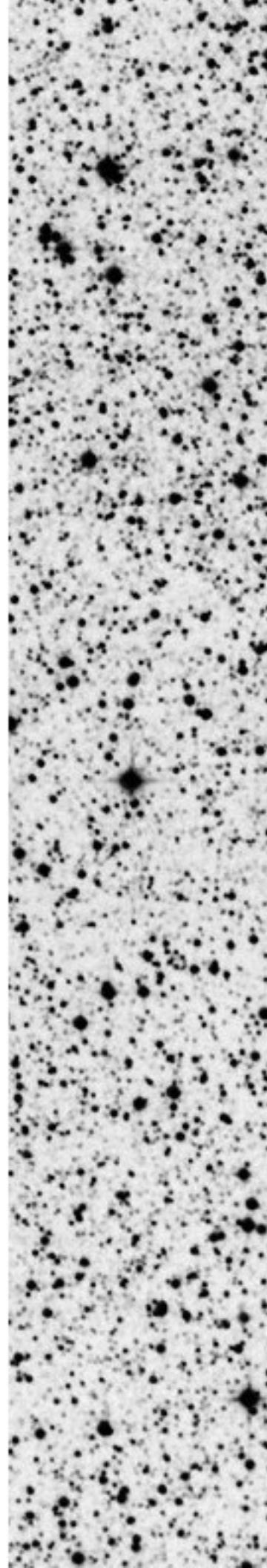


EDWARD BROWN

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# STELLAR ASTROPHYSICS

OPEN ASTROPHYSICS BOOKSHELF



*About the cover:* The image shows the star field in the direction of the second closest star to Earth, Proxima Centauri.

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## Preface

These notes were written while teaching a graduate-level astronomy course on stars at Michigan State University. The only background preparation for this course is undergraduate physics and a course on radiative processes, and so portions of these notes could be useful for upper-level undergraduates. The text layout uses the `tufte-book` (<https://tufte-latex.github.io/tufte-latex/>) L<sup>A</sup>T<sub>E</sub>X class: the main feature is a large right margin in which the students can take notes; this margin also holds small figures and sidenotes. Exercises are embedded throughout the text. These range from “reading exercises” to longer, more challenging problems. In addition, there are several numerical exercises that use the MESA stellar evolution code, available from <http://mesa.sourceforge.net/>. These numerical exercises are prefaced with the logo **MESA**, used by kind permission of the MESA council. Because the exercises are spread throughout the text, there is a “List of Exercises” in the front matter to help with looking for specific problems.

The course notes were originally meant as a supplement to the main text, Hansen et al.<sup>1</sup>; in some editions of the course I also drew from Clayton<sup>2</sup> and Kippenhahn and Weigert<sup>3</sup>. These notes therefore tend to expand upon topics not already covered there. In the second half of the course, the students typically gave presentations on current topics in stellar evolution, and I supplemented those with readings from the MESA instrument papers<sup>4</sup>. As a result, however, my notes on topics of stellar evolution have lagged behind the rest of the text and are not yet ready for posting.

Some of the material was inspired by three courses at UC-Berkeley in the mid-90’s: “Stars with Lars”, taught by Professor L. Bildsten; Statistical Physics, taught by Professor E. Commins, and Fluid Mechanics, taught by Professor J. Graham. I am also indebted to the students who took the MSU stellar physics course for their questions, feedback, and encouragement. Additional thanks go to MSU graduate students Dana Koeppe and Wei Jia Ong for reading late stages of the drafts and testing the numerical exercises.

PLEASE BE ADVISED THAT THESE NOTES ARE UNDER ACTIVE DEVELOPMENT; to refer to a specific version, use the eight-character stamp labeled “git version” on the copyright page.

<sup>1</sup> Carl J. Hansen, Steven D. Kawaler, and Virginia Trimble. *Stellar Interiors*. Springer-Verlag, 2d edition, 2004

<sup>2</sup> Donald D. Clayton. *Principles of Stellar Evolution and Nucleosynthesis*. University of Chicago Press, 1983

<sup>3</sup> R. Kippenhahn and A. Weigert. *Stellar Structure and Evolution*. Springer-Verlag, 1994

<sup>4</sup> Bill Paxton, Lars Bildsten, Aaron Dotter, Falk Herwig, Pierre Lesaffre, and Frank Timmes. Modules for experiments in stellar astrophysics (MESA). *ApJS*, 192:3, January 2011; and Bill Paxton, Matteo Cantiello, Phil Arras, Lars Bildsten, Edward F. Brown, Aaron Dotter, Christopher Mankovich, M. H. Montgomery, Dennis Stello, F. X. Timmes, and Richard Townsend. Modules for experiments in stellar astrophysics (MESA): Planets, oscillations, rotation, and massive stars. *ApJS*, 208:4, 2013



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# 1

## Preliminaries and Stellar Zoology

Much of these notes will be concerned with the essential *physics* of stars. We will focus on the questions: What determines the structure and evolution of stars of varied initial properties? What are the end states of stars? and How do stars affect cosmic chemical evolution?

We shall see that much of the salient physics of stars separates into two broad categories: *macrophysics* and *micriphysics*, the former characterized by, e.g., the mechanical equations of stellar structure, and the latter characterized by, e.g., the physics of nuclear reactions. We will also see that these two limits, the large and small, are fundamentally coupled and that, in large part, the macrophysical equations are simply a continuum limit of the governing microphysical equations (analogous to how classical physics is a particular limit of quantum physics). In an extremely reductionist sense, a star's life can be described as a struggle between the inexorable inward pull of its own gravity and the competing outward push of various forms of pressure, the latter generated by different types of microphysical processes. As a star evolves and changes, the physics that dominates the generation of this outward pressure changes frequently, altering the essential character of the star in important ways. But gravity remains every constant in its pull, seeking to make anything and everything it can into a black hole.<sup>1</sup>

Before digging into the physics, it is useful to review some observational and nomenclatural aspects of stars. Figure 1.1 shows a field of stars in Sagittarius as imaged by the Hubble Space Telescope (HST). These stars clearly span a range in color, from red to blue, and brightness. Also apparent is that the bluer stars appear brighter. Stars are classified according to their *spectral type*, denoted by letters OBAFGKM. We now know that this seemingly random arrangement of letters reflects the effect temperatures at which stars radiate. The original spectral classification scheme developed at Harvard Observatory by Annie Jump Cannon was based solely on spectral features (Balmer lines) and was, indeed, alphabetical. Once it was under-

<sup>1</sup> anthropomorphism: *noun*, the attribution of human characteristics or behavior to a animal, object, or force of nature.

stood that different Balmer lines appear as a function of effective temperature, the spectral classes were rearranged in order of decrease temperature.

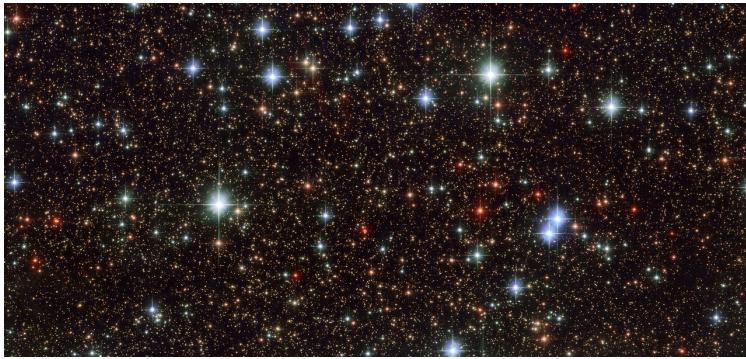


Figure 1.1: Hubble Space Telescope image of stars in the constellation Sagittarius. *Image credit: NASA/ESA.*

Table 1.1 lists the modern Harvard spectral types and some of their salient properties, including Main Sequence mass, radius, and luminosity. Immediately apparent is the strong dependence of all these properties on the mass of the stars. The most important parameter in determining the evolution of a star is its initial mass, usually termed as the Zero-age Main Sequence (ZAMS) mass. But this is not the only important evolutionary parameter. A star's *metallicity*, or content of elements heavier than H and He, as well as the initial angular momentum content and magnetic field strength are also critical. These latter two are inextricably linked and subject to substantial uncertainties, both in their typical values for stars and in their impact on stellar evolution. The roles of rotation and magnetic fields in stellar structure and evolution will be explored in Chapter 19. Another critical component in determining the fate of a star, particularly a massive star, is whether or not it has a close binary companion.

Important aspects of binarity will be discussed in Chapter 18.

Type	Effective temperature	Main sequence mass	Main sequence radius	Main sequence bolometric luminosity
O	$\gtrsim 30,000$ K	$\gtrsim 16 M_{\odot}$	$\gtrsim 6.6 R_{\odot}$	$\gtrsim 30,000 L_{\odot}$
B	10,000–30,000 K	2.1–16 $M_{\odot}$	1.8–6.6 $R_{\odot}$	25–30,000 $L_{\odot}$
A	7,500–10,000 K	1.4–2.1 $M_{\odot}$	1.4–1.8 $R_{\odot}$	5–25 $L_{\odot}$
F	6,000–7,500 K	1.04–1.4 $M_{\odot}$	1.15–1.4 $R_{\odot}$	1.5–5 $L_{\odot}$
G	5,200–6,000 K	0.8–1.04 $M_{\odot}$	0.96–1.15 $R_{\odot}$	0.6–1.5 $L_{\odot}$
K	3700–5200 K	0.45–0.8 $M_{\odot}$	0.7–0.96 $R_{\odot}$	0.08–0.6 $L_{\odot}$
M	2400–3700 K	0.08–0.45 $M_{\odot}$	$\lesssim 0.7 R_{\odot}$	$\lesssim 0.08 L_{\odot}$

Table 1.1: The Modern Harvard Stellar Spectral Types

The spectral classifications in Table 1.1 are supplemented by *luminosity classes*, denoted by Roman numerals, with “I” being the brightest supergiants and “V” being main sequence stars. The spectral classes are also typically augmented by an additional Arabic numeral further specifying the star’s precise effective temperature. So, for example, our Sun is a “G2 V,” being a main sequence star with an effective temperature of 5770 K.

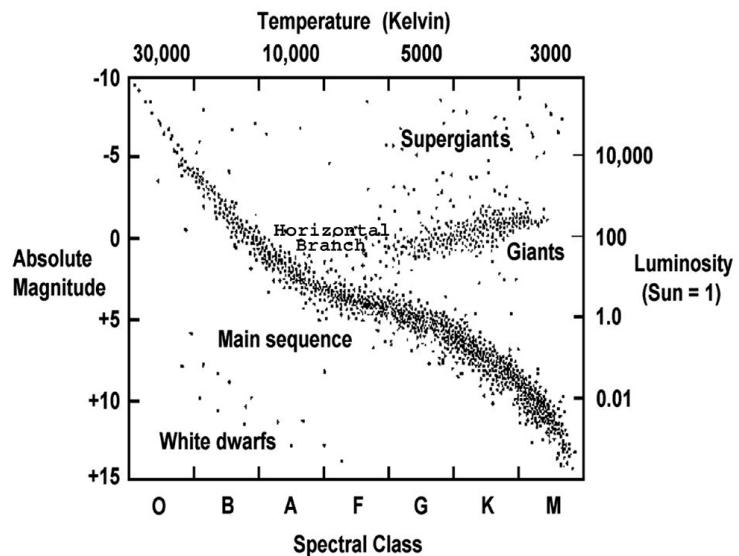


Figure 1.2: Example HR diagram for a subset of stars in the Milky Way.

The principle graphical tool for understanding stellar evolution is the *Hertzsprung-Russel (HR) diagram*. Stars are placed on the HR diagram depending on their effective temperature (i.e., spectral type) and their absolute luminosity. An example HR diagram is shown in Figure 1.2. Observationally, it is generally easier to produce a *color-magnitude* diagram in which the absolute magnitude in a given broadband filter is plotted against the *color*, usually the difference of magnitudes between two broadband filters.



## 2

### *The Sun on a Blackboard*

To begin our study of stellar structure, let us first consider the star that we know best, our sun. The planetary orbits and the gravitational constant  $G$  tell us its mass; our knowledge of the earth-sun distance and observations tell us its radius; measurements of the solar radiant flux and spectra tell us its luminosity and temperature; and radiometric dating of meteorites tells us the age of the solar system. In summary:

$$\begin{aligned}M_{\odot} &= 1.99 \times 10^{33} \text{ g} \\R_{\odot} &= 6.96 \times 10^{10} \text{ cm} \\L_{\odot} &= 3.86 \times 10^{33} \text{ erg s}^{-1} \\T_{\text{eff}} &= 5780 \text{ K} \\\tau_{\odot} &= 4.6 \text{ Gyr.}\end{aligned}$$

Moreover, the composition of the sun is well known<sup>1</sup>; the five most abundant elements are H, He(−1.07), N(−4.22), O(−3.34), and C(−3.61), where the number in parentheses is  $\log(n_{\text{el}}/n_{\text{H}})$ , the abundance relative to hydrogen.

Another salient feature of our sun is its stability: the power output is remarkably constant, varying by less than 0.1% over several solar cycles<sup>2</sup>, with inferred variations over 2,000 yr on a similar scale<sup>3</sup>. On longer timescales, evidence for liquid water over much of Earth history suggest that the power output of the sun cannot have varied greatly over its life. The first task, then, is to investigate the mechanical and thermal stability of a self-gravitating fluid.

---

**EXERCISE 2.1**— What is the mean density of the sun? What is the luminous flux (energy/area/time) at 1 AU? What is the orbital period of a test mass just exterior to the radius of the sun?

---

<sup>1</sup> E. Anders and N. Grevesse. Abundances of the elements - meteoritic and solar. *Geochim. Cosmochim. Acta*, 53:197–214, 1989; and M. Asplund, N. Grevesse, and A. J. Sauval. The Solar Chemical Composition. In T. G. Barnes, III and F. N. Bash, editors, *Cosmic Abundances as Records of Stellar Evolution and Nucleosynthesis*, volume 336 of *Astronomical Society of the Pacific Conference Series*, page 25, September 2005

<sup>2</sup> R. C. Willson and H. S. Hudson. The sun's luminosity over a complete solar cycle. *Nature*, 351:42–44, May 1991

<sup>3</sup> C. Fröhlich and J. Lean. Solar radiative output and its variability: evidence and mechanisms. *A&A Rev.*, 12:273–320, December 2004

## 2.1 Fluid equation of motion

We can describe a star by deriving differential equations for the density, pressure, temperature, luminosity and composition. Over scales that are large compared to the collisional mean free paths between particles, we can treat the fluid as a continuous medium. That is, we suppose that we can find a scale that is infinitesimal compared to the macroscopic scales, but still much larger than the scales for microscopic interactions. As a consequence, we can define thermodynamic quantities (which only make sense if we have a large sample of particles) at a location (which means our sample is small compared to the macroscopic lengthscales in the star).

