

# Compact Object Astrophysics

Jacopo Tissino

2020-10-13

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## Introduction

Tuesdays and Wednesday at 14.30 PM in room P1A, Paolotti building. 22 people.

This course overlaps with “Computational Astrophysics” by professor Mapelli.

The examination is an oral one, done either online or live.

We start with a brief overview of the final fates of massive stars. We have white dwarfs, neutron stars and black holes under the category of “compact objects”, but white dwarfs are not really that compact.

We then discuss accretion onto compact objects, and neutron stars. An open question: what is the EOS of ultradense neutron matter?

“Accretion power in astrophysics”, “The physics of Compact Objects”, “Astrofisica Relativistica I & II”, “Astrofisica delle Alte Energie”.

## 0.1 A journey into the life of a massive star

Stars whose mass is  $M \gtrsim 8M_{\odot}$  go supernova at the end of their life. During their lifetime, hydrogen fuses through two channels: the p-p chain and the CNO cycle.

In the CNO cycle, four protons turn into a  ${}^4\text{He}$  nuclide, two positrons, two electron neutrinos using heavier nuclides as catalysts.

The critical temperature above which the CNO cycle dominates is around  $T_c \sim 2 \times 10^7$  K. For the Sun, less than 8 % of energy production is through the CNO cycle.

When the temperature of the core reaches a value around  $1 \div 2 \times 10^8$  K and the density is around  $10^8 \div 10^9$  g/cm<sup>3</sup>, helium starts to burn in the  $3\alpha$  process, becoming  ${}^{12}\text{C}$ . The Q-value here is around 7.27 MeV.

As soon as we have carbon, this can fuse with an  $\alpha$  particle giving rise to a nucleus of oxygen with  $Q \approx 7.16$  MeV. This oxygen can further catch an  $\alpha$  particle, making a  ${}^{20}\text{Ne}$  nuclide.

Then, we have a temperature around  $5 \times 10^8$  K and a density around  $3 \times 10^6$  g/cm<sup>3</sup>. Carbon starts to fuse with itself, making sodium, magnesium, and more neon (plus an  $\alpha$  particle).

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If you want carbon burning to proceed in a steady way, it must occur in a nondegenerate electron gas. This occurs only if the star is quite massive, more than  $8M_{\odot}$ . Otherwise, it is an explosive process.

Now the core temperature reaches  $10^9$  K. The energy of a typical photon is quite high,  $h\nu \sim k_B T \sim 100$  keV. Suppose there are neon nuclei in the core (this will be the case since they are a product of fusion).

It is not hard for a Neon to lose an  $\alpha$  particle through photodissociation, this produces an Oxygen. The energy required for this is of the order 4.7 MeV, at the high energy tail we have a few photons at this energy.

This is the “neon burning phase”, after which we have oxygen and magnesium. Oxygen is the next candidate for nuclear burning, and after a further contraction the star starts burning it. It fuses with itself to produce  $^{28}\text{Si}$  plus an  $\alpha$ , or  $^{32}\text{S}$ .

Sulfur cannot fuse with itself, the potential barrier is too high. Through successive  $\alpha$  captures, the star synthesizes elements in the “iron peak”: iron, nickel, cobalt.

The core tries to contract in the attempt to get them to burn, but they have the maximum possible binding energy per nucleon. So, the contraction continues.

If the mass of the contracting core exceeds the Chandrasekhar limit, it cannot become an electron-degenerate object. The mass of the core is always in excess of this limit mass for the stars which are massive enough to reach this stage of stellar burning.

Iron is photodissociated to make helium nuclei first, then bare protons, electrons and neutrons. Protons and electrons can combine into neutrons. The core becomes more and more neutrons rich, but the reaction also produces neutrinos, which can fly away.

The *neutron* degeneracy pressure can stop the collapse in certain cases: this is how a neutron star is formed. The threshold between neutron stars and black holes is hard to determine, but generally speaking with  $8M_{\odot} < M < 25M_{\odot}$  a neutron star is formed, while for larger masses the core collapses further to form a black hole.

The freefall velocity is a significant fraction of the speed of light. What are the statistics? how many NS and BH are there in our galaxy?

