are the equations of gas dynamics and the physics of plasmas. We shall give a brief introduction to gas dynamics in Sections 2.1–2.4 of the next chapter, and treat some aspects of plasma physics in Chapter 3. In addition, the elements of the theory of radiative transfer are summarized in the Appendix. The reader who is already familiar with these subjects can omit these parts of the text. The rest of the book divides into three somewhat distinct parts. First, in Chapters 4 to 6 we consider accretion by stellar mass objects in binary systems. In these cases, we often find that observations provide fairly direct evidence for the nature of the systems. For example, there is sometimes direct evidence for the importance of angular momentum and the existence of accretion discs. This contrasts greatly with the subsequent discussion of active galactic nuclei in Chapters 7 to 10. Here, the accretion theory arises at the end of a sequence of plausible, but not unproblematic, inductions. Furthermore, there appears to be no absolutely compelling evidence for, or against, the existence of accretion *discs* in these systems. Thus, whereas we normally use the observations of stellar systems to test the theory, for active nuclei we use the theory, to some extent, to illustrate the observations. This is particularly apparent in the final part of the book, where, in Chapters 9 and 10, we discuss two quite different models for powering an active nucleus by an accretion disc around a supermassive black hole. Finally, in Chapter 11 we review all possible accretion flows, most of which have already been studied in earlier chapters, classifying them according to which physical effects dominate their properties and behaviour. We also describe in some detail recent advances in our understanding of accretion flows, with particular emphasis on the class of advection dominated accretion flows or ADAFs.