Consider such a macroscopically small volume  $V$ . Its mass is  $M = \int_V \rho dV$ , where  $\rho$  is the mass density. If  $\mathbf{u}(x, t)$  is the velocity, then the flux of mass into the element is

$$-\int_{\partial V} \rho \mathbf{u} \cdot d\mathbf{S} = \frac{\partial}{\partial t} \int_V \rho dV$$

where the right-hand side follows from mass conservation. Using Gauss's law to transform the left-hand side into an integral over  $V$  and combining terms, we have

$$\int_V \left\{ \frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{u}) \right\} dV = 0.$$

Since this equation holds for any  $V$ , the integrand must vanish, and we have our first equation,

$$\partial_t \rho + \nabla \cdot (\rho \mathbf{u}) = 0. \quad (2.1)$$

Our next equation is to get the analog of  $\mathbf{F} = m\mathbf{a}$ . Ignoring viscous effects, the net force on our fluid element (with volume  $V$ ) is due to the pressure over its surface  $P$  and the gradient of the gravitational potential  $\Phi$ :

$$\int_V \rho \frac{d^2 \mathbf{r}}{dt^2} dV = \int_V \mathbf{F} dV = - \int_V \rho \nabla \Phi dV - \int_{\partial V} P d\mathbf{S}.$$

Transforming the second integral on the right-hand side to a volume integral, and assuming that  $\nabla \Phi$  and  $\nabla P$  vary on macroscopic lengthscales, we arrive at an equation for the acceleration,

$$\frac{d^2 \mathbf{r}}{dt^2} = -\nabla \Phi - \frac{1}{\rho} \nabla P. \quad (2.2)$$

where  $\mathbf{r}(t)$  is the position of the particle so that the left-hand side is the acceleration. Here we must be careful: the velocity of the fluid is specified by a field  $\mathbf{u}(x, t)$  that refers to the velocity of the fluid at a

given point in space and a given instance of time, *not* to the velocity of a given particle. A fluid element can still accelerate even if  $\partial_t \mathbf{u} = 0$  by virtue of moving to a different location. At time  $t$  this particle has the velocity

$$\frac{d\mathbf{r}}{dt} \Big|_t = \mathbf{u}(\mathbf{x} = \mathbf{r}|_{t,t}) \quad (2.3)$$

where we use the fact that the particle is moving along a streamline of the fluid. At a slightly later time  $t+h$ , the particle has moved to a location  $\mathbf{r}(t+h) \approx \mathbf{r}(t) + h\mathbf{u}$ , and its velocity is then

$$\frac{d\mathbf{r}}{dt} \Big|_{t+h} = \mathbf{u}(\mathbf{x} = \mathbf{r}|_{t+h}, t+h) \approx \mathbf{u} + h(\mathbf{u} \cdot \nabla \mathbf{u} + \partial_t \mathbf{u}), \quad (2.4)$$

where we evaluate the derivatives at time  $t$ . Subtracting equation (2.3) from equation (2.4) and dividing by  $h$  gives us the acceleration; inserting this into Newton's law and dividing by volume gives us Euler's equation of motion,

$$\partial_t \mathbf{u} + \mathbf{u} \cdot \nabla \mathbf{u} = -\nabla \Phi - \frac{1}{\rho} \nabla P. \quad (2.5)$$

Equations (2.1) and (2.5) form the first two equations we need to describe stellar structure.

**EXERCISE 2.2** — Using equation (2.1), show that equation (2.5) can be written as

$$\partial_t(\rho u_i) + \partial_j(\rho u_i u_j) = -\rho \partial_i \Phi - \partial_i P, \quad (2.6)$$

where the subscripts  $i$  denote components and repeated subscripts are understood to be summed over. Interpret the terms on the left-hand side in terms of conservation of momentum.

## 2.2 Estimates of solar properties

From equations (2.1) and (2.5) we are in a position to estimate, in an order-of-magnitude sense, many of the stellar properties. First, let's consider the scale for each term in equation (2.5),

$$\underbrace{\partial_t \mathbf{u}}_{\text{I}} + \underbrace{\mathbf{u} \cdot \nabla \mathbf{u}}_{\text{II}} = -\underbrace{\nabla \Phi}_{\text{III}} - \underbrace{\frac{1}{\rho} \nabla P}_{\text{IV}}$$

For a "characteristic" velocity  $U$  and lengthscale  $R$ , we see that terms I and II are both of order  $\sim U^2/R$  (the timescale is  $R/U$ ). For term III, we note that  $GM/R^2 = (GM/R)/R \sim U_{\text{esc}}^2/R$ , where  $U_{\text{esc}}$  is the escape velocity. Finally, for term IV,  $(P/\rho)/R \sim c_s^2/R$ , where  $c_s$  is the speed of sound. Hence the typical scales of the terms are

$$\text{I : II : III : IV} \sim U^2 : U^2 : U_{\text{esc}}^2 : c_s^2$$

Unless we are dealing with stellar explosions, the terms on the left-hand side are quite negligible; in this case we must have the two terms on the right-hand side balance, and the star is in hydrostatic balance,

$$\frac{dP}{dr} = -\rho \frac{Gm(r)}{r^2}. \quad (2.7)$$

Note that this does not mean that  $u$  and  $a$  are zero; it simply means that they are not important for establishing the mechanical structure of the star.

**EXERCISE 2.3**— Equation (2.7) must in general be solved numerically for a real equation of state  $P = P(\rho)$ , but it is useful to construct a toy model to gain insight. Suppose the sun has a density profile

$$\rho(r) = \rho_0 \left(1 - \frac{r}{R_\odot}\right)$$

where  $\rho_0$  is the central density. Further suppose that the equation of state is that of an ideal gas with mean molecular weight  $\mu$ . Find the central density, pressure, and temperature in terms of  $M_\odot$ ,  $R_\odot$ , and  $\mu$ . How do they compare with the values for a constant density star? Evaluate them numerically for a solar composition (hydrogen mass fraction of 0.7). Keeping  $M$  and  $R$  fixed, what happens to the central temperature if the composition is transformed to pure helium? If the nuclear reaction rate depends on temperature, what would this do to the luminosity, in the absence of any other changes?

A side benefit of our argument about the scaling of the terms is that  $c_s \sim U_{\text{esc}} \sim (GM_\odot/R_\odot)^{1/2}$ . We can use this to get an estimate of the central temperature of the sun in terms of  $M_\odot$  and  $R_\odot$ :  $T_{\odot,\text{center}} \sim 10^7 \text{ K}$ , assuming that the equation of state is that of an ideal gas,  $P = (n_{\text{ion}} + n_e)k_B T$  (see exercise 2.4).

**EXERCISE 2.4**— Use this scaling to get an estimate of the central temperature of the sun in terms of  $M_\odot$  and  $R_\odot$ , assuming the composition is an ideal ionized hydrogen plasma. What is the numerical value of the temperature?

### *A worked example: free-fall collapse*

It's worthwhile to imagine what would happen if we suddenly turned off pressure support in the sun, say by having a demon replace each particle with a non-interacting cold particle. For spherically symmetric collapse, let's follow the motion of an observer on the surface. The mass interior to the observer is  $M = M_\odot$ , so her equation of motion is

$$\frac{du}{dt} = -\frac{GM}{r(t)^2}. \quad (2.8)$$

Multiplying both sides by  $u = dr/dt$  and integrating gives

$$\frac{1}{2}u^2 = GM \left( \frac{1}{r} - \frac{1}{R} \right),$$

where  $R = r(t=0)$ . Defining  $x = r/R$  gives

$$\frac{dx}{dt} = \left[ 2 \frac{GM}{R^3} \left( \frac{1}{x} - 1 \right) \right]^{1/2}. \quad (2.9)$$

Now,  $GM/R^3$  has dimension [time<sup>-2</sup>]; furthermore,  $M/R^3 = 4\pi\bar{\rho}/3$ , where  $\bar{\rho}$  is the average density at the start of collapse. Hence, we can define the **dynamical timescale** as  $t_{\text{dyn}} \equiv (G\bar{\rho})^{-1/2}$ . For the sun,  $t_{\text{dyn}} \approx 1 \text{ hr}$ . Defining  $\tau = t/t_{\text{dyn}}$  in equation (2.9) gives us a math problem,

$$\frac{dx}{d\tau} = \left( \frac{8\pi}{3} \right)^{1/2} \left( \frac{1}{x} - 1 \right)^{1/2}$$

which can be integrated from  $x = 1$  to  $x = 0$  to give

$$t_{\text{collapse}} = \left( \frac{3\pi}{32} \right)^{1/2} t_{\text{dyn}} \approx 0.5 \text{ hr}$$

as the time for the sun to collapse if all pressure support were removed.

This is another way of looking at the derivation of eq. (2.7): if terms III and IV are out of balance by even a small amount, the characteristic time for the star to mechanically adjust is very rapid.

For the sun,  $\bar{\rho} = 1.4 \text{ g cm}^{-3}$ , just a bit denser than you.

**EXERCISE 2.5—** Consider a planar atmosphere, in which  $-\nabla\Phi = g = -ge_z$  with  $g$  constant. Thus the equation of hydrostatic equilibrium (eq. [2.7]) is

$$\frac{dP}{dz} = -\rho g. \quad (2.10)$$

Suppose we have an isothermal ideal gas,  $P = \rho k_B T / (\mu m_u)$ , where  $T$  is the temperature,  $k_B$  is Boltzmann's constant, and  $\mu m_u$  is the mass of particles in the gas ( $m_u$  is the atomic mass unit), so that the number of particles per unit volume is  $N/V = \rho / (\mu m_u)$ . Show that for such a gas the density decreases as

$$\rho(z) = \rho(0) \exp(-z/H)$$

and find an expression for the *scale height*  $H$ . Evaluate  $H$  for conditions at sea level on Earth. Does the value make sense? Now evaluate  $H$  under conditions appropriate for the solar photosphere; in this case what is  $H/R_\odot$ ?

### A closer look at hydrostatic equilibrium

If the center of the sun is indeed at a temperature  $\sim 10^7 \text{ K}$ , then most of the gas should be ionized. Now electrons are much lighter than

ions, so we might worry that the charges might separate (their scale heights are different). If that were the case, an electric field would be established. For a pure hydrogen plasma, then, we would have *two* equations of hydrostatic equilibrium, one for the electrons and one for the protons,

$$\nabla P_p = n_p m_p g + n_p e E \quad (2.11)$$

$$\nabla P_e = n_e m_e g - n_e e E. \quad (2.12)$$

Here  $g = -ge_r$  is the gravitational acceleration and  $E$  is the electric field. Notice that if we *presume* that the plasma is charge-neutral, then  $\nabla(P_p + P_e) = \rho g$ , and we can solve for the electric field  $E$ .

**EXERCISE 2.6—** Consider a fully ionized hydrogen plasma in a gravitational field in planar geometry. You may assume that both the protons and electrons each have an ideal, non-degenerate equation of state.

1. Argue that in the absence of an electric field, the protons would sink to the bottom of the atmosphere. Show that if the atmosphere is to remain charge neutral, then an electric field

$$E = -\frac{1}{2} \frac{m_u}{e} g,$$

must be present. Compare this field to that between the proton and electron in an atom. Could this external field be detectable, by Stark effect for example?

2. Suppose a trace ion of charge  $Z'e$  and mass  $A'm_u$  is introduced. What is the net force on this ion?
3. In order to have an electric field, there must be some charge separation. Quantify this: define a parameter

$$\delta \equiv \frac{n_e - n_p}{n_e + n_p}$$

and estimate its magnitude. *Hint:* Use Poisson's equation for both the gravitational and electrostatic potentials, and the results of part (1).

Of course, we must have some charge separation in order to establish the electric field in the first place, but one can show that the fractional charge separation needed is self-consistently small.

### 2.3 Energy considerations

For a spherically symmetric gaseous body in hydrostatic equilibrium, the mass enclosed by radius  $r$  satisfies the differential equation  $dm/dr = 4\pi r^2 \rho$ . Solving for  $\rho$ , substituting into the equation for hydrostatic balance, eq. (2.7), and rearranging terms gives

$$4\pi r^3 \frac{dP}{dr} = -\frac{Gm(r)}{r} \frac{dm}{dr}.$$

Integrating both sides from  $r = 0$  to  $r = R$ , and changing variables on the right hand side from  $r$  to  $m$  gives

$$\int_0^R 4\pi r^3 \frac{dP}{dr} dr = -3 \int_V P dV = - \int_0^M \frac{Gm}{r(m)} dm = E_{\text{grav}}, \quad (2.13)$$

where we integrated the left-hand side by parts, used the fact  $P(R) \ll P(0)$ , and replaced  $4\pi r^2 dr$  with  $dV$ . Now the pressure is related to the internal thermal (kinetic) energy per unit volume  $U$ . For a non-relativistic ideal gas,  $P = 2U/3$ ; for a relativistic gas, such as photons,  $P = U/3$ . Defining  $\gamma = (P + U)/U$ , we can write the total energy of our gaseous sphere as

$$\begin{aligned} E &= E_{\text{th}} + E_{\text{grav}} = \int U dV - 3 \int P dV \\ &= \frac{1 - 3(\gamma - 1)}{\gamma - 1} \int P dV = \frac{3(\gamma - 1) - 1}{3(\gamma - 1)} E_{\text{grav}}. \end{aligned} \quad (2.14)$$

This is just an application of the virial theorem to our star.

As a first example, consider a star with the pressure provided by a non-relativistic ideal gas. Then  $\gamma = 5/3$  and the total energy is<sup>4</sup>

$$E = \frac{1}{2} E_{\text{grav}} < 0.$$

The star is bound. As a second example, consider a star that is so luminous that radiation pressure dominates. In this case, the pressure is that of a relativistic ideal gas. Then  $\gamma = 4/3$  and  $E = 0$ : the star is marginally bound. We must worry about the stability of very luminous stars!

Now suppose the sun were to slowly contract, such that we can still assume hydrostatic equilibrium. How long would this take? The time needed to radiate away the thermal energy defines the **Kelvin-Helmholtz timescale**,

$$t_{\text{KH}} \equiv \frac{E_{\text{th}}}{L} \approx \frac{GM_{\odot}^2}{2R_{\odot}L_{\odot}} = 16 \text{ Myr}. \quad (2.15)$$

We have written “approximately” because we made the approximation that  $E_{\text{grav}} = -GM_{\odot}^2/R_{\odot}$ ; in reality, the density profile of the sun is such that  $E_{\text{grav}}$  is closer to  $-(3/2)GM_{\odot}^2/R_{\odot}$ . The estimated timescale is much less than the age of the earth, and fossils indicate that the sun has not changed dramatically on this timescale. Hence there is an energy source needed to maintain the star’s interior in thermal steady-state. To estimate the scale of energy released per, integrated over the lifetime of the sun, is

$$\frac{\Delta E}{N} \approx \frac{L_{\odot} \times 4.6 \text{ Gyr}}{N} \approx 0.2 \text{ MeV}.$$

This is much larger than chemical reactions<sup>5</sup> could provide. The sun must be powered by nuclear reactions.

<sup>4</sup> This is true even if the matter is degenerate.

<sup>5</sup> typical energy scale is 1 eV

## 2.4 Some analytical limits

We can use the virial theorem of the previous section to set a few limits on the interior pressure and temperature of any star. First, the mass  $m(r)$  inside a volume of radius  $r$  is

$$m(r) = 4\pi \int_0^r \rho r^2 dr,$$

so  $dm/dr = 4\pi r^2 \rho$ . Combining this with the equation of hydrostatic equilibrium gives,

$$\frac{dP}{dm} = -\frac{Gm}{4\pi r^4}.$$

Integrating this equation from the center, where  $P = P_c$ , to some radius  $r$  gives

$$P_c - P(r) = \frac{G}{4\pi} \int_0^r \frac{m dm}{r^4}. \quad (2.16)$$

Now, the average density enclosed in a sphere of radius  $r$  is  $\bar{\rho}(r) = 3m(r)/(4\pi r^3)$ ; solving for  $r$  and inserting in equation (2.16) gives

$$P_c - P(r) = \left(\frac{4\pi}{3}\right)^{4/3} \frac{G}{4\pi} \int_0^{m(r)} \bar{\rho}(r)^{4/3} m^{-1/3} dm. \quad (2.17)$$

Now, the density must decrease outward if the system is to be stable (you can't have heavy fluid on top of light!) and so the average density  $\bar{\rho}(r)$  must also decrease outward. Hence,

$$\rho_c \geq \bar{\rho}(r) \geq \bar{\rho}(R) = \frac{3M}{4\pi R^3}.$$

Inserting this inequality into equation (2.17) and evaluating at  $r = R$  gives a constraint on the central pressure,

$$\frac{3}{8\pi} \frac{GM^2}{R^4} \leq P_c \leq \frac{1}{2} \left(\frac{4\pi}{3}\right)^{1/3} \rho_c^{4/3} GM^{2/3}. \quad (2.18)$$

The critical point here is to notice the order-of-magnitude scale:  $P_c \sim GM^2/R^4$ . The only assumption in setting the limits (eq. [2.18]) is that the density decreases outward.