We can model the distribution of star masses in our galaxy with the distribution, the IMF, as a Salpeter IMF,<sup>1</sup> which is given by

$$N(m) \propto m^{-\alpha}, \quad (0.1.1)$$

where  $\alpha \approx 2.35$ . Then, we can calculate the number of stars which have more than  $8M_{\odot}$  by integrating: we find something proportional to  $8^{-1.35}$ , while the number of stars which have more than  $25M_{\odot}$  we get something proportional to  $25^{-1.35}$ . These will give us the amount of compact objects. The proportionality constant depend on the minimum and maximum mass of stars, but we can calculate the ratio of the two without concern for it. We find

$$\frac{N_{BH}}{N_{NS} + N_{BH}} = \left( \frac{8}{25} \right)^{-1.35} \approx 0.2. \quad (0.1.2)$$

The present rate of supernova explosions in the galaxy is around 1 per century, or  $10^{-2} \text{ yr}^{-1}$ . In the age of the galaxy (around  $10^{10} \text{ yr}$ ), we will then have had around  $10^8$  compact objects.

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<sup>1</sup> See the evil organization in Mission Impossible.

How much can we trust this figure? Kind of, the true number is closer to  $10^9$ , about 1 % of the number of stars in the galaxy.

The galaxy roughly looks like a cylinder with radius  $R \sim 60$  kpc and height  $H \sim 1$  kpc. Its volume will then be  $V = 2\pi R^2 H \approx 10^{13} \text{ pc}^3$ . Then, the number density of compact objects is around  $0.1 \text{ pc}^{-3}$ .

The typical separation between them will be something like 20 pc. Compact objects are close, common! Beware!

The closest compact object we know of is a neutron star 60 pc away: this is on the same order of magnitude.

**Compactness** The gravitational radius characterizing an object is

$$R_g = \frac{GM}{c^2}, \quad (0.1.3)$$

while the Schwarzschild radius is  $2R_g$ . For the Sun, this is approximately 1.5 km. It's small.

The value  $R_g/R$  is 0.5 for black holes, 0.15 for neutron stars,  $10^{-4}$  for white dwarfs.

As we said earlier, massive stars go type-2 supernova: this corresponds to Core-Collapse. Compact objects are quite common in the galaxy.

A compact object is one for which the ratio of the gravitational radius  $R_g = GM/c^2$  is comparable to the radius of the true object. For a white dwarf, the ratio is of the order of  $10^3$ .

We then need GR in order to deal with them. Let us quickly go over exact solutions of the Einstein Field Equations.

This lecture, we consider the vacuum Schwarzschild solution. The most general line element which is spherically symmetric (invariant under spatial rotations) must be made up of elements which are themselves invariant under spatial rotations. We will use spherical coordinates:  $r, \theta, \varphi, t$ .

In flat spacetime, the line element reads

$$ds^2 = -c^2 dt^2 + dr^2 + r^2 (d\theta^2 + \sin^2 \theta d\varphi^2), \quad (0.1.4)$$

and our Schwarzschild solution will need to reduce to this in some limit.

The spatial line element is given by

$$d\vec{r} \cdot d\vec{r} = dr^2 + r^2 (d\theta^2 + \sin^2 \theta d\varphi^2) = g_{ij} dx^i dx^j. \quad (0.1.5)$$

Then, the most general spherically symmetric line element will read

$$ds^2 = F(r, t) dt^2 + M(r, t) dr^2 + G(r, t) dr dt + C(r, t) r^2 (d\theta^2 + \sin^2 \theta d\varphi^2), \quad (0.1.6)$$

however, we can redefine the radial coordinate in order to remove the function multiplying the angular term, so we get

$$ds^2 = F dt^2 + M dr^2 + G dr dt + r^2 (d\theta^2 + \sin^2 \theta d\varphi^2). \quad (0.1.7)$$

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We can also introduce a new time variable:

$$dt' = dt + \psi(r, t) dr. \quad (0.1.8)$$

If  $\psi = rG/F$ , then we remove the mixed term, and then we are left with the expression

$$ds^2 = -B(r, t) dt'^2 + A dr^2 + r^2 d\Omega^2. \quad (0.1.9)$$

However, we have not yet determined the two functions, and we have not said anything about the Einstein Field Equations, which are

$$R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = 8\pi GT_{\mu\nu}. \quad (0.1.10)$$

In vacuo, the stress-energy tensor vanishes. The curvature scalar must vanish (we can show this by contracting the EFE with the inverse metric), so the equations reduce to  $R_{\mu\nu} = 0$ . We restrict ourselves to the static case.