**EXERCISE 2.7**— Compute the mean kinetic (thermal) energy per hydrogen nucleus in the sun, and express it in electron volts. How does this compare to the gravitational binding energy of a hydrogen atom?

## 2.5 Transport of energy

We derived that at the current luminosity, the sun would take  $\sim 16$  Myr to radiate away its internal (thermal) energy. This raises an

interesting question: what sets the luminosity? To develop this idea further, let us write the luminosity as

$\text{luminosity} \sim (\text{radiation energy stored in sun}) / (\text{photon escape time}).$

To get the radiation energy stored in the sun, we multiply the energy density of a thermal distribution of photons, at the central temperature of the sun, by the volume of the sun:

$$E_\gamma = aT_c^4 \times \frac{4\pi}{3}R_\odot^3. \quad (2.19)$$

What about the photon escape time? As a first try, suppose the sun were transparent, so that photons could freely stream out. Then the escape time would simply be  $R_\odot/c$ . This gives a ridiculously large luminosity.

Suppose now instead that each photon can only travel a short distance  $\ell$  before it is scattered into some random direction. In a random walk, the total distance the photon travels to escape is  $R_\odot(R_\odot/\ell)$ . In this case, the flux from the sun would be

$$F = \frac{L_\odot}{4\pi R_\odot^2} \sim \frac{(4\pi/3)R_\odot^3 aT_c^4}{4\pi R_\odot^2} \frac{c\ell}{R_\odot^2} = \frac{1}{3}\ell c \frac{aT_c^4}{R_\odot}. \quad (2.20)$$

This is very crude, but we can use it to estimate that  $\ell \sim 10^{-3}$  cm. The average distance a photon can travel before being absorbed or scattered is called its **mean free path**. Given this value of  $\ell$  estimated from eq. (2.20), we can estimate the total number of scatterings a photon must suffer in escaping; it is a very large number, and the sun is quite opaque.

Of course, a transparent sun would also not produce a thermal spectrum, since there would be no way for the photons to come into thermal equilibrium with the matter.

---

**EXERCISE 2.8**— If we regard the sun as a large cavity filled with photons, estimate the total energy stored in the radiation field. If the sun were suddenly to become completely transparent, what would be the resulting luminosity?

---

## 2.6 Summary

In summary, we've taken the observed gross properties of the sun, the equation of motion for a fluid, and the ideal gas equation of state; from these we've deduced that the sun is in hydrostatic balance, that its interior temperature is of order  $10^7$  K, and that it would radiate away its thermal energy and contract within about 10 Myr if there were no nuclear reactions in its core. We have developed now a crude picture of the sun: it is a mass of plasma that is in hydrostatic equilibrium, with pressure gradients supporting the inward pull

of gravity. It is very opaque, and thus it acts as a reservoir for photons with a thermal, Planckian distribution. Because it is opaque, the photons leak out very slowly. This slow leakage represents a loss of thermal energy; the thermal energy is replenished by heat liberated from nuclear reactions.

What comes next is fleshing out the detailed physics implied by these considerations: an equation of state to relate the pressure to the density and temperature; photon scattering and absorption cross sections to compute the heat transport; nuclear reaction rates to determine the thermal steady-state and the gradual change in composition of the interior.

### **MESA** *A contracting pre-main-sequence star*

#### *Setting up your workspace*

To install MESA on your personal linux or mac computer do the following.

1. Download and install the appropriate SDK, available from <http://www.astro.wisc.edu/~townsend/static.php?ref=mesasdk>.
2. Use svn to checkout the latest release verion of MESA. Instructions are available from <http://mesa.sourceforge.net>.
3. From a terminal window, go into the top-level MESA directory and execute the following commands: `'./clean; ./install'`. This may take some time as MESA will be building fairly large datasets for the equation of state and opacity. If all goes well, you will get a message at the end indicating that MESA installation was successful.

Once you have MESA installed, you then set up your work environment. To do this, you create a project directory and set environment variables so that your FORTRAN compilers can find the MESA libraries.

#### *Running your first MESA project*

You are now ready to compile and run your first MESA project for this course. This is a model of a  $1 M_{\odot}$  pre-main-sequence star. The code stops when the luminosity from hydrogen fusion first exceeds 0.95 of the luminosity from the surface.

Download the folder `1M-pms` and place it into your projects folder.

Now execute `'./mk'`. This will call the fortran compiler to build the executable.

If everything compiles okay, then an executable file called 'star' will be placed in the directory. To run MESA, type `'./rn'` at the prompt.

If all goes well, after some time a window should appear with an animated plot looking something like Figure 2.1.

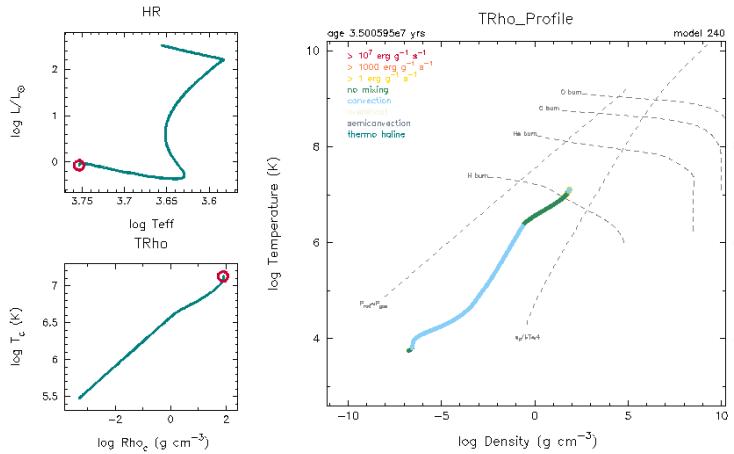


Figure 2.1: Graphical output from a MESA run of a  $1 M_{\odot}$  PMS star.

The large plot labeled ‘TRho\_profile’ shows the run of temperature  $T$  versus density  $\rho$  in the star, with the various colors indicating mixing and energy generating regions. The small plot labeled ‘HR’ traces the history of luminosity  $L$ , in units of the solar luminosity  $L_{\odot}$ , versus effective temperature  $T_{\text{eff}}$ . The other small plot labeled ‘TRho’ traces the history of the central temperature  $T_c$  versus the central density  $\rho_c$ .

#### EXERCISE 2.9—

1. How long did the star take to contract to the main sequence?
2. What are  $L$ ,  $T_{\text{eff}}$ ,  $T_c$ , and  $\rho_c$  when the star begins H burning?
3. Describe how the fraction of the star that is convective changes during the run.

#### *A peak under the hood*

To understand what just happened, we start with the command ‘./rn’. This is just a script—you can open it with a text editor—containing the following.

```

1 #!/bin/bash
2
3 rm -f restart_photo
4 date "+DATE: %Y-%m-%d%TIME: %H:%M:%S"
5 if [[ -e star.exe ]];then
6     ./star.exe
7 else

```

```

8      ./star
9  fi
10 date "+DATE: %Y-%m-%d%nTIME: %H:%M:%S"

```

This just gets some information about the version of MESA being used (the ‘svn info’ directive), removes any pre-existing file with restart information, prints the date, runs star, and then prints the date again.

The command ‘star’ is built from the source code in ‘src/run.f’. This is a very short program. The relevant lines

```

11      inlist_fname = 'inlist'
12
13      call do_read_star_job(inlist_fname, ierr)
14      if (ierr /= 0) stop 1
15
16      call do_run_star(inlist_fname)

```

direct the program to read in parameters from a file ‘inlist’ and then hand control to a subroutine, ‘do\_run\_star’, within the MESA library.

The file ‘inlist’, is divided into three sections (each section begins with an ‘&’ followed by the section name and ends with a ‘/’): ‘star\_job’, ‘controls’, and ‘pgstar’. The bang ‘!’ denotes the start of a comment. This list is rather simple. The first section

```

1 &star_job
2
3     read_extra_star_job_inlist1 = .true.
4     warn_run_star_extras=.false.
5     extra_star_job_inlist1_name = '1M_pms.inlist'
6
7 / ! end of star_job namelist
8
9 &controls

```

tells MESA to read another file, ‘1M\_pms.inlist’. The second section, ‘&controls’ also tells MESA to read in ‘1M\_pms.inlist’. It is in ‘1M\_pms.inlist’ that all of the settings are placed, so take a quick peek at that file. For example, in section ‘&star\_job’ the lines

```

21     pause_before_terminate = .true.
22     pgstar_flag = .true.

```

tell MESA to pause and wait for the user to hit ‘return’ before ending the code, and to activate the plotting windows. In section ‘&controls’ the lines

```

31     initial_mass = 1.0
32     initial_z = 0.02d0

```

tell MESA the initial mass and metallicity of the star, while the lines

```
42     stop_near_zams = .true. ! stop if Lnuc/L > Lnuc_div_L_zams_limit
43     Lnuc_div_L_zams_limit = 0.95d0
```

tell MESA to stop when the total power from nuclear burning,  $L_{\text{nuc}}$ , exceeds  $0.95L$ , where  $L$  is the total surface luminosity.

The third section of ‘inlist’

```
18     read_extra_pgstar_inlist1 = .true.
19     extra_pgstar_inlist1_name = 'basic_plot.inlist'
20
21     read_extra_pgstar_inlist2 = .false.
22     extra_pgstar_inlist2_name = 'track_scaled_vars.inlist'
23
24 / ! end of pgstar namelist
```

reads in the parameters for the plot from ‘basic\_plot.inlist’. There is another parameter file, ‘track\_scaled\_vars.inlist’, but the flag to read that file is ‘read\_extra\_pgstar\_inlist2 = .false.’ As you might guess, you’ll be using this later on in the assignment.

**Question:** There is a reason for nesting the parameter inlist files. Can you discern what that reason is?

Now that you’ve seen how the code in action, we are going to look a bit more at the architecture of MESA. If you do ‘ls \$MESA\_DIR’, you will see that MESA is divided into modules: ‘eos’ computes the equation of state, ‘kap’ computes the opacity, and so forth. Within each module are two folders, ‘public’ and ‘private’. The ‘public’ folder contains the interface of that module. The source file ending with ‘\_def.f’ contains the data structures used by that module, and the source file ending with ‘\_lib.f’ contains the routines for that module. The private directory contains the inner machinery of the module.

All of these modules can be used by themselves; the ‘star’ module puts everything together to simulate stellar evolution. What star does is to evolve a stellar model—a complete description of a star at a given instant of time—forward in time by some amount  $\Delta t$ . String together a sequence of such models and you have a representation of the star’s evolution. These models are not evenly spaced in time; rather, MESA adjusts  $\Delta t$  to keep the models accurate within specified tolerances. The star module contains, in addition to the public interface and private machinery, additional routines in the folder job for starting a run from some initial model and stopping that run when a specified condition is met.

### *A MESA project*

The MESA code that you installed is a library, a collection of routines that when combined simulate the evolution of a stellar-like object. To put everything together, you create a directory, such as ‘1M-pms’. A template for such a directory is contained in ‘\$MESA\_DIR/star/work’—consult the ‘README’ file there for instructions.

The working directory is organized into several sub-directories. The ‘make’ folder contains the ‘makefile’ script for compiling the code. The ‘src’ folder contains, in addition to the top-level ‘run.f’ code, a collection of customizable routines in the file ‘run\_star\_extras.f’. In addition to these folders, the working directory contains a set of inlist files; these, as mentioned above, contain all of the parameters necessary to control the MESA run and its output. The complete listings of parameters and their default settings are contained in the directory ‘\$MESA\_DIR/star/defaults’ in the three files ‘\*.defaults’. The inlists in the work directory only need to contain those parameters that differ from the defaults.

The final components of the working directory are sub-directories to hold the output of MESA. The names of these are customizable and can be set in the inlists; by default, the main two are called ‘LOGS’ and ‘photos’. Within ‘LOGS’ are the file ‘history.data’ and the files ‘profiledd.data’. The ‘history.data’ file contains the time evolution of global stellar properties, such as luminosity, radius, surface effective temperature, and so on. The ‘profiledd.data’ files contain “snapshots” of the star’s structure: the run of temperature, density, pressure, and so on with location within the star.

### *Exercise: customizing MESA output*

After that brief overview of the MESA architecture, let’s do something concrete: we’ll customize MESA to have it generate a plot of a variable that we define. In this chapter, we’ve argued that the central pressure of a star should scale as  $P_c \propto GM^2/R^4$ . We also found that the central density should scale as the mean density,  $\rho_c \propto \bar{\rho} = 3M/(4\pi R^3)$ . Exercise 2.4 asks you to find the central temperature in terms of  $M$ ,  $R$ , and mean molecular weight  $\mu$ . What the derivation in the chapter doesn’t tell us is the coefficient  $\rho_c/\bar{\rho}$  and its counterparts for pressure and temperature. We can use MESA, however, to test these scalings and extract these constants of proportionality.

To do this, we modify the code in ‘src/run\_star\_extras.f’. What we want is for MESA to calculate  $P_{\text{scale}} = GM^2/R^4$ ,  $\rho_{\text{scale}} = \bar{\rho}$ , and  $T_{\text{scale}}$ , and then write out the values of  $P_c/P_{\text{scale}}$ ,  $\rho_c/\rho_{\text{scale}}$ , and  $T_c/T_{\text{scale}}$  in the file ‘history.data’. To do this, we first tell MESA how many extra columns in ‘history.data’ we need:

```

96     integer function how_many_extra_history_columns(id, id_extra)
97         integer, intent(in) :: id, id_extra
98         integer :: ierr
99         type (star_info), pointer :: s
100        ierr = 0
101        call star_ptr(id, s, ierr)
102        if (ierr /= 0) return
103        how_many_extra_history_columns = 1 !3
104    end function how_many_extra_history_columns

```

Here I am adding one column, which will be  $P_c/P_{\text{scale}}$ . When you implement the other scalings, you'll change the variable `how_many_extra_history_columns` to reflect that we need three columns.

Next, we need to compute the data for these columns. We therefore modify the following routine.

```

107    subroutine data_for_extra_history_columns(id, id_extra, n, names, vals, ierr)
108        integer, intent(in) :: id, id_extra, n
109        character (len=maxlen_history_column_name) :: names(n)
110        real(dp) :: vals(n)
111        integer, intent(out) :: ierr
112        type (star_info), pointer :: s
113        real(dp) :: G, M, R, mu, Pscale, rhoscale, Tscale
114
115        ierr = 0
116        call star_ptr(id, s, ierr)
117        if (ierr /= 0) return
118
119        !note: do NOT add the extras names to history_columns.list
120        ! the history_columns.list is only for the built-in log column options.
121        ! it must not include the new column names you are adding here.
122
123
124        names(1) = 'Pc_scaled'
125        ! names(2) = 'Tc_scaled'
126        ! names(3) = 'rhoc_scaled'
127
128        ! Newton's constant, defined in module const_def (which
129        ! is included at the top)
130        G = standard_cggrav
131        ! mass of the star, in g; s is a data structure and is
132        ! described in $MESA_DIR/star/public/star_data.inc
133        M = s% mstar
134        ! log_surface_radius is in units of Rsol, so we convert; rsol
135        ! is in const_def

```

```

136      R = rsol*10.0**s% log_surface_radius
137      ! central mean molecular weight
138      mu = s% center_mu
139      ! scalings for pressure, density, and temperature
140      Pscale = G*M**2/R**4
141      Tscale = ?
142      rhoscale = ?
143      ***HINT***: the combination kB*NA, where kB = Boltzmann's constant
144      and NA = Avogadro's number, is defined in const_def is given the name
145      cgas. You will need this to compute Tscale.
146
147
148      vals(1) = 10.0**s% log_center_pressure / Pscale
149      !     vals(2) = ?
150      !     vals(3) = ?
151
152
153
154 end subroutine data_for_extra_history_columns

```

In detail, we first declare the extra variables we need.

```
113 real(dp) :: G, M, R, mu, Pscale, rhoscale, Tscale
```

Next, we give each column a name.

```

124 names(1) = 'Pc_scaled'
125 !     names(2) = 'Tc_scaled'
126 !     names(3) = 'rhoc_scaled'

```

Notice that the latter two are commented out ('!'); for this example; you'll need to uncomment them to print out the other variables.