The linearly independent components of the Ricci tensor read

$$R_0^0 = \frac{B''}{2AB} - \frac{A'B'}{4A''B} - \frac{B'^2}{4AB^2} + \frac{B'}{rAB} = 0 \quad (0.1.11)$$

$$R_1^1 = \frac{B''}{2AB} - \frac{A'B'}{4A''B} - \frac{B'^2}{4AB^2} + \frac{A'}{rA^2} = 0 \quad (0.1.12)$$

$$R_2^2 = \frac{1}{4rA} \left( \frac{B'}{B} - \frac{A'}{A} \right) + \frac{1}{r^2} \left( \frac{1}{A} - 1 \right). \quad (0.1.13)$$

Computing  $R_0^0 - R_1^1 = 0$  we find

$$\frac{1}{rA} \left( \frac{B'}{B} + \frac{A'}{A} \right) = 0 \quad (0.1.14)$$

$$\frac{d \log(AB)}{dr} = 0, \quad (0.1.15)$$

so  $AB$  is constant. Without losing generality we can take  $A = 1/B$ , since if this is not the case we can just rescale the radial or temporal coordinate until it is.

Then, we can compute

$$R_2^2 = \frac{B}{2r} \left( \frac{B'}{B} + \frac{B'}{B} \right) + \frac{1}{r^2} (B - 1) = 0 \quad (0.1.16)$$

$$B' + \frac{B}{r} - \frac{1}{r} = 0 \quad (0.1.17)$$

$$\frac{d}{dr}(rB) = 1, \quad (0.1.18)$$

so  $rB(r) = r + C$ , or equivalently

$$B(r) = \frac{C}{r} + 1. \quad (0.1.19)$$

After this, we can already substitute into the metric:

$$ds^2 = -\left(1 + \frac{C}{r}\right) dt^2 + \frac{1}{1 + C/r} dr^2 + r^2 d\Omega^2. \quad (0.1.20)$$

For any value of  $C$ ,  $B \rightarrow 1$  as  $r \rightarrow \infty$ : the metric reduces to the flat one asymptotically. Right now  $C$  is an arbitrary constant, however in the weak field limit it is known that

$$g_{00} = -\left(1 + 2\frac{\phi}{c^2}\right), \quad (0.1.21)$$

where  $\phi = -GM/r$  is the Newtonian gravitational field. Equating this expression to the one for  $g_{00}$ , we find

$$g_{00} = -\left(1 - \frac{2GM}{rc^2}\right) \implies C = -\frac{2GM}{c^2}. \quad (0.1.22)$$

The constant  $M$  in the classical case is the mass of the source, however we are computing a vacuum solution. This is the mass we would compute if we were to measure the orbits of objects around the compact object.

This is then surely a *gravitational* mass, is it also an *inertial* mass? Can we show this in GR?

Then, we can write the Schwarzschild metric:

$$ds^2 = -\left(1 - \frac{2GM}{c^2 r}\right) dt^2 + \left(1 - \frac{2GM}{c^2 r}\right)^{-1} dr^2 + r^2 d\Omega^2. \quad (0.1.23)$$

This is derived by assuming time-independence, however the result is the same even in the time-dependent case by the Jebsen-Birkhoff theorem (which we will prove in a moment). The element  $g_{rr}$  diverges as  $r \rightarrow R_g = 2GM/c^2$ , however this does not represent any physical divergence: the scalar  $R^{\mu\nu\rho\sigma}R_{\mu\nu\rho\sigma} \propto r^{-6}$  does not diverge there.

There are coordinates which do not diverge near the horizon: one classical choice employs the “tortoise” coordinates, which are the same for  $r$ ,  $\theta$ ,  $\phi$  as the Schwarzschild ones, while the time becomes (setting  $G = c = 1$ )

$$t = t' - 2M \log \left(1 - \frac{r'}{2M}\right). \quad (0.1.24)$$

Substituting this into the metric yields (dropping the primes for clarity):

$$ds^2 = -\left(1 - \frac{2M}{r}\right) dt^2 + \frac{4M}{r} dr dt + \left(1 + \frac{2M}{r}\right) dr^2 + r^2 d\Omega^2. \quad (0.1.25)$$

There is no pathology at  $r = 2M$  anymore, so it was not a physical divergence. The temporal coefficient  $g_{00}$  is the same: it can be shown that it is an invariant under coordinate transformations.

If we take two points which are very close along a particle trajectory, they must be separated by an interval  $ds^2 < 0$ .

If we consider a (nongeodesic!) radial path described by  $r(t)$ , we can compute the corresponding line element by neglecting the angular part:

$$ds^2 = -\left(1 - \frac{2M}{r}\right) dt^2 + \frac{4M}{r} dr dt + \left(1 + \frac{2M}{r}\right) dr^2 \quad (0.1.26)$$

$$\left(\frac{ds}{dt}\right)^2 = -\left(1 - \frac{2M}{r}\right) + \frac{4M}{r} \frac{dr}{dt} + \left(1 + \frac{2M}{r}\right) \left(\frac{dr}{dt}\right)^2. \quad (0.1.27)$$

Now, the question we ask is: is it possible for the particle trajectory to be timelike or lightlike ( $ds^2 \leq 0$ ) and outgoing ( $dr/dt > 0$ ) under these conditions? If this is the case, the signs of the three terms read

$$\underbrace{\left(\frac{ds}{dt}\right)^2}_{<0?} = -\left(1 - \frac{2M}{r}\right) + \underbrace{\frac{4M}{r} \frac{dr}{dt}}_{>0} + \underbrace{\left(1 + \frac{2M}{r}\right) \left(\frac{dr}{dt}\right)^2}_{>0}, \quad (0.1.28)$$

so we can see that the equality can be satisfied (a positive number cannot equal a negative one!) as long as the first term on the right-hand side is negative, which means  $r > 2M$ . If  $r \leq 2M$ , on the other hand, this cannot be the case: a radial trajectory below the horizon *cannot* be outward.

This is what “horizon” means: it is a *semi-permeable* membrane, particles can surpass it only in one direction.

**Jebsen-Birkhoff** This theorem states that the Schwarzschild solution also describes the spacetime around an object in the spherically-symmetric but *time-dependent* case. Let us give a sketch of its proof, omitting some tedious calculations. If we write out the components of the Ricci tensor, we find something in the form

$$R_0^0 = R_0^0 \Big|_{\text{static}} + \dot{A}(\dots) \quad (0.1.29)$$

$$R_1^1 = R_1^1 \Big|_{\text{static}} + \dot{A}(\dots), \quad (0.1.30)$$

while  $R_2^2$  and  $R_3^3$  are the same. Also, the term  $R_0^1$  does not vanish unlike the static case, and is equal to

$$R_0^1 = -\frac{\dot{A}}{rA^2} = 0, \quad (0.1.31)$$

so  $\dot{A} = 0$ : the equations then are the same as the static case ones! This, however, is not the end, since now the equation

$$\frac{A'}{A} + \frac{B'}{B} = 0 \quad (0.1.32)$$

is not necessarily solved by  $\log A = -\log B$ , since a prime denotes a *partial* derivative with respect to  $r$ , so in general we will have  $\log A + \log B = f(t)$ , some generic function of time. The metric will then read

$$ds^2 = -\left(1 - \frac{2M}{r}\right) f(t) dt^2 + \left(1 - \frac{2M}{r}\right) dr^2 + r^2 d\Omega^2, \quad (0.1.33)$$

but we can simply rescale the time coordinate to  $t \rightarrow \sqrt{f}t$  in order for this to reduce to the usual expression. This theorem was originally discovered by the Norwegian physicist Jebsen, and only later popularized in a textbook by Birkoff [JR05].

The source of the geometry can change in time while leaving the outside spacetime unperturbed, however this holds only as long as the variation remains spherically symmetric: collapse, expansion or pulsation. Any asymmetry can lead to the emission of gravitational radiation.

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2020-10-7,  
compiled  
2020-10-13

“Schwarzschild” means “Black Shield”: no relation though, it is the name of a German scientist. We have discussed the properties of this metric.

Now, let us consider **geodesic motion** in the vacuum Schwarzschild spacetime.

We cannot choose an arbitrary radius for an orbit around a Black Hole: there is an Innermost Stable Circular Orbit, and this has a direct impact on the accretion efficiency.

We will use Latin letters for spacetime indices. The four-velocity, as usual, is  $u^i = dx^i/d\tau$ , where  $d\tau$  is the proper time. Also, for massive particles we also have the four-momentum  $p^i = mu^i$ .