We are then ready to compute our values. Note that MESA defines many physical constants in '\$MESA\_DIR/const/public/const\_def.f'; we therefore use this in line 130 for  $G$ , and you can read the hint in lines 143ff. Next, we need to get the values of  $M$ ,  $R$ ,  $\mu_c$ ,  $P_c$ ,  $\rho_c$ , and  $T_c$ . These are provided in a data structure  $s$ , which the routine fetches

```
116 call star_ptr(id, s, ierr)
```

from the variable  $id$  that is passed to the routine. To see a complete list of what is in this data structure, look at '\$MESA\_DIR/star/public/star\_data.inc'. I've already taken care of computing  $M$ ,  $R$ ,  $\mu_c$ , and  $P_{\text{scale}}$  for you. The values of our scaled variables are then stored in the array 'vals'; you will fill in the second and third members of the array.

If you look in 'LOGS/history.data', you should see that the last column is indeed 'Pc\_scaled', as expected. Of course, we'd like

to display it graphically, and MESA does predefine plots that can display values in ‘history.data’. We customized the output in ‘track\_scaled\_vars.inlist’:

```

1  &pgstar

2

3      History_Panels2_win_flag = .false.
4      History_Panels2_win_width = 6
5      History_Panels2_win_aspect_ratio = 0.75 ! aspect_ratio = height/width

6

7      History_Panels2_title = 'Scaled central values'

8

9      History_Panels2_num_panels = 1

10

11     History_Panels2_file_flag = .false.
12     History_Panels2_file_dir = 'frames'
13     History_Panels2_file_prefix = 'scaled_center_vals_'

14

15     History_Panels2_yaxis_name(1) = 'Pc_scaled'
16 !     History_Panels2_yaxis_name(2) = 'Tc_scaled'
17 !     History_Panels2_yaxis_name(3) = 'rhoc_scaled'

18

19 / ! end of pgstar namelist

```

To use this, we activate the window by setting the flag in line 3 to ‘.true.’. If we want to save our output, we set the ‘file\_flag = .true.’ in line 11.



# 3

## *The Lagrangian Equations of Stellar Structure*

### 3.1 The conservation laws

After our rapid overview, we now gather the tools needed to tackle stellar evolution. The first is to get the macroscopic equations for stellar structure. We will start from the equations expressing conservation of mass<sup>1</sup>, momentum, and energy. We already derived the continuity (conservation of mass) equation,

$$\partial_t \rho + \nabla \cdot (\rho \mathbf{u}) = 0, \quad (3.1)$$

and the Euler equation,

$$\partial_t \mathbf{u} + \mathbf{u} \cdot \nabla \mathbf{u} = -\nabla \Phi - \frac{1}{\rho} \nabla P. \quad (3.2)$$

Note that if we multiply eq. (3.2) by  $\rho$ , we can rewrite it, using eq. (3.1), as

$$\partial_t(\rho \mathbf{u}) + \nabla \cdot [\mathbf{u}(\rho \mathbf{u})] = -\rho \nabla \Phi - \nabla P. \quad (3.3)$$

The left-hand side is interpreted as expressing the conservation of momentum ( $\rho \mathbf{u}$ ) in the absence of forces, analogous to eq. (3.1) for the conservation of mass ( $\rho$ ).

Note the general form of a conservation equation:

$$\begin{aligned} & \partial_t(\text{conserved quantity}) \\ & + \nabla \cdot (\text{flux of conserved quantity}) = (\text{sources}) - (\text{sinks}). \end{aligned}$$

Because the momentum density  $\rho \mathbf{u}$  is a vector, its flux is a tensor:

$$[\mathbf{u}(\rho \mathbf{u})]_{ij} \equiv \rho u_i u_j.$$

THE NEXT EQUATION IS THAT OF ENERGY CONSERVATION. Here we must consider both the internal energy per unit volume  $E/V = \rho \varepsilon$  and the kinetic energy per unit volume  $\rho u^2/2$ . In this section  $\varepsilon$  represents the internal energy per unit mass of the fluid. In a fixed

<sup>1</sup> In a relativistic system, we would instead start from conservation of baryon number, since mass is not invariant.

By  $\rho u_i u_j$ , we mean the momentum along direction  $i$  being transported along direction  $j$ .

volume of the fluid the total energy is thus

$$\int_V \left( \rho \frac{1}{2} u^2 + \rho \varepsilon \right) dV.$$

The flux of energy into this volume will clearly include

$$-\int_{\partial V} \left( \frac{1}{2} \rho u^2 + \rho \varepsilon \right) \mathbf{u} \cdot d\mathbf{S}.$$

But wait, there's more! In addition, we have a conductive heat flux  $\mathbf{F}$ ; the total heat conducted through the surface  $\partial V$  is

$$-\int_{\partial V} \mathbf{F} \cdot d\mathbf{S}.$$

Moreover, the pressure acting on fluid flowing into our volume does work on the gas at a rate

$$-\int_{\partial V} P \mathbf{u} \cdot d\mathbf{S}.$$

As a result, the net change of energy in our volume is

$$\begin{aligned} \partial_t \int_V \left( \frac{1}{2} \rho u^2 + \rho \varepsilon \right) dV = \\ - \int_{\partial V} d\mathbf{S} \cdot \left[ \mathbf{u} \left( \frac{1}{2} \rho u^2 + \rho \varepsilon + P \right) + \mathbf{F} \right] \\ + \int_V (\rho \mathbf{u} \cdot \mathbf{g} + \rho q) dV. \end{aligned} \quad (3.4)$$

On the right-hand side we've added in the work done by gravity and the heating evolved by nuclear reactions (this could also involve sinks, such as neutrinos with a long mean free path). Expressed in differential form, equation (3.4) is

$$\partial_t \left( \frac{1}{2} \rho u^2 + \rho \varepsilon \right) + \nabla \cdot \left[ \rho \mathbf{u} \left( \frac{1}{2} u^2 + \varepsilon + \frac{P}{\rho} \right) \right] + \nabla \cdot \mathbf{F} = \rho q + \rho \mathbf{u} \cdot \mathbf{g}. \quad (3.5)$$

You are possibly wondering why I didn't put gravity, which can be expressed as a potential, on the left hand side of this equation. The reason is that the gravitational stresses cannot be expressed in a *locally* conservative form; it is only when integrating over all space that the conservation law appears.

EQUATIONS (3.1), (3.3), AND (3.5) ARE SUPPLEMENTED BY AN EQUATION OF STATE, which allows one to get from the pressure  $P$ , the temperature  $T$ , and the mass fractions  $X_i$  of the species present, the remaining thermodynamical quantities, such as mass density  $\rho$  and specific energy  $\varepsilon$ . In addition, Poisson's equation

$$\nabla^2 \Phi = 4\pi G \rho, \quad (3.6)$$

specifies the gravitational acceleration  $\mathbf{g} = -\nabla\Phi$ . We then need one more equation to specify the heat flux  $\mathbf{F}$ . We argued in §2.5 that the typical length over which a photon travels before scattering is very small compared to the lengthscale over which the macroscopic properties of the star vary. In this case, we expect the flux to obey a conduction equation of the form

$$\mathbf{F} = -K\nabla T. \quad (3.7)$$

This assumption is clearly questionable near the stellar surface, and we have left unspecified the form of  $K$ . Such an equation does, however, close the system of equations; all of the physics is then contained in the equation of state  $P(\rho, T, \{X_i\})$ , the rate of heating from nuclear reactions  $q(\rho, T, \{X_i\})$ , and the thermal conductivity  $K(\rho, T, \{X_i\})$ . Here  $\{X_i\}$  are the mass fractions of the isotopes composing the solar plasma. We will also need a system of equations to describe how the  $X_i$  change as a result of nuclear reactions and diffusion.

### 3.2 Thermodynamics of a mixture: A digression

#### Specifying the composition

In this section we'll look at how one describes the composition for a multi-component plasma. To make things concrete, let's imagine a box containing a mixture of nuclei, of many different isotopes, and electrons. (To keep things simple, we'll assume complete ionization.) Each isotope species  $i$  has  $N_i$  nuclei present, and is characterized by charge number  $Z_i$  and nucleon number  $A_i$ . Charge neutrality then specifies the number of electrons,

$$N_e = \sum_i Z_i N_i. \quad (3.8)$$

The total mass of the box is

$$M = m_e N_e + \sum_i m_i N_i, \quad (3.9)$$

where  $m_e$  and  $m_i$  are respectively the mass of an electron and a nucleus of species  $i$ . Now what is  $m_i$ ? Breaking a nucleus  $i$  into  $Z_i$  protons and  $A_i - Z_i$  neutrons takes a certain amount of energy, the *binding energy*  $B_i$ . We can therefore write  $m_i = Z_i m_p + (A_i - Z_i) m_n - B_i/c^2$ , where  $m_p$  and  $m_n$  are respectively the proton and neutron rest masses.

Inserting our expression for  $m_i$  into equation (3.9), dividing by the volume of the box  $V$ , and rearranging terms gives us the mass

density,

$$\rho = \frac{M}{V} = \sum_i n_i \left[ (A_i - Z_i) m_{\text{n}} + Z_i (m_{\text{p}} + m_{\text{e}}) - B_i/c^2 \right]. \quad (3.10)$$

Here  $n_i$  is the number density of isotope species  $i$ , and we have used equation (3.8) to eliminate  $N_{\text{e}}$ . The numbers  $n_i$  are, of course, fantastically<sup>2</sup> large, so we scale the numbers by *Avogadro's constant*,

$$N_{\text{A}} = 6.0221367 \times 10^{23} \text{ mol}^{-1}. \quad (3.11)$$

If we multiply and divide the right-hand side of equation (3.10) by  $N_{\text{A}}$ , we then have

$$\rho = \sum_i \left( \frac{n_i}{N_{\text{A}}} \right) \mathcal{A}_i, \quad (3.12)$$

where

$$\mathcal{A}_i = \left[ (A_i - Z_i) m_{\text{n}} + Z_i (m_{\text{p}} + m_{\text{e}}) - B_i/c^2 \right] \times N_{\text{A}} \quad (3.13)$$

is the *gram-molecular weight* of species  $i$  with dimensions  $[\mathcal{A}] \sim [\text{g} \cdot \text{mol}^{-1}]$ . Strictly speaking, the gram-molecular weight actually refers to the mass of a mole of the isotope in *atomic* form; the right-hand side of eq. (3.13) is the gram-molecular weight neglecting the electronic binding energy.

Now you may wonder where the numerical value of  $N_{\text{A}}$  came from. It is not pulled out of thin air, but rather is defined so that 1 mol of  $^{12}\text{C}$  has a mass of exactly 12 g. In other words, for  $^{12}\text{C}$   $\mathcal{A} \equiv A \text{ g mol}^{-1}$ . In fact for all nuclei,  $\mathcal{A} \approx A \text{ g mol}^{-1}$  to better than about 1%, as demonstrated in Table 3.1. Because in CGS  $\mathcal{A} \approx A$ , it is customary to write  $\mathcal{A} = A \times (1 \text{ g mol}^{-1})$ , so that equation (3.12) is

$$\rho = \sum_i \left( \frac{n_i}{N_{\text{A}}} \times 1 \frac{\text{g}}{\text{mol}} \right) \mathcal{A}_i. \quad (3.14)$$

This only works if our unit of mass is the gram: in SI units  $^{12}\text{C}$  has a mass of  $0.012 \text{ kg mol}^{-1}$ . Equation (3.14) would be exact if  $A$  were a real number, but the custom is to just keep it as the nucleon number, which introduces an error of order one percent. Astronomers typically then *redefine*  $N_{\text{A}}$  to mean  $N_{\text{A}}(\text{astronomy}) \equiv N_{\text{A}}/(1 \text{ g mol}^{-1}) = 6.0221367 \times 10^{23} \text{ g}^{-1}$ . Alternatively, one can use the atomic mass unit (symbol u) defined as 1/12 the mass of an atom of  $^{12}\text{C}$ , so that  $1 \text{ u} = (1 \text{ g mol}^{-1})/N_{\text{A}} = 1.66054 \times 10^{-24} \text{ g}$ . This puts equation (3.14) into the more obvious form  $\rho = \sum n_i \times A_i m_{\text{u}}$ , with  $m_{\text{u}}$  having a mass of 1 u.

With the redefinition of  $N_{\text{A}}$ , equation (3.14) can be rewritten as

$$1 = \sum_i \left( \frac{n_i}{N_{\text{A}} \rho} \right) \mathcal{A}_i \equiv \sum_i Y_i \mathcal{A}_i \quad (3.15)$$

<sup>2</sup> astronomically?

Recall that a **mole** is an amount of something: in 1 mol there are  $N_{\text{A}}$  items.

nuclide	$A$	$\mathcal{A}$	$( \mathcal{A} - A /A) \times 100$
n	1	1.00865	0.865
$^1\text{H}$	1	1.00783	0.783
$^4\text{He}$	4	4.00260	0.065
$^{16}\text{O}$	16	15.99491	0.032
$^{28}\text{Si}$	28	27.97693	0.082
$^{56}\text{Fe}$	56	55.93494	0.116

Table 3.1: Selected gram-molecular weights.

where  $Y_i \equiv n_i/(\rho N_A)$  is the *molar fraction*. It is customary to call  $Y_i A_i$  the *mass fraction*  $X_i$ , with  $\sum X_i = 1$ . We can then define the mean atomic mass number,

$$\bar{A} = \frac{\sum A_i Y_i}{\sum Y_i} = \frac{1}{\sum Y_i}, \quad (3.16)$$

and mean charge number

$$\bar{Z} = \frac{\sum Z_i Y_i}{\sum Y_i} = \bar{A} \sum Z_i Y_i. \quad (3.17)$$

The molar fraction of electrons is

$$Y_e = \sum Z_i \frac{n_i}{\rho N_A} = \sum Z_i Y_i = \frac{\bar{Z}}{\bar{A}}. \quad (3.18)$$

In stellar structure work, it is common to use the *mean molecular weight*, defined so that the total number of particles, including electrons, per unit volume is

$$\sum_i n_i + n_e \equiv \frac{\rho N_A}{\mu}. \quad (3.19)$$

Yes, this is still the redefined  $N_A$ :  $\mu$  is dimensionless. From the definition,

$$\mu = \left( \sum_i Y_i + Y_e \right)^{-1} = \left[ \sum_i (Z_i + 1) Y_i \right]^{-1}; \quad (3.20)$$

sometimes astronomers also define the mean ion molecular weight,  $\mu_I = (\sum Y_i)^{-1}$ , and the mean electron weight,  $\mu_e = Y_e^{-1}$ .