Analogously to classical mechanics, the Lagrangian of a free particle is

$$L = -\sqrt{-g_{ij}p^ip^j} = -m\sqrt{-g_{ij}u^iu^j}, \quad (0.134)$$

and the corresponding Euler-Lagrange equations read

$$\frac{d}{d\tau} \left( \frac{\partial L}{\partial u^i} \right) - \frac{\partial L}{\partial x^i} = 0, \quad (0.135)$$

however we will not study these in the general case, it is too complicated. It can be shown that along a geodesic the Lagrangian itself is conserved:  $L = \text{const}$ . Therefore,  $g_{ij}u^iu^j$  is also a constant, a negative one since the velocity is timelike. We can then set it to  $-1$  by reparametrizing, so we will have  $u^2 = -1$  and  $p^2 = -m^2$ .

Suppose that there is a certain coordinate  $x^k$  such that

$$\frac{\partial g_{ij}}{\partial x^k} = 0, \quad (0.136)$$

that is, the metric does not change along the  $x^k$  direction. Therefore,

$$\frac{\partial L}{\partial x^k} = 0, \quad (0.137)$$

which tells us that

$$\frac{d}{d\tau} \left( \frac{\partial L}{\partial u^k} \right) = 0, \quad (0.138)$$

meaning that this cyclic variable gives us a conserved quantity. If we do the calculation, the constant comes out to be

$$mg_{kj}u^j = p_k = \text{const}. \quad (0.139)$$



This is the case in Schwarzschild spacetime: the metric only depends on  $r$  and  $\theta$ , while  $t$  and  $\varphi$  are cyclic. Therefore, we will have two constants of motion: in terms of the momentum,  $p_t$  and  $p_\varphi$ .

We can denote them as  $E = -p_t$  and  $L = p_\varphi$ . Let us start from  $L$ :

$$L = g_{\varphi k} p^k = g_{\varphi\varphi} m \frac{dx^\varphi}{d\tau} \quad (0.1.40)$$

$$= r^2 \sin^2 \theta m \frac{d\varphi}{d\tau} \quad (0.1.41)$$

$$\frac{d\varphi}{d\tau} = \frac{L}{mr^2 \sin^2 \theta} . \quad (0.1.42)$$

Now, geodesic motion is in general planar since the metric is spherically symmetric: we do not lose any generality by setting  $\theta = \pi/2$ , so we can simplify  $\sin^2 \theta = 1$ .

In general, even in Newtonian mechanics, when discussing orbits we can decompose the velocity into the radial and perpendicular direction. The angular momentum, in classical Newtonian mechanics, has a modulus  $|\vec{L}| = |\vec{r} \wedge m\vec{v}| = mrv_\varphi = mr^2\dot{\varphi}$ . This is the same relation we have found here.

The energy of an object with four-momentum  $p_k$  as measured by an observer with four-velocity  $u^k$  is given by  $E = -p_k u^k$ . We choose an observer which is static with respect to the coordinates:<sup>2</sup> then, we have  $dt/d\tau = \gamma/\sqrt{-g_{tt}}$ , so

$$E = -g_{tt} m \frac{dx^t}{d\tau} = + \left(1 - \frac{2M}{r}\right) \frac{m}{\sqrt{1-v^2}} \frac{1}{\sqrt{-g_{tt}}} \quad (0.1.43)$$

$$= \sqrt{1 - \frac{2M}{r}} m \gamma \quad (0.1.44)$$

$$\approx m \left(1 - \frac{M}{r} + \frac{v^2}{2}\right) \quad (0.1.45)$$

so in the Newtonian limit ( $M \ll r$ ,  $v \ll 1$ ) this is regular expression for the conservation of energy, gravitational plus kinetic.

Now we have the tools to study geodesic motion: let us write down  $p^2 = -m^2$  explicitly,

$$g_{tt} (p^t)^2 + g_{rr} (p^r)^2 + \underbrace{g_{\theta\theta} (p^\theta)^2}_{=0 \text{ if } \theta \equiv \pi/2} + g_{\varphi\varphi} (p^\varphi)^2 = -m^2, \quad (0.1.46)$$

and if we substitute we must be careful: the constants of motion are related to the *covariant* components of the quantities, so we have

$$g_{tt} (g^{tt} p_t)^2 + g_{rr} (g^{rr} p_r)^2 + g_{\varphi\varphi} (g^{\varphi\varphi} p_\varphi)^2 = -m^2 \quad (0.1.47)$$

$$g_{tt} (g^{tt})^2 E^2 + g_{rr} (p^r)^2 + g_{\varphi\varphi} (g^{\varphi\varphi})^2 L^2 = -m^2 \quad (0.1.48)$$

$$g^{tt} E^2 + g^{rr} (p_r)^2 + g^{\varphi\varphi} L^2 = -m^2 \quad (0.1.49)$$

<sup>2</sup> Its four velocity will be given by  $k^i = (1/\sqrt{-g_{tt}}, \vec{0})$  by normalization.

$$-\left(1 - \frac{2M}{r}\right)^{-1} E^2 + \left(1 - \frac{2M}{r}\right)^{-1} (p^r)^2 + \frac{L^2}{r^2} = -m^2 \quad (0.1.50)$$

$$-\left(1 - \frac{2M}{r}\right)^{-1} E^2 + \left(1 - \frac{2M}{r}\right)^{-1} m^2 \left(\frac{dr}{d\tau}\right)^2 + \frac{L^2}{r^2} = -m^2, \quad (0.1.51)$$

which is an ODE for the single function  $r(\tau)$ : we can integrate this to find the shape of the trajectory. We now divide everything by  $m^2$ , and define the specific energy and angular momentum  $\epsilon = E/m$  and  $\ell = L/m$ :

$$\left(1 - \frac{2M}{r}\right)^{-1} \left[ \left(\frac{dr}{d\tau}\right)^2 - \epsilon^2 \right] + \frac{\ell^2}{r^2} = -1. \quad (0.1.52)$$

We can then express this as

$$\left(\frac{dr}{d\tau}\right)^2 = -\left(1 + \frac{\ell^2}{r^2}\right) \left(1 - \frac{2M}{r}\right) + \epsilon^2, \quad (0.1.53)$$

which can give us  $r(\tau)$  or  $r(\varphi)$  in a rather simple way numerically, however the answer is not analytic: it is given by an elliptic integral.

We can study it analytically by looking at the turning points, those with  $\frac{dr}{d\tau} = 0$ . These will tell us about the shape of the orbit. They obey the equation

$$0 = -\left(1 + \frac{\ell^2}{r^2}\right) \left(1 - \frac{2M}{r}\right) + \epsilon^2, \quad (0.1.54)$$

which means

$$\frac{\ell^2}{r^2} = \left(\epsilon^2 - 1 + \frac{2M}{r}\right) \left(1 - \frac{2M}{r}\right)^{-1}, \quad (0.1.55)$$

which has the solutions

$$l_{\pm} = \pm r \left[ \left(\epsilon^2 - 1 + \frac{2M}{r}\right) \left(1 - \frac{2M}{r}\right) \right]^{1/2}. \quad (0.1.56)$$

Let us define  $\lambda = \ell/2M$  and  $x = r/2M$ . These are dimensionless in our units. Also, we define  $\Gamma = \epsilon^2 - 1$ . They satisfy

$$\lambda_{\pm} = \pm x \sqrt{\left(\Gamma + \frac{1}{x}\right) \left(1 - \frac{1}{x}\right)^{-1}} = \pm x \sqrt{\frac{x\Gamma + 1}{x - 1}}. \quad (0.1.57)$$

We want to draw these as curves in the  $x, \lambda$  plane. We fix a value of  $\ell$ , which means we have chosen  $\lambda$ . This depends on the initial conditions of the motion of the particle. If this  $\lambda$  is a  $\lambda_+$ , then  $x$  can move in all the region  $x(\lambda_+) < x < \infty$ .

This is the relativistic analog of the hyperbolic trajectories in the Keplerian case.

Let us try to compute  $\frac{d\lambda_{\pm}}{dx}$ : this yields

$$2x^2\Gamma + x - 3x\Gamma - 2 = 0 \quad (0.1.58)$$

$$\Gamma(2x^2 - 3x) + x - 2 = 0 \quad (0.1.59)$$

$$\Gamma = \frac{2 - x}{2x^2 - 3x}. \quad (0.1.60)$$

This means we have a single curve, which will relate  $\Gamma$  and  $x$ . The physical cases are only the ones with  $\Gamma > -1$ . The minimum is at  $x = 3$ , corresponding to  $r = 6M$ , where we have  $\Gamma = -1/9$ . This corresponds to  $\epsilon^2 - 1$ : so,  $\epsilon = \sqrt{2} \times 2/3$ .

For  $-1/9 < \Gamma < 0$  we have two choices for  $x$ , for  $\Gamma > 0$  we have only one.

Coming back to the  $\lambda_{\pm}$  curves, we can have  $\lambda_{\pm}$  as two distinct curves, and there is a region around  $\lambda = 0$  for which the particle will cross the horizon.