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**EXERCISE 3.1**— Consider a gas of  $^1\text{H}$  and  $^4\text{He}$  with molar hydrogen fraction  $Y_{\text{H}}$ . Derive expressions for the molar fraction of  $^4\text{He}$ ,  $Y_{\text{He}}$ ,  $\bar{A}$ ,  $\bar{Z}$ , and  $\mu$ . What are the numerical value of these quantities for  $Y_{\text{H}} = 0.7$ , i. e., solar?

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**EXERCISE 3.2**— Assume that we can describe this plasma as an ideal gas. What is the sound speed and the average kinetic energy of a particle, for a given mass density and temperature?

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### 3.3 The Equation of State of an Ideal Gas

Equations (3.1), (3.3), and (3.5) must be closed by a suitable *equation of state* that relates the thermodynamic quantities to each other. A simple equation of state is that of an ideal gas. An *ideal gas* is assumed to be made of vast numbers of point-like particles that interact only via elastic collisions. This is generally a pretty good approximation for a monatomic, single-species gas but also works fairly well for molecular real gasses such as the terrestrial atmosphere, within certain limits. Typically, a gas behaves more ideal at higher temperatures at which the kinetic motion of the particles dominate any intermolecular forces (i.e., van der Waals) that the gas may experience. Ideal gasses obey the *ideal gas law*. A *perfect gas* is a special case of an ideal gas that has constant specific heat capacities at constant volume and at constant pressure (see Section 3.4). In general, an ideal gas can have specific heat capacities that temperature- and/or density-dependent.

Consider a closed volume containing a large number of point-like collisional particles. Elastic collisions of these particles with the walls of the volume will result in a force exerted on the walls. In the absence of any external forces (i.e., gravity), the force exerted will be the same for all the walls of the volume. The force per wall area is the *pressure*  $P$  on the walls.

The ideal gas law was first determined empirically by chemists in the 19th century and relates the pressure of a gas to its volume and temperature:

$$PV = N_m RT, \quad (3.21)$$

where  $V$  is the volume of the gas,  $T$  its temperature,  $N_m$  the number of *moles* of gas, and  $R$  is the universal gas constant. It turns out, experimentally, the number of particles in a mole of gas is a universal constant: Avogadro's number  $N_A$ . The universal gas constant is the product of Avogadro's number and the Boltzmann constant  $k_B$ :  $R = N_A k_B$ . Thus, we can write the ideal gas law in a form more familiar to physicists,

$$P = nk_B T, \quad (3.22)$$

where  $n = N_M N_A / V$  is the number density of particles in the gas.

If the particles can be described by a mean atomic mass  $\bar{A}$  (Eq. 3.16) then we can express the mass density in terms of the number density as  $\rho = nAm_u$ , where  $m_u$  is the atomic mass unit. Then another form of the ideal gas law is

$$P = \frac{\rho}{\bar{A}m_u} k_B T. \quad (3.23)$$

For a mixture of different particle types, including free electrons as

we would have in a stellar plasma, we can express the ideal gas law in yet another equivalent form using the mean molecular weight,

$$P = \frac{N_A \rho}{\mu} k_B T = \frac{\rho}{\mu m_u} k_B T, \quad (3.24)$$

where we recognize that  $N_A \approx m_u^{-1}$ . Using the definition of the mean molecular mass in terms of the molar fractions (Eq. 3.20), we see that the total pressure is the sum of partial pressures,

$$P = \left( \sum_i Y_i + Y_e \right) N_A k_B \rho T. \quad (3.25)$$

Equations (3.21-3.25) are all equivalent for a gas of identical particles, with Eqs. (3.24) and (3.25) being more general for a mixture that includes free electrons as well as ions.

### 3.4 Thermodynamics of Stellar Plasmas

#### The First Law

Consider a volume of gas that can exchange energy with its surrounding either by the transfer of heat or by doing *work*. Now assume the gas undergoes a truly *reversible* process, in which the state of the gas changes slowly enough that it is always in equilibrium. I.e., throughout the entire process the gas moves through a series of infinitesimally different equilibrium states. In such a process, the change in the gas internal energy will be

$$dE = dQ - dW, \quad (3.26)$$

where  $dQ$  is the heat gained or lost and  $dW$  is the work done *by* the gas (positive sign) or *on* the gas by its surroundings (negative sign).

First, let's consider how to deal with the work term. If the gas is "contained" within some volume  $V$ , the gas exerts a differential force on the "wall" of this volume,

$$\delta F = P \delta A \hat{n},$$

where  $\delta A$  is some differential area element of the volume walls. Now consider that the force exerted by the gas causes an infinitesimal displacement  $dx$  in the volume walls doing some differential *work* (i.e., force through a distance),

$$\delta(dW) = \delta F \cdot dx = P(\hat{n} \cdot dx)\delta A = P\delta V.$$

Now, integrating over all area of the volume walls gives us  $dW = PdV$ , which you will recall elementary dynamics. Then we have the First Law of Thermodynamics,

$$dE = dQ - PdV. \quad (3.27)$$

## First Law of Thermodynamics

In Eq. (3.27) the energy  $E$  and heat  $Q$  are *extensive* quantities that scale with the number of particles  $N$  in our sample. In a fluid, however, these quantities are all functions of position. By  $E(r)$ , we mean that we can define a small portion of the star about the coordinate  $r$  that is large enough particles to ensure that quantities such as pressure and temperature are well-defined, but small enough that we can treat  $E(r)$  as a continuous function of position when integrating over the whole star.

Using extensive quantities in fluid mechanics is cumbersome, so we instead use quantities like the energy per unit mass  $\varepsilon = E/(\rho V)$  or the heat per unit mass  $q = Q/(\rho V)$ . Since a fixed mass of fluid  $M$  occupies a volume  $V = M/\rho$ , we can divide the first law, eq. (3.27), by  $M$  to obtain

$$d\varepsilon = dq - Pd\left(\frac{1}{\rho}\right) = dq + \frac{P}{\rho^2}d\rho. \quad (3.28)$$

The other extensive variables can be re-defined into mass-specific forms in a similar fashion.

The definition of an *adiabatic* process is one in which there is no change in the heat content of the system,  $dQ \equiv 0$ . Consider the Joule-Kelvin experiment in which a gas is allowed to expand adiabatically ( $dQ = 0$ ) into a larger volume initially a vacuum (i.e.,  $V_1 < V_2$ ). Since the gas experiences no resistance while expanding into the new volume it does not do any work while expanding,  $dW = PdV = 0$ . Therefore, by the First Law (Eq. 3.27), we have  $dE = 0$  implying that the temperature remains constant during the expansion,  $T_2 = T_1$ . Thus, the internal energy of the gas is a function only of temperature,  $E(T)$ . In general, we would expect that the internal energy is also a function of density (or volume),  $E(V, T)$ . This somewhat surprising result only holds for perfect gasses.

## Extensive vs. intensive quantities

## Adiabatic process

*Heat Capacity and Thermodynamic Derivatives*

Now consider a non-adiabatic process in which some heat,  $dq$ , is added to the system resulting in a temperature change  $dT$ . The *specific heat capacity*, or simply specific heat, is then defined as

$$C \equiv \frac{dq}{dT},$$

where, again, the heat per mass is  $q$ . Specific heats are usually specified under conditions in which a thermodynamic quantity is assumed constant. For instance, the specific heat at constant pressure is  $C_P$  and the specific heat at constant volume is  $C_V$ .

A key feature of thermodynamics is that any thermodynamic quantity can be expressed as a function of any other *two* thermodynamic quantity. Consider, for instance, the ideal gas law (e.g., Eq. 3.23), which can obviously be arranged as we like to give  $P(\rho, T)$ , or  $T(\rho, P)$ , or  $\rho(P, T)$ . Thus, it is useful to define derivatives of thermodynamic quantities with respect to others while keeping a third quantity fixed. For variables  $x, y, z$  we would write this using the notation,

$$\left( \frac{\partial z}{\partial x} \right)_y,$$

where the subscript  $y$  is the variable held fixed while calculating the derivative of  $z$  with respect to  $x$ .

Now assume that our variables  $(x, y, z)$  can be connected by some function  $\mathcal{F}(x, y, z) = 0$  (such as you might get by moving the RHS of the ideal gas law to the LHS). Then, we can derive several useful relations (see Appendix A.1):

$$\left( \frac{\partial x}{\partial y} \right)_z = \frac{1}{(\partial y / \partial x)_z}, \quad (3.29)$$

$$\left( \frac{\partial x}{\partial y} \right)_z \left( \frac{\partial y}{\partial z} \right)_x \left( \frac{\partial z}{\partial x} \right)_y = -1. \quad (3.30)$$

Now if some other quantity,  $w$ , is a function of any two of  $(x, y, z)$ , then

$$\left( \frac{\partial x}{\partial y} \right)_w \left( \frac{\partial y}{\partial z} \right)_w = \left( \frac{\partial x}{\partial z} \right)_w. \quad (3.31)$$

Consider as independent variables  $(\rho, T, \varepsilon)$ . Expanding the total differential of  $\varepsilon$  in terms of its partial derivatives w.r.t.  $\rho$  and  $T$ , we have

$$d\varepsilon = \left( \frac{\partial \varepsilon}{\partial T} \right)_\rho dT + \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T d\rho.$$

Using the First Law (Eq. 3.28) we then find

$$dq = \left( \frac{\partial \varepsilon}{\partial T} \right)_\rho dT + \left[ \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T - \left( \frac{P}{\rho^2} \right) \right] d\rho. \quad (3.32)$$

For a process occurring at constant volume, the density is also constant<sup>3</sup>,  $d\rho = 0$ . The change in heat is given by  $dq = C_V dT$ . Then by inspection, we find that the specific heat at constant volume is

$$C_V = \left( \frac{\partial \varepsilon}{\partial T} \right)_\rho. \quad (3.33)$$

This result is *completely general* for any thermodynamic expression relating  $(\rho, T, \varepsilon)$ .

<sup>3</sup> Unless we are creating and destroying mass!

Similarly, at constant pressure the heat change is  $dq = C_P dT$ , and the specific heat is

$$C_P = C_V + \left[ \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T - \left( \frac{P}{\rho^2} \right) \right] \left( \frac{\partial \rho}{\partial T} \right)_P. \quad (3.34)$$

For an adiabatic process, we have  $dq = 0$  and can rearrange Eq. (3.32) to find

$$\begin{aligned} \left( \frac{\partial \varepsilon}{\partial T} \right)_\rho dT + \left[ \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T - \left( \frac{P}{\rho^2} \right) \right] d\rho &= 0 \\ \left( \frac{\partial \varepsilon}{\partial T} \right)_\rho dT &= \left[ \left( \frac{P}{\rho^2} \right) - \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T \right] d\rho \\ \left( \frac{\partial \varepsilon}{\partial T} \right)_\rho \left( \frac{\partial T}{\partial \rho} \right)_s &= \left[ \left( \frac{P}{\rho^2} \right) - \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T \right] \\ C_V \left( \frac{\partial T}{\partial \rho} \right)_s &= \left[ \left( \frac{P}{\rho^2} \right) - \left( \frac{\partial \varepsilon}{\partial \rho} \right)_T \right] \end{aligned} \quad (3.35)$$

where we have used the definition of the specific heat at constant volume (Eq. 3.33) and transformed the full derivative ( $dT/d\rho$ ) into a partial derivative at constant entropy,  $s$ .

The specific heat at constant pressure, Eq. (3.34), is not nearly as simple and neat an expression as that for constant volume, Eq. (3.33). We can gain some physical insight into the specific heat at constant pressure by defining the specific *enthalpy*:

$$h = \varepsilon + PV = \varepsilon + \frac{P}{\rho}. \quad (3.36)$$

The change in enthalpy in some thermodynamic process is then

$$dh = d\varepsilon + Pd(1/\rho) + dP/\rho$$

Substituting into the First Law, Eq. (3.28), we find

$$dq = dh - \frac{dP}{\rho}.$$

Therefor, at constant pressure ( $dP = 0$ ) we recognize that we can write the specific heat by replace  $q$  with  $h$ :

$$C_P = \left( \frac{\partial h}{\partial T} \right)_P, \quad (3.37)$$

a much simpler expression than Eq. (3.34). This implies that for isobaric processes, the enthalpy takes the role of internal energy for isovolumetric processes.

Neat!

### The Second Law

Empirically, it is found that not all energy-conserving cyclic processes are possible. That is, energy in such processes is transferred into, e.g., its surroundings and is rendered unavailable to do work or heat. In a sense, the original energy of the system is *degraded* and, even if total energy is conserved, it is not possible to return the system to its original state.

To understand this behavior of gasses, we must first introduce a new thermodynamic state function, *entropy S*. If in a reversible process a system exchanges some heat  $dQ$  with a reservoir at temperature  $T$ , the change in entropy is

$$dS \equiv \frac{dQ}{T}. \quad (3.38)$$

Using this definition of entropy, we can re-write the First Law (Eqs. 3.27 and 3.28) as

$$TdS = dE + PdV \quad (3.39)$$

$$Tds = de + \frac{P}{\rho^2} d\rho, \quad (3.40)$$

where the second version is the intensive form as we have entropy the *specific entropy s*.

For any cyclic process, the integral of Eq. (3.38) over an entire cycle must be less than equal to zero. That is

$$\oint \frac{dQ}{T} \leq 0. \quad (3.41)$$

This is an elementary state of the Second Law of Thermodynamics.

To elucidate the meaning of this further, consider a reversible process in which heat  $dQ_1(T)$  is delivered to the system. In order to reverse the process and return the system to its original state, heat  $dQ_2(T) \equiv -dQ_1(T)$  must be given to the system. Equation (3.38) requires that any cycle of the system satisfies both

$$\oint \frac{dQ_1(T)}{T} \leq 0$$

and

$$\oint \frac{dQ_2(T)}{T} = - \oint \frac{dQ_1(T)}{T} \leq 0.$$

In order for both conditions to be simultaneously true, the heat transfer must be zero so that for a reversible process

$$\oint_{\text{rev}} \frac{dQ}{T} = 0. \quad (3.42)$$

This then implies that  $dS = 0$  by Eq. (3.38). I.e., in a reversible process, the entropy remains constant.

The process is irreversible.

It will turn out that Eq. (3.38) is an *exact* differential, i.e.  $\int dS$  is path-independent. In other words, the entropy difference between two states is independent of the details of the reversible process connecting them. If  $A$  and  $B$  are two states connected by a reversible cycle made up of path 1 ( $A$  to  $B$ ) and path 2 ( $B$  to  $A$ ) then Eq. (3.42) implies

$$\oint \frac{dQ}{T} = \left( \int_A^B \frac{dQ}{T} \right)_1 + \left( \int_B^A \frac{dQ}{T} \right)_2 = 0. \quad (3.43)$$

Now that we've specified some limits we can write these integrals in terms of the entropies of the states via Eq. (3.38)

$$\left( \int_A^B \frac{dQ}{T} \right)_1 = S(B) - S(A),$$

which gives us

$$S(B) - S(A) = - \left( \int_B^A \frac{dQ}{T} \right)_2 = \left( \int_A^B \frac{dQ}{T} \right)_2, \quad (3.44)$$

which proves that the entropy difference of the states doesn't depend on the details of the processes.

Now let's consider that an *irreversible* process  $I$  connects  $A$  to  $B$  then reversible process  $R$  returns  $B$  to  $A$ . Now via Eq. (3.41)

$$\int_I \frac{dQ}{T} + \int_R \frac{dQ}{T} \leq 0, \quad (3.45)$$

or

$$\int_I \frac{dQ}{T} \leq \left( - \int_B^A \frac{dQ}{T} \right)_R = \left( \int_A^B \frac{dQ}{T} \right)_R = S(B) - S(A). \quad (3.46)$$

This implies that for an isolated system (i.e.,  $dQ = 0$ ) entropy can only increase:

$$S(B) \geq S(A)$$

where the equality holds for reversible processes only. All processes in the real world are actually irreversible, at least to some extent. Thus, the entropy of any real system always increases, even if that system is perfectly thermally insulated ( $dQ = 0$ ).

### Some Thermal Properties of Ideal Gasses

We often approximate stellar plasmas as ideal gasses, fraught as that may be. As we argued above, the internal energy is a function of temperature only. So using the definition of the specific heat (Eq. 3.34) we have

$$\varepsilon = \int C_V dT. \quad (3.47)$$

While the proof will require some more development of kinetic theory, we will find that  $C_V$  is a constant for an ideal gas, thus

$$\varepsilon = C_V T. \quad (3.48)$$

An exact differential is one that can be express  $df = (\partial f / \partial x)_y dx + (\partial f / \partial y)_x dy$ .

### 3.5 The equations in Lagrangian form

The fluid equations (3.1), (3.3), and (3.5) are in **Eulerian** form; that is, they describe everything in terms of spatial coordinates and time. This is not necessarily the most convenient form for practical calculations. For example, the star can expand and contract, making the radius a function of time. Moreover, the velocity  $u$  is *not* the velocity of a given fluid element, which is why the equation of motion (eq. [2.5]) is non-linear. It is often desirable to put the fluid equations into **Lagrangian** form, in which the coordinates are some label for a fluid element and time.

In one-dimension, the transformation to Lagrangian equations is easy. At some reference time, we label the mass enclosed by a shell of radius  $r$

$$m(r, t) = \int_0^r \rho(r', t) 4\pi r'^2 dr', \quad (3.49)$$

as a Lagrangian coordinate  $m$ ; we then transform coordinates from  $(r, t)$  to  $(m, t)$ . To do this, differentiate eq. (3.49) w.r.t.  $r$ ,

$$\partial_r m = 4\pi r^2 \rho,$$

and substitute for  $\rho$  in the equation of continuity (eq. [3.1]). The first term becomes

$$\partial_t \rho = \partial_t \left( \frac{1}{4\pi r^2} \partial_r m \right) = \frac{1}{4\pi r^2} \partial_r (\partial_t m),$$

while the second term becomes

$$\frac{1}{4\pi r^2} \partial_r (u \partial_r m);$$

the equation of continuity therefore becomes

$$\frac{1}{4\pi r^2} \partial_r (\partial_t m + u \partial_r m) = 0. \quad (3.50)$$

We can integrate this over  $r$  to find that  $\partial_t m + u \partial_r m = f(t)$ , where  $f(t)$  is some as-yet-unspecified function; to fix  $f(t)$ , we note that since  $m(0, t) = 0, \forall t$ , we must have  $f(t) = 0$ . Now  $\partial_t m + u \partial_r m = Dm/Dt = 0$ , so along a streamline,  $m$  is a constant. We can therefore transform from coordinates  $(r, t)$  to  $(m, t)$  by setting

$$\left. \frac{\partial}{\partial t} \right|_r + u \left. \frac{\partial}{\partial r} \right|_t = \left. \frac{\partial}{\partial t} \right|_m \equiv \frac{D}{Dt} \quad (3.51)$$

$$\left. \frac{\partial}{\partial r} \right|_t = 4\pi r^2 \rho \left. \frac{\partial}{\partial m} \right|_t. \quad (3.52)$$

Here  $D/Dt \equiv (\partial/\partial t)_m$  is the Lagrangian time derivative. In deriving this change, we used the equation of continuity, which becomes

$$\frac{\partial r}{\partial m} = \frac{1}{4\pi r^2 \rho}. \quad (3.53)$$

Our equation for momentum (eq. [3.2]) becomes

$$\frac{\partial P}{\partial m} = -\frac{Gm}{4\pi r^4} - \frac{1}{4\pi r^2} \frac{Du}{Dt}. \quad (3.54)$$

In hydrostatic balance the second term on the right-hand side is negligible. The flux equation, (eq. [3.7]) can be transformed to

$$\frac{\partial T}{\partial m} = -\frac{1}{16\pi^2 r^4 \rho K} L_r \quad (3.55)$$

Here  $L_r$  is the luminous flux at a radius  $r$ .

The energy equation (eq. [3.5]) is more complicated. We can expand the time derivative as

$$\begin{aligned} \partial_t \left( \frac{1}{2} \rho u^2 + \rho \varepsilon \right) &= \left( \frac{1}{2} u^2 + \varepsilon \right) \partial_t \rho + \rho \partial_t \left[ \frac{1}{2} (\mathbf{u} \cdot \mathbf{u}) + \varepsilon \right] \\ &= - \left( \frac{1}{2} u^2 + \varepsilon \right) \nabla \cdot (\rho \mathbf{u}) + \rho \mathbf{u} \cdot \nabla \partial_t \mathbf{u} + \rho \partial_t \varepsilon, \end{aligned}$$

using equation (3.1) to substitute for  $\partial_t \rho$ . We then use equation (3.2) to replace  $\partial_t \mathbf{u}$ , and recognizing that  $\mathbf{u}(\mathbf{u} \cdot \nabla) \mathbf{u} = \mathbf{u} \cdot \nabla[(1/2)u^2]$ , rewrite equation (3.5) as

$$\rho (\partial_t + \mathbf{u} \cdot \nabla) \varepsilon + P \nabla \cdot \mathbf{u} = -\nabla \cdot \mathbf{F} + \rho q. \quad (3.56)$$

We've canceled all common factors here. Finally, we once again use equation (3.1) to set

$$P \nabla \cdot \mathbf{u} = -(P/\rho)(\partial_t \rho + \mathbf{u} \cdot \nabla \rho) = \rho P (\partial_t + \mathbf{u} \cdot \nabla) \left( \frac{1}{\rho} \right).$$

Substituting this into the left-hand of equation (3.56) and using the first law of thermodynamics (see eq. [3.28]), we obtain

$$\rho (\partial_t + \mathbf{u} \cdot \nabla) \varepsilon + P \nabla \cdot \mathbf{u} = \rho T (\partial_t + \mathbf{u} \cdot \nabla) s. \quad (3.57)$$

For the right-hand side of equation (3.56), we expand the divergence operator in spherical symmetry and use equation (3.52) to obtain

$$-\nabla \cdot \mathbf{F} = -\frac{1}{r^2} \frac{\partial(r^2 F)}{\partial r} = -\rho \frac{\partial L_r}{\partial m}.$$

Putting everything together, we finally have our heat equation in Lagrangian form,

$$\frac{\partial L_r}{\partial m} = q - T \frac{Ds}{Dt}. \quad (3.58)$$

This has a simple interpretation: the change in luminosity across a mass shell is due to sources or sinks of energy and the change in the heat content of the shell.

---

EXERCISE 3.3— Show that equation (3.58) can be written as

$$\frac{\partial L_r}{\partial m} = q - \frac{c_p T}{\chi_T} \left\{ \frac{D \ln P}{Dt} - [\chi_\rho + \chi_T (\Gamma_3 - 1)] \frac{D \ln \rho}{Dt} \right\} \quad (3.59)$$

$$= q - \frac{P}{\rho(\Gamma_3 - 1)} \frac{D}{Dt} \ln \left( \frac{P}{\rho^{\Gamma_1}} \right), \quad (3.60)$$

where

$$\begin{aligned} \chi_T &\equiv \frac{T}{P} \left( \frac{\partial P}{\partial T} \right)_\rho, \\ \chi_\rho &\equiv \frac{\rho}{P} \left( \frac{\partial P}{\partial \rho} \right)_T, \\ \Gamma_1 &\equiv (\partial \ln P / \partial \ln \rho)_s, \text{ and} \\ \Gamma_3 - 1 &\equiv \left( \frac{\partial \ln T}{\partial \ln \rho} \right)_s \end{aligned}$$

are defined in Appendix A.1.

---

It is more useful, however, to work with temperature and pressure instead of entropy. Write

$$T \frac{Ds}{Dt} = T \left( \frac{\partial s}{\partial T} \right)_P \frac{DT}{Dt} + T \left( \frac{\partial s}{\partial P} \right)_T \frac{DP}{Dt},$$

and use the identity (see Appendix A.1)

$$\left( \frac{\partial s}{\partial P} \right)_T = - \left( \frac{\partial s}{\partial T} \right)_P \left( \frac{\partial T}{\partial P} \right)_s$$

to obtain

$$\frac{\partial L_r}{\partial m} = q - c_P \left[ \frac{DT}{Dt} - \left( \frac{\partial T}{\partial P} \right)_s \frac{DP}{Dt} \right]. \quad (3.61)$$

Equations (3.53), (3.54), (3.55), and (3.61), when supplemented by an equation of state, a prescription for the thermal conductivity, and the equations for nuclear heating and neutrino cooling, form the equations for stellar structure and evolution in spherical symmetry.



# 4

## *Polytropes and the Lane-Emden Equation*

In the previous chapters, we derived the basic equations of stellar structure. To construct detailed models is clearly an involved task because there are many linked physical variables. For example, the equation of hydrostatic balance, eq. (2.7), contains both  $\rho$  and  $P$ ; in order to connect these two variables we have to know the temperature, so in general we need to include an equation for heat transport and heat flux. Before diving into all that, however, we'll take a little digression to solve some simplified stellar models known as **polytropes**; these are useful not only for historical reasons, but they also allow for quick analytical calculations.

### *4.1 Historical Background*

To understand where the term polytropic comes from, let's first consider an ideal gas in hydrostatic equilibrium, and furthermore suppose that the temperature and density lie along an adiabat. In that case we have the following relations:

$$T\rho^{1-\gamma} = \text{const}; \quad P\rho^{-\gamma} = \text{const}; \quad TP^{(1-\gamma)/\gamma} = \text{const}, \quad (4.1)$$

which you will recall from elementary thermodynamics. Here  $\gamma = C_P/C_\rho$  is the ratio of specific heats. The equation of state is

$$P = \left( \frac{N_A k_B}{\mu} \right) \rho T, \quad (4.2)$$

where  $\mu$  is the mean molecular weight and the quantity in parenthesis is  $C_P - C_\rho = N_A k_B / \mu$ . We might imagine, for example, a rising plume of hot air in Earth's atmosphere. There is one snag with this analysis for the troposphere, however: the condensation of water vapor means that one cannot hold  $dS = 0$  in a rising plume of hot, moist air. Attempting to model this moist convection in the Earth's troposphere motivated work in the early 1900's by Kelvin, Lane, Emden, and others to consider a more general problem, in

which  $TdS = CdT$ , where  $C$  is a constant. A configuration for which this is true is called **polytropic**. An adiabat is a special case of a polytrope with  $C = 0$ . Writing the first law of thermodynamics as  $TdS = C_p dT - (P/\rho^2)d\rho$ , substituting for  $TdS$ , and using equation (4.2), we obtain

$$(C_p - C) \frac{dT}{T} = (C_p - C_\rho) \frac{d\rho}{\rho}.$$

This equation has a solution  $T \propto \rho^{(C_p - C_\rho)/(C_p - C)}$ . Comparing this solution with equation (4.1), we can define a *polytropic exponent*,  $\gamma' = (C_p - C)/(C_\rho - C)$ . Then equation (4.1) holds with  $\gamma$  replaced by  $\gamma'$ . The advantage of this approximation is that it relates density to pressure so that one can solve the equation of hydrostatic equilibrium without simultaneously having to solve for  $T(r)$ .

## 4.2 The Lane-Emden Equation and Solution

To use the polytropic equation of state, write the pressure  $P$  as

$$P(r) = K\rho^{1+1/n}(r) \quad (4.3)$$

where  $n$  and  $K$  are constants. Further define the dimensionless variable  $\theta$  via

$$\rho(r) = \rho_c \theta^n(r), \quad (4.4)$$

where the subscript  $c$  denotes the central value at  $r = 0$ . Note that since

$$P(r) \propto \rho \times \rho^{1/n} \propto \rho\theta,$$

the quantity  $\theta$  plays the role of a dimensionless temperature for an ideal non-degenerate gas.

Substitute equations (4.3) and (4.4) into Poisson's equation,

$$\nabla^2 \Phi = 4\pi G\rho, \quad (4.5)$$

and the equation for hydrostatic equilibrium,

$$\nabla P = -\rho \nabla \Phi, \quad (4.6)$$

to obtain the *Lane-Emden* equation for index  $n$ ,

$$\xi^{-2} \frac{d}{d\xi} \left( \xi^2 \frac{d\theta}{d\xi} \right) = -\theta^n. \quad (4.7)$$

Here  $\xi = r/r_n$  is the dimensionless coordinate, and

$$r_n = \left[ \frac{(n+1)P_c}{4\pi G\rho_c^2} \right]^{1/2} \quad (4.8)$$

is the radial length scale.

For a stellar model described by a single polytropic relation, the appropriate boundary conditions are

$$\theta(\xi)|_{\xi=0} = 1, \quad (4.9)$$

$$\theta'(\xi)|_{\xi=0} = 0. \quad (4.10)$$

From the form of equation (4.7), it follows that  $\theta(-\xi) = \theta(\xi)$ , that is, the solution is *even* in  $\xi$ . A power-series solution of  $\theta$  out to order  $\xi^6$  is

$$\theta(\xi) = 1 - \frac{1}{6}\xi^2 + \frac{n}{120}\xi^4 - \frac{n(8n-5)}{15120}\xi^6 + \mathcal{O}(\xi^8) \quad (4.11)$$

There are analytical solutions for  $n = 0, 1$ , and  $5$ :

$$\theta_0(\xi) = 1 - \frac{\xi^2}{6} \quad (4.12)$$

$$\theta_1(\xi) = \frac{\sin \xi}{\xi} \quad (4.13)$$

$$\theta_5(\xi) = \left( \frac{3}{3 + \xi^2} \right)^{1/2}. \quad (4.14)$$

Finally, the radius of the stellar model is determined by the location of the first zero,  $\xi_1$ , where  $\theta(\xi_1) = 0$ . For example, if  $n = 0$  (eq. [4.12]),  $\xi_1 = \sqrt{6}$ . Note that if  $n = 5$  there is no root;  $\theta_5(\xi) > 0, \forall \xi > 0$ . A sample of Lane-Emden solutions for various indices is shown in Figure 4.1.

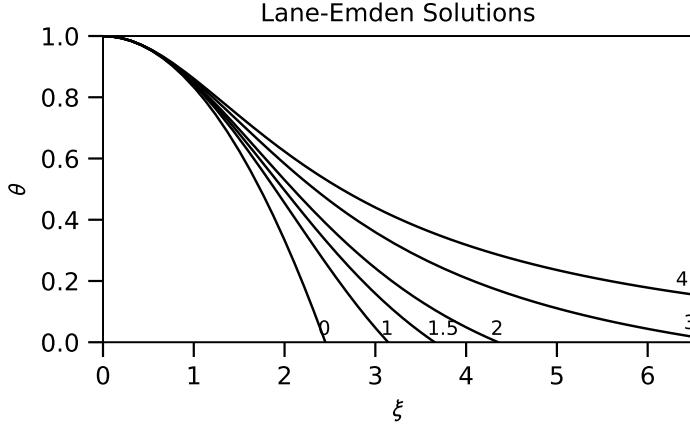


Figure 4.1: Solutions of the Lane-Emden equation for selected values of the index  $n$ .