This allows us to compute the cross-section for gravitational capture, which is nonvanishing *even if* we neglect the size of the black hole.

Insert plots

We were discussing geodesics in the external Schwarzschild solution. We wrote down the locus of the inversion point  $\dot{r} = 0$ : this is

$$\lambda_{\pm} = \pm x \sqrt{\frac{\Gamma x + 1}{x - 1}}. \quad (0.1.61)$$

1. For  $\Gamma > 0$  there is one extremum;
2. for  $-1/9 < \Gamma < 0$  there are two extrema;
3. for  $-1 < \Gamma < -1/9$  there are no extrema.

Let us consider the two-extrema case. Then, there is an intersection with the  $x$  axis at  $x = -1/\Gamma$  for the  $\lambda_{\pm}$ . Then the curve looks like a doorknob: it is a closed curve.

In this case, we are considering *bound states*: the particle can start on the LHS of the curve, and if this is the case it must reach the inversion point and then go back and eventually cross the horizon. If it is in the “doorknob region”, then it is trapped there.

These orbits are similar to the Newtonian elliptic ones, but there are differences. These are radially bound, however they are not in general periodic: they will precess.

We can also have circular orbits, both stable and unstable.

In the no-extrema case there are no orbits. The limiting case  $\Gamma = -1/9$  is interesting: the minimum becomes an inflection point, with zero second derivative. This happens at  $x = 3$ . This is the ISCO: the Innermost Stable Circular Orbit.

The corresponding energy is given by  $\epsilon^2 - 1 = -1/9$ :  $\epsilon = 2\sqrt{2}/3$ , so

$$E_{\text{ISCO}} = \frac{2\sqrt{2}}{3} mc^2. \quad (0.1.62)$$

If we can bring a particle from infinity to the ISCO (which is what happens as an accretion disk is formed) it can release an energy equal to  $mc^2 - E_{\text{ISCO}} = (1 - 2\sqrt{2}/3)mc^2$ . This is why an accretion disk can produce a large amount of energy.

Tuesday  
2020-10-13,  
compiled  
2020-10-13

The efficiency can then be computed as

$$\text{efficiency} = \frac{mc^2 - 2\sqrt{2}mc^2/3}{mc^2} \approx 6\% . \quad (0.1.63)$$

This is a *very large* efficiency: an order of magnitude more than the efficiency of nuclear burning.

## 0.2 Schwarzschild internal solution

Now we try to solve the EFE under spherical symmetry *with* a source:

$$G_{ij} = R_{ij} - \frac{1}{2}g_{ij}R = 8\pi T_{ij} , \quad (0.2.1)$$

with a metric

$$ds^2 = -B(r) dt^2 + A(r) dr^2 + r^2 d\Omega^2 . \quad (0.2.2)$$

Birkhoff does not apply here: we are *assuming* stationarity. We will use a very simple  $T_{ij}$ : a perfect fluid, where

$$T_{ij} = (\rho + P)u_i u_j + P g_{ij} , \quad (0.2.3)$$

where the energy density is given by  $\rho = \rho_0(1 + \epsilon)$ :  $\rho_0$  is the rest energy density of the fluid, while  $\epsilon$  accounts for thermal motion and other kinds of internal energy.

We will assume that the fluid is at rest with respect to the interior: we want to find an analog of the hydrostatic equilibrium equation. So,  $dr = d\theta = d\varphi = 0$ , while only  $dt \neq 0$ . This means that only  $u^0 \neq 0$ , and normalization requires  $g_{ij}u^i u^j = -1$ : this means that  $u^0 = 1/\sqrt{B}$ , and the whole vector reads  $u^i = \delta_0^i / \sqrt{-g_{00}}$ .

Let us then write the nonzero mixed components of the stress-energy tensor:

$$T_0^0 = -\rho \quad T_1^1 = T_2^2 = T_3^3 = P , \quad (0.2.4)$$

since if we have one upper and one lower index the normalization in  $u^i u_j$  simplifies.