### 4.3 Some Useful Relations

First, let's get the mass of our polytropic sphere. To do this, we write the integral

$$M = \int_0^R 4\pi r^2 \rho dr$$

and make the substitutions  $r = r_n \xi$ ,  $R = r_n \xi_1$ , and  $\rho = \rho_c \theta^n(\xi)$  to obtain

$$M = 4\pi r_n^3 \rho_c \int_0^{\xi_1} \xi^2 \theta^n(\xi) d\xi.$$

Using equation (4.7), the integrand can be written as a perfect differential, so we get

$$M = 4\pi r_n^3 \rho_c \left( -\xi_1^2 \theta'_1 \right). \quad (4.15)$$

Here I define the shorthand  $\theta'_1 \equiv [d\theta(\xi)/d\xi]_{\xi=\xi_1}$ . Substituting  $r_n = R/\xi_1$  and dividing by 3 allows us to get a formula relating the central density to the mean density,

$$\rho_c = \frac{3M}{4\pi R^3} \left( -\frac{\xi_1}{3\theta'_1} \right). \quad (4.16)$$

For the solutions shown in Figure 4.1, we have the following values of  $\rho_c/\bar{\rho}$ , as shown in Table 16.1. As the index  $n$  increases, the configuration becomes more and more concentrated toward the center.

$n$	0	1.0	1.5	2.0	3.0	4.0
$\xi_1$	2.449	3.142	3.654	4.353	6.897	14.972
$-\theta'_1$	0.8165	0.3183	0.2033	0.1272	0.04243	0.008018
$\rho_c/\bar{\rho}$	1.00	3.29	5.99	11.41	54.18	622.4

Table 4.1: Properties of the Lane-Emden solutions.

Starting from equation (4.15), we can substitute for  $r_n$  using equation (4.8) and  $\rho_c$  using equation (4.16) to get an equation for the central pressure,

$$P_c = \frac{GM^2}{R^4} \frac{1}{4\pi(n+1)(-\theta'_1)^2}. \quad (4.17)$$

For an ideal gas,  $P_c = (N_A k_B / \mu) \rho_c T_c$  with  $\mu$  being the mean molecular weight, we can solve for the central temperature,

$$T_c = \left( \frac{\mu}{N_A k_B} \right) \left( \frac{GM}{R} \right) \frac{1}{(n+1)\xi_1(-\theta'_1)}. \quad (4.18)$$

Finally, starting from equation (4.17), substituting for  $P_c$  using equation (4.3), and eliminating  $\rho_c$  using equation (4.16), we obtain a relation between mass and radius in terms of  $K$  and  $n$ ,

$$M^{1-1/n} = \left[ \frac{K(n+1)}{G(4\pi)^{1/n}} \xi_1^{1+1/n} (-\theta'_1)^{1-1/n} \right] R^{1-3/n}. \quad (4.19)$$

Alternatively, one could use this equation to fit  $K$  to a star of known  $M$  and  $R$ .

---

EXERCISE 4.1 — Derive equations (4.15)–(4.19). Explain what the mass-radius relation, eq. (4.19), means for the cases  $n = 1$  and  $n = 3$ .

---

Finally, we can derive a formula for the gravitational energy of our polytropic sphere. First, we can integrate the equation for the energy,

$$E_{\text{grav}} = -G \int_0^M \frac{m}{r} dm,$$

by parts to obtain

$$E_{\text{grav}} = -\frac{GM^2}{2R} - \frac{1}{2} \int_0^R \frac{Gm^2}{r^2} dr = -\frac{GM^2}{2R} - \frac{1}{2} \int_0^R \frac{d\Phi}{dr} m dr. \quad (4.20)$$

If we define our zero of energy to be such that  $\Phi(R) = 0$ , then we can integrate by parts again to obtain

$$E_{\text{grav}} = -\frac{GM^2}{2R} + \frac{1}{2} \int_0^R \Phi dm. \quad (4.21)$$

We can rewrite the equation of hydrostatic equilibrium as

$$\frac{dP}{dr} = -\frac{d\Phi}{dr}\rho$$

and use equation (4.3) to eliminate  $P$  to obtain

$$\frac{d\Phi}{dr} = (1+n)K\rho^{1/n}.$$

Integrating from a point in the star  $r$  to  $R$ , and again using choosing  $\Phi(R) = 0$ , we obtain

$$\Phi(r) = -(1+n)K\rho(r)^{1/n} = -(1+n)\frac{P(r)}{\rho(r)}. \quad (4.22)$$

Inserting equation (4.22) into equation (4.21), we have

$$E_{\text{grav}} = -\frac{GM^2}{2R} - \frac{1+n}{2} \int_0^M \frac{P}{\rho} dm = -\frac{GM}{2R} + \frac{1+n}{6} E_{\text{grav}}, \quad (4.23)$$

where we used equation (2.13) to relate the integral of  $P/\rho$  to  $E_{\text{grav}}$ .

Solving equation (4.23) for  $E_{\text{grav}}$  gives us the desired result,

$$E_{\text{grav}} = -\frac{3}{5-n} \frac{GM^2}{R}. \quad (4.24)$$

Note that solutions with  $n > 5$  have a positive gravitational energy.

---

**EXERCISE 4.2—** For a fully convective star with  $S(r) = \text{const}$  and an ideal gas equation of state, how is the polytropic constant  $K$  related to the entropy? Derive a formula for  $R$  in terms of  $M$  and  $s$  in this case. What happens to the star if the entropy increases, i.e., heat is added to it?

*Hint:* Recall from thermodynamics that the entropy per unit mass of an ideal gas is

$$s = \frac{k_B N_A}{\mu} \left\{ \frac{5}{2} + \ln \left[ \frac{\mu}{\rho N_A} \left( \frac{\mu k_B T}{2\pi N_A \hbar^2} \right)^{3/2} \right] \right\}. \quad (4.25)$$

Use the Lane-Emden solution to compute the specific entropy, per unit mass, in terms of the central temperature  $T$  and the stellar mass  $M$ :  $s = s(T_c, M)$ . From this expression, compute the “gravothermal” specific heat

$$c_* = T_c \frac{\partial s(T_c, M)}{\partial T_c}, \quad (4.26)$$

and comment on its physical significance.

---

### **MESA Entropy, radius, and gravothermal specific heat**

Download the folder `Lane-Emden/0.3M-entropy` and place it into your projects folder. This project evolves a  $0.3 M_\odot$  through its pre-main-sequence phase and the first 3 Gyr of its main-sequence life. A star with initial mass  $M \lesssim 0.35 M_\odot$  remains fully convective, so that the entire star lies along an adiabat and has a single value of entropy  $s$  at a given instant. This star therefore makes a good test case for checking your work in the first warm-up exercise.

You now have analytical expressions for how the radius scales with entropy, and for the “gravothermal” specific heat  $c_*$ . As you evolve the low-mass star through its pre-main-sequence phase, its radius and entropy change. Make a plot that demonstrates the scaling of radius with entropy. Make a second plot that shows the scaling of entropy with central temperature  $T_c$  and from this plot obtain an estimate of  $c_*$ . Compare your findings against the results of the first warm-up exercise.

**NB.** It is unnecessary to modify ‘`run_star_extras.f`’ for this project. Although **MESA** computes the central entropy, it does not print it to the ‘`history.data`’ file by default. In the ‘`&star_job`’ section of ‘`project.inlist`’ we therefore set

25        `history_columns_file = 'entropy_variables'`

and make the file ‘`entropy_variables`’:

```
1        include ''
2        center_entropy ! in units of kerg per baryon
```

This gives us a history of  $s/(N_A k_B)$  as the star evolves. You will note that in the top-level inlist, there is the following reference:

```
26     read_extra_pgstar_inlist2 = .false.  
27     extra_pgstar_inlist2_name = 'track_entropy.inlist'
```

You will need to create the file ‘track\_entropy.inlist’ to make these plots. I recommend using a custom version of the ‘History\_Track’ plots, a complete description of which you can find in ‘\$MESA\_DIR/star/defaults/pgstar.defaults’.



# 5

## *Kinetic Theory for Stars*

Kinetic theory deals with collections of particles that evolve by undergoing collisions. We will find that for the study of stellar physics, kinetic theory comes up in several aspects including fluid dynamics, radiative transport, equations of state, and nuclear reactions. Therefore, in this chapter we will develop some of the fundamentals of kinetic theory and some key relations. Much of the development in this chapter follows that in Mihalas and Mihalas [1984] and we refer interested readers to that work for a more thorough discussion of some of these, and related, topics.

### *Distribution Functions and the Boltzmann Equation*

The physical state of a gas can be described statistically by a *distribution function*  $f(\mathbf{x}, \mathbf{v}, t)$  which describes the occupation of *phase space* element  $d^3x d^3v$ , where  $\mathbf{x}$  and  $\mathbf{v}$  the position and velocity vectors.

Distribution function

The average number of particles contained in the volume  $d^3x$  at  $\mathbf{x}$  and in the velocity-space element  $d^3v$  about the velocity  $\mathbf{v}$  is  $f d^3x d^3v$ . The distribution function cannot be negative and we assume that as the velocity goes to infinity (in either the positive or negative direction) that  $f \rightarrow 0$  sufficiently fast that there is a finite number of particles and a finite total energy.

Note that a more general phase space is  $d^3x d^3p$ , where  $\mathbf{p}$  is the momentum vector. Working in momentum space allows for particles of differing masses, or even massless particles such as photons. We will switch to momentum space eventually, but for now it is simpler to assume identical massive particles.

Several key macroscopic properties of a gas can be computed from the distribution function. The particle number density is

$$n(\mathbf{x}, t) = \frac{dN}{d^3x} = \int_{-\infty}^{\infty} f(\mathbf{x}, \mathbf{v}, t) d^3v, \quad (5.1)$$

from which we can easily compute the mass density,

$$\rho(\mathbf{x}, t) = mn(\mathbf{x}, t), \quad (5.2)$$

where  $m$  is some appropriate (mean) particle mass. The average velocity at  $x$  can be computed as

$$\mathbf{u}(\mathbf{x}, t) = \langle \mathbf{v} \rangle \equiv n^{-1} \int_{-\infty}^{\infty} f(\mathbf{x}, \mathbf{v}, t) \mathbf{v} d^3 v. \quad (5.3)$$

Indeed, the average value of any quantity  $Q^1$  at some location  $x$  can be calculated using the distribution function:

$$\langle Q \rangle = \frac{\int Q f d^3 v}{\int f d^3 v} = n^{-1} \int Q f d^3 v. \quad (5.4)$$

In considering the microscopic properties of a gas it will often be convenient to decompose a given particles velocity into *bulk* and *random* parts:

$$\mathbf{v} = \mathbf{u} + \mathbf{U}, \quad (5.5)$$

where  $\mathbf{u}$  is the bulk average velocity at  $x$  as given in Eq. (5.3) and we have introduced the *random velocity*  $\mathbf{U}$  relative to the mean flow. By definition, the mean of the random velocity is zero:  $\langle \mathbf{U} \rangle \equiv 0$ .

Similarly we can obtain the rate at which particles are crossing some unit surface of the  $y - z$  plane (i.e., the particle flux in the  $x$ -direction) by

$$S_x = \frac{dN}{dydzdt} = \int \frac{dN}{dxdydzd^3v} \frac{dx}{dt} d^3v = \int f(\mathbf{x}, \mathbf{v}, t) v_x d^3v. \quad (5.6)$$

Generalizing to all directions, the particle flux can be computed as

$$\mathbf{S} = \int f(\mathbf{x}, \mathbf{v}, t) \mathbf{v} d^3v = n \mathbf{u}, \quad (5.7)$$

which we recognize as just the mean velocity times the number density. Multiplying this by the mean particle mass again, we obtain the momentum density vector,

$$\mathcal{P} = m \mathbf{S} = \rho \mathbf{u} = \int m \mathbf{v} f(\mathbf{x}, \mathbf{v}, t) d^3v. \quad (5.8)$$

Define a *stress tensor*  $T_{jk}$ , the meaning of which is the  $j$  component of the force per unit area across a surface perpendicular to  $\hat{k}$ . This has the simultaneous meaning of the  $j$  component of the momentum that crosses a unit area perpendicular to  $\hat{k}$ .<sup>2</sup> Now armed with this definition we can write the rate at which the  $j$  component of the momentum crossing the  $y - z$  plane as

$$\begin{aligned} T_{jx} &= \int \frac{dN}{dydzdt d^3v} m v_j d^3v = \int \frac{dN}{dxdydzd^3v} \frac{dx}{dt} m v_j d^3v \\ &= \int m v_j v_x f d^3v. \end{aligned} \quad (5.9)$$

Generalizing this to all directions and components, we find

$$T_{jk} = \int m v_j v_k f d^3v. \quad (5.10)$$

<sup>1</sup> Not to be confused with the extensive heat from 3.4.

<sup>2</sup> See Section 1.9 of Thorne and Blandford [2017].

Comparing Eqs. (5.2), (5.8), and (5.10) we can see that these physical quantities of the fluid actually arise from a hierarchy of *moments* in velocity space of the distribution function. A moment of the distribution function is formed by multiplying the mass-weighted distribution function  $mf$  by the velocity  $M^{\text{th}}$  times then integrating over all velocity space. Or,

$$\text{Moment}(f, M) = m \int (v_i)^M f d^3 v. \quad (5.11)$$

Let's now consider how the distribution function evolves in time. We will at first ignore particle collisions. Consider a group of particles at the phase-space location  $(x_0, y_0, z_0, u_0, v_0, w_0)$ <sup>3</sup> in the phase-space volume  $(dx_0, dy_0, dz_0, du_0, dv_0, dw_0)$ . If an external force  $F(x, t)$  acts on the particles, the acceleration is  $a(x, t) = m^{-1}F(x, t)$ , where  $m$  is the particle mass.<sup>4</sup> After some time  $dt$  the particles will evolve into a new phase-space element centered on position  $x = x_0 + v_0 dt$  and velocity  $v = v_0 + adt$  with volume  $(dx, dy, dz, du, dv, dw)$ . We can relate the new volume to the original volume via a Jacobian  $J$ :

$$d^3 x d^3 v = J(x, y, z, u, v, w / x_0, y_0, z_0, u_0, v_0, w_0) d^3 x_0 d^3 v_0. \quad (5.12)$$

To first order in  $dt$ , this Jacobian is one:  $J = 1 + \mathcal{O}(dt^2)$ . So for sufficiently small  $dt$  the volume is unchanged,

$$d^3 x d^3 v = d^3 x_0 d^3 v_0.$$

The total number of particles initially is  $\delta N_0 = f(x_0, v_0, t_0) d^3 x_0 d^3 v_0$  and, in the absence of any collisions, all the same particles that started in  $d^3 x_0 d^3 v_0$  will end up in  $d^3 x d^3 v$ , so

$$\delta N = f(x = x_0 + v_0 dt, v = v_0 + adt, t_0 + dt) d^3 x d^3 v = \delta N_0. \quad (5.13)$$

Now since both the number of particles and the phase space volume are the same, the phase-space density of the group of particles is unchanged. I.e., the distribution function remains the same

$$f(x = x_0 + v_0 dt, v = v_0 + adt, t_0 + dt) = f(x_0, v_0, t_0). \quad (5.14)$$

Expanding Eq. (5.14) to first order in  $dt$  we find

$$\frac{\partial f}{\partial t} + v^i \frac{\partial f}{\partial x^i} + a^i \frac{\partial f}{\partial v^i} = 0, \quad (5.15)$$

which is the *collisionless Boltzmann equation*.

<sup>3</sup> Where  $(u, v, w)$  represent the components of the 3-velocity,  $u$  not to be confused with the mean velocity.

<sup>4</sup> Again, all particles are, for now, assumed to have the same mass  $m$ .

in and out of the phase-space element  $d^3x d^3v$ . We can represent this as a source term on the RHS of Eq. (5.15) resulting in the *Boltzmann transport equation*:

Boltzmann transport equation

$$\frac{\partial f}{\partial t} + v^i \frac{\partial f}{\partial x^i} + a^i \frac{\partial f}{\partial v^i} = \left( \frac{Df}{Dt} \right)_{\text{coll}}. \quad (5.16)$$

The collision source term is expressed as a Lagrangian derivative because here we are following the evolution of some particular group of particles. The exact nature of the collision source term can be complex and depends sensitively on the microphysical nature of the gas under consideration.