Let us also skip the computations: the Einstein tensor reads

$$G_0^0 = \frac{1}{A} \left( \frac{1}{r^2} - \frac{A'}{rA} \right) - \frac{1}{r^2} = -8\pi\rho \quad (0.2.5)$$

$$G_1^1 = \frac{1}{A} \left( \frac{1}{r^2} + \frac{B'}{rB} \right) - \frac{1}{r^2} = 8\pi P \quad (0.2.6)$$

$$G_2^2 = G_3^3 = -\frac{1}{2A} \left[ \frac{A'B'}{2AB} + \left( \frac{B'}{B} \right)^2 - \frac{B'}{Br^2} + \frac{A'}{r^2 A} - \frac{B''}{2B} \right] = 8\pi P . \quad (0.2.7)$$

We need to solve for  $\rho$ ,  $P$ ,  $A$  and  $B$ : the equations of motion of the particles,  $\nabla_i T^{ij} = 0$ , are a consequence of the Einstein equations.

We have three equations for four variables: we need an additional one, typically we combine them with an *equation of state* for  $P$ : a simple case, a *barotropic EOS*, is in the form  $P = P(\rho)$ .

The first equation can be written as

$$\frac{1}{A} - \frac{rA'}{A^2} = 1 - 8\pi r^2 \rho, \quad (0.2.8)$$

which we can integrate to find

$$\frac{r}{A} = r - \int_0^r 8\pi \tilde{r}^2 \rho \, d\tilde{r}. \quad (0.2.9)$$

We introduce the quantity

$$m(r) = \int_0^r 4\pi \tilde{r}^2 \rho \, d\tilde{r}, \quad (0.2.10)$$

so that

$$A(r) = \left(1 - \frac{2m(r)}{r}\right)^{-1}. \quad (0.2.11)$$

The problem is that  $m(r)$  is *not* the total mass-energy enclosed in the region below  $r$ : we are not integrating with respect to a covariant volume form.

With the second two equations, we end up with

$$\frac{1}{2B} \frac{dB}{dr} = -\frac{1}{P + \rho} \frac{dP}{dr}. \quad (0.2.12)$$

We have essentially solved for  $B$ . We can also manipulate the equations to find:

$$\frac{dP}{dr} = -\frac{m(r)\rho}{r^2} \left(1 + \frac{P}{\rho}\right) \left(1 + \frac{4\pi r^3 P}{m(r)}\right) \left(1 - \frac{2m(r)}{r}\right)^{-1}, \quad (0.2.13)$$

the **Tolman-Oppenheimer-Volkov** (TOV) equation. This reduces to the hydrostatic equilibrium equation in weak gravity.

The three other terms are basically three relativistic corrections, from velocity and radius: the first are applied based on whether the gas is relativistic *in its own rest frame*; if this is the case, then  $P \sim \rho$ . The first term is a local correction, the second one is a global one; the third term on the other hand is the usual GR correction, based on  $r_{\text{Schw}}/r$ .

We can also write

$$\frac{dm}{dr} = 4\pi r^2 \rho. \quad (0.2.14)$$

This equation, the TOV one and the equation of state form a closed system for  $m$ ,  $\rho$  and  $P$ .

As for boundary conditions, we require that  $m(0) = 0$  and  $P(R) = 0$ . This is hard to do numerically: it is a boundary value problem, not an initial condition one.

Finally, we discuss the meaning of  $m(r)$ . The total  $M$  is given by

$$M = \int_0^R 4\pi r^2 \rho \, dr , \quad (0.2.15)$$

where  $R$  is the radius of the star. This is *not* the total mass, since  $dV$  is *not*  $4\pi r^2 dr$ . The *proper* three-volume element is instead given by  $dV = 4\pi r^2 \sqrt{A(r)} \, dr$ , since the true radial distance is calculated through the metric:  $ds^2 = A(r) \, dr^2$  if the measured length is radial.

We call  $M_*$  the mass which is calculated including the  $\sqrt{A}$  term. Their difference is

$$M - M_* = \int_0^R 4\pi r^2 \rho \left(1 - \sqrt{A}\right) \, dr . \quad (0.2.16)$$

If  $r \gg 2m$  (the weak-field limit) then we can approximate  $A = 1 - 2m/r$ :

$$M - M_* \approx \int_0^R 4\pi r^2 \rho \left(1 - 1 - \frac{2m}{r}\right) \, dr \quad (0.2.17)$$

$$\approx - \int_0^R \underbrace{4\pi r^2 \rho \, dr}_{dm} \frac{m}{r} \quad (0.2.18)$$

$$\approx - \int_M \frac{Gm}{r} \, dm = E_g . \quad (0.2.19)$$

So, in the weak field limit, the mass defect is given by the gravitational potential energy! This is a physical consequence of the *nonlinearity* of the EFE.

The electromagnetic field behaves linearly at low energies, while in QED we can see nonlinearities.

# Bibliography

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