### Fluid Equations from the Boltzmann Equation

An interesting and useful result of kinetic theory is that we can derive the fluid equations (Eqs. 3.1, 3.3, & 3.5) from the Boltzmann transport equation, Eq. 5.16, connecting the macroscopic directly to the microscopic. We begin with taking a *moment* of the Boltzmann equation by multiplying by some physical quantity  $Q(x, v)$  and then integrating over velocity space:

$$\int Q \left( \frac{\partial f}{\partial t} + v^i \frac{\partial f}{\partial x^i} + a^i \frac{\partial f}{\partial v^i} \right) d^3v = \int Q \left( \frac{Df}{Dt} \right)_{\text{coll}} d^3v \equiv I(Q), \quad (5.17)$$

where we have introduced the collision integral  $I(Q)$ . Consider now a collision between two particles initial with velocities  $v$  and  $v_1$ . The collision will change the particle velocities:  $(v, v_1) \rightarrow (v', v'_1)$ . If  $Q(x, v)$  is a conserved quantity during such a collision, then

$$Q(x, v) + Q(x, v_1) = Q(x, v') + Q(x, v'_1), \quad (5.18)$$

and the collision integral  $I(Q) = 0$ . This can be understood simply as the collisions at  $x$  cannot create or destroy  $Q$  only change its distribution in velocity space. Then, if we are integrating over all velocity space the integral must be zero.

Let's consider each term on the LHS of Eq. (5.17) separately. The first term is

$$\begin{aligned} \int Q \frac{\partial f}{\partial t} d^3v &= \partial_t \int Q f d^3v - \int \frac{\partial Q}{\partial t} f d^3v \\ &= \partial_t (n \langle Q \rangle) - n \left\langle \frac{\partial Q}{\partial t} \right\rangle, \end{aligned} \quad (5.19)$$

where we have used the definition of averages, Eq. (5.4). The second term is

$$\begin{aligned} \int Q v^i \frac{\partial f}{\partial x^i} d^3v &= \frac{\partial}{\partial x^i} \int Q v^i f d^3v - \int v^i \frac{\partial Q}{\partial x^i} f d^3v \\ &= \frac{\partial}{\partial x^i} (n \langle Q v^i \rangle) - n \left\langle v^i \frac{\partial Q}{\partial x^i} \right\rangle. \end{aligned} \quad (5.20)$$

Finally, the third term on the LHS of Eq. (5.17) becomes

$$\begin{aligned}\int Q a^i \frac{\partial f}{\partial v^i} d^3 v &= \int \frac{\partial(Q a^i f)}{\partial v^i} d^3 v - \int \frac{\partial(Q a^i)}{\partial v^i} f d^3 v \\ &= \sum_i \int_{-\infty}^{\infty} dv_l dv_k |_{-\infty}^{\infty} (Q a^i f) - n \left\langle \frac{\partial(Q a^i)}{\partial v^i} \right\rangle \\ &= -n \left\langle \frac{\partial(Q a^i)}{\partial v^i} \right\rangle,\end{aligned}\quad (5.21)$$

where we recognize that as the first integral is now even and that as  $v^i \rightarrow \pm\infty$ ,  $f \rightarrow 0$  rapidly meaning that  $(Q a^i f)$  must vanish.

Combining all the terms Eqs. (5.19) to (5.21) again, we arrive at a completely general conservation theorem:

$$\frac{\partial}{\partial t}(n\langle Q \rangle) + \frac{\partial}{\partial x^i}(n\langle Q v^i \rangle) - n \left[ \left\langle \frac{\partial Q}{\partial t} \right\rangle + \left\langle v^i \frac{\partial Q}{\partial x^i} \right\rangle + \left\langle \frac{\partial(Q a^i)}{\partial v^i} \right\rangle \right] = 0,\quad (5.22)$$

for any conserved physical quantity  $Q$ . We can simplify this expression if we assume that all external forces are velocity-independent and that  $Q$  is a function of  $v$  only (and not  $x$  and  $t$ ) then our conservation law is

$$\frac{\partial}{\partial t}(n\langle Q \rangle) + \frac{\partial}{\partial x^i}(n\langle Q v^i \rangle) - n a^i \left\langle \frac{\partial Q}{\partial v^i} \right\rangle = 0.\quad (5.23)$$

Now for a gas without internal degrees of freedom in which to store energy, the physical quantities conserved in collisions are mass, momentum, and energy. And so we can construct a vector of conserved quantities

$$\mathbf{Q} = [m, mv_1, mv_2, mv_3, \frac{1}{2}mv^2]^T.\quad (5.24)$$

Taking each of these quantities in turn in Eq. (5.23) yields our usual equations of fluid dynamics.

THE CONTINUITY EQUATION is somewhat obvious when  $Q = m$  in Eq. (5.23). This gives

$$\frac{\partial}{\partial t}(nm) + \frac{\partial}{\partial x^i}(nm\langle v^i \rangle) = 0.$$

The mass density is just  $\rho = nm$  and the mean, or bulk, velocity is  $\langle v^i \rangle = u^i$  (Eq. 5.3), so we have

$$\rho_{,t} + (\rho u^i)_{,i} = 0,\quad (5.25)$$

in indicial notation. Recall that repeated indices imply summation<sup>5</sup> and so we recognize the second term on the LHS as the divergence of the mass flux  $\rho u^i$ . Therefore, back in vector notation,

$$\partial_t \rho + \nabla \cdot (\rho \mathbf{u}) = 0,$$

which is precisely the continuity equation represent conservation of mass (Eq. 3.1).

<sup>5</sup> Thanks, Einstein.

THE MOMENTUM EQUATION is a bit more involved. Setting  $\mathbf{Q} = mv^i$  in Eq. (5.23) gives

$$(nm\langle v^i \rangle)_t + (nm\langle v^i v^j \rangle)_{,j} - nma^j \delta_j^i = 0, \quad (5.26)$$

where we have used the Kroenecker delta to ensure that the external force term is only applied when  $i = j$ . Using the decomposition of the velocity into mean and random parts  $v = u + \mathbf{U}$  (Eq. 5.5), and recalling that  $\langle \mathbf{U} \rangle = 0$ , we can rewrite the tensorial term as

$$\begin{aligned} nm\langle v^i v^j \rangle &= \rho \langle (u^i + U^i)(u^j + U^j) \rangle = \rho(u^i u^j + u^i \langle U^j \rangle + u^j \langle U^i \rangle + \langle U^i U^j \rangle) \\ &= \rho u^i u^j + \rho \langle U^i U^j \rangle. \end{aligned} \quad (5.27)$$

This is the *momentum flux density tensor*. Now define the *stress tensor* as

$$T^{ij} \equiv -\rho \langle U^i U^j \rangle, \quad (5.28)$$

meaning that *macroscopic* fluid stresses arise from momentum exchange at a *microscopic* level. With this, Eq. (5.26) is now

$$(\rho u^i)_{,t} + (\rho u^i u^j - T^{ij})_{,j} = \rho a^i, \quad (5.29)$$

making the same substitutions we made above for the continuity equation.

So what is the stress tensor? Well, by definition in fluid dynamics the *pressure* is the negative of the average of the normal stresses experienced by the fluid. Or,

$$P = -\frac{1}{3}T_{ii} = \frac{1}{3}\rho \langle U_x^2 + U_y^2 + U_z^2 \rangle = \frac{1}{3}\rho \langle U^2 \rangle, \quad (5.30)$$

where  $\langle U^2 \rangle$  is the mean translational kinetic energy of the gas and is essentially the gas temperature (see below). So, in the absence of any *viscous* stresses, we can redefine the stress tensor in terms of the more familiar pressure as  $T_{ij} = -P\delta_{ij}$ , we can rewrite Eq. (5.26) as

$$(\rho u^i)_{,t} + (\rho u^i u^j + P\delta^{ij})_{,j} = \rho a^i, \quad (5.31)$$

which is the same as Eq. (3.2) if we take  $a^i = g = -\nabla\Phi$ .

THE ENERGY EQUATION is found by setting  $\mathbf{Q} = \frac{1}{2}mv^2$  making Eq. (5.26)

$$(\frac{1}{2}nm\langle v^2 \rangle)_{,t} + (\frac{1}{2}nm\langle v^2 v^i \rangle)_{,j} = nma^i v_j. \quad (5.32)$$

The mean-squared particle velocity is

$$\begin{aligned} \langle v^2 \rangle &= \langle v_i v^i \rangle = \langle (u_i + U_i)(u^i + U^i) \rangle \\ &= u_i u^i + 2u_i \langle U^i \rangle + \langle U_i U^i \rangle = u^2 + \langle U^2 \rangle, \end{aligned} \quad (5.33)$$

and

$$\begin{aligned}\langle v^2 v^i \rangle &= \langle (u_i + U_i)(u^i + U^i)(u^j + U^j) \rangle \\ &= u_i u^i (u^i + \langle U^i \rangle) + u^i \langle U_i U^i \rangle + 2u^i u_i \langle U^i \rangle + 2u_i \langle U^i U^j \rangle + \langle U_i U^i U^j \rangle \\ &= u^i (u^2 + \langle U^2 \rangle) + 2u_i \langle U^i U^j \rangle + \langle U^2 U^j \rangle.\end{aligned}\quad (5.34)$$

Substituting these expressions in Eq (5.32) we have

$$[\rho(\frac{1}{2}\langle U^2 \rangle + \frac{1}{2}u^2)]_t + [\rho(\frac{1}{2}\langle U^2 \rangle + \frac{1}{2}u^2)u^j + \rho u_i \langle U^i U^j \rangle + \rho \langle \frac{1}{2}U^2 U^j \rangle]_j = \rho u_j a^j. \quad (5.35)$$

Considering each term in this equation we recognize that the specific internal energy of the gas is  $\varepsilon = \frac{1}{2}\langle U^2 \rangle$ . As we identified above, the stress tensor is  $\rho \langle U^i U^j \rangle = -T^{ij} = P\delta^{ij}$ . Finally the heat flux  $F$  as

$$F = \rho \langle \frac{1}{2}U^2 U^j \rangle, \quad (5.36)$$

which is the energy flux in the gas resulting from microscopic particle motions. Thus we have

$$(\rho\varepsilon + \frac{1}{2}\rho u^2)_{,t} + [\rho u^j (\varepsilon + \frac{P}{\rho} + \frac{1}{2}u^2) + F^j]_j = \rho u_i g^i \quad (5.37)$$

which is the same as Eq. (3.5) except neglecting the local source term  $\rho q$  from, e.g., nuclear reactions.

This approach to deriving the fluid equations is identical to constructing a hierarchy of moments of the Boltzmann equation in a completely analogous fashion to how we constructed moments of the distribution function to arrive at physical quantities such as the density, momentum, and stress in the previous section. Inspecting our resulting equations in light of the moment formalism, we see that each evolution equation depends on the next higher moment of the distribution function. For instance, the continuity equation, which is just an evolution equation for the mass density, depends on the divergence of the momentum density  $\rho u^i$ . Similarly, the momentum evolution equation depends on the divergence of the energy density  $\rho u^i u^j$ . This is a common feature of moment formalisms, requiring that the hierarchy be “closed” by some equation or equations that give the higher moments in terms of the lower moments which are being directly solved for. In the case of the fluid equations, this is achieved by an equation of the state which gives the pressure  $P$  in terms of the the energy and density  $\varepsilon$ ,  $\rho$  (or any other two thermodynamic variables which we are solving for). Additionally, we supply an equation for solving for the heat flux  $F$  [c.f. Eq. (3.7)].



# 6

## *The Equation of State*

In statistical equilibrium, we can describe a system of particles by a distribution function  $f(\mathbf{p}, \mathbf{x}) d^3 p d^3 x$ , such that the number of particles is

$$N = \int d^3 p d^3 x f(\mathbf{p}, \mathbf{x}), \quad (6.1)$$

where the integration is over the phase spaces of momentum and position coordinates  $(\mathbf{p}, \mathbf{x})$ . In an ideal gas, the particles do not interact. In such a case, the distribution function  $f = f(\mathbf{p})$  does not depend on position. The integration over  $d^3 x$  just gives a factor of the volume, so the number density is  $n = \int d^3 p f(\mathbf{p})$ . From equation (6.1), we can get our other thermodynamic quantities: for example.

$$\frac{E}{V} \equiv u = \int \varepsilon(\mathbf{p}) f(\mathbf{p}) d^3 p \quad \text{energy per unit volume} \quad (6.2)$$

$$P = \int (\mathbf{p} \cdot \mathbf{e}_z) (\mathbf{v} \cdot \mathbf{e}_z) f(\mathbf{p}) d^3 p \quad \text{pressure.} \quad (6.3)$$

Here  $\varepsilon$  is the particle energy and  $\mathbf{v}$  the velocity.

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EXERCISE 6.1 — For an isotropic momentum distribution, show that

$$P = \frac{1}{3} \int |\mathbf{p}| |\mathbf{v}| f(\mathbf{p}) d^3 p.$$

Then show that for a non-relativistic gas,  $P = (2/3)E/V$ , and that for a relativistic gas,  $P = (1/3)E/V$ .

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### *6.1 Connection to thermodynamics*

Once we have the distribution functions, we can get all of the other thermodynamic properties from the thermodynamic relations: in what follows let  $N = nV$  be the total number of particles, with  $V$  being the volume of the system. The total energy is then  $E = uV$ , and

the entropy is  $S$ , and we have

$$F = E - TS, \quad \text{Helmholtz free energy,} \quad (6.4)$$

$$H = E + PV, \quad \text{Enthalpy,} \quad (6.5)$$

$$\mu N = G = F + PV, \quad \text{Gibbs free energy.} \quad (6.6)$$

For example, we have in the non-degenerate limit that

$$\mu = k_B T \ln K = k_B T \ln \left[ \frac{n}{g} \left( \frac{2\pi\hbar^2}{mk_B T} \right)^{3/2} \right], \quad (6.7)$$

and so we could write the entropy per unit mass as

$$\begin{aligned} s \equiv \frac{S}{Nm} &= \frac{1}{Nm} \frac{E + PV - \mu N}{T} \\ &= \frac{k_B}{m} \left\{ \frac{5}{2} + \ln \left[ \frac{g}{n} \left( \frac{mk_B T}{2\pi\hbar^2} \right)^{3/2} \right] \right\}. \end{aligned} \quad (6.8)$$

In this equation I have used the ideal non-degenerate values  $E = (3/2)Nk_B T$ ,  $PV = Nk_B T$  and have denoted the mass per particle as  $m$  and the degeneracy of the spin-states as  $g$ .

**EXERCISE 6.2** — In an external field (i.e., gravitational) the chemical potential, which is the change in energy when the number of particles is increased, must include the potential. Consider an ideal gas in a planar atmosphere of constant gravitational acceleration  $g$ . Write the chemical potential as  $\mu(z) = \mu_{\text{id}} + \Phi$ , where  $\mu_{\text{id}}$  is the chemical potential for an ideal gas in the absence of gravity, and  $\Phi$  is the gravitational potential. For an atmosphere in complete equilibrium ( $\mu = \text{const}$ ,  $T = \text{const}$ ), calculate the pressure as a function of position,  $P = P(z)$ , and show that it agrees with considerations from hydrostatic balance, equation (2.10).

## 6.2 An ideal Fermi gas

For fermions, particles with half-integer spin, it can be shown that

$$f(p) = \frac{g}{(2\pi\hbar)^3} \left[ \exp \left( \frac{\varepsilon - \mu}{k_B T} \right) + 1 \right]^{-1}. \quad (6.9)$$

In this equation  $\varepsilon(p)$  is the energy of a particle,  $\mu$  is the *chemical potential*,  $T$  is the temperature, and  $g$  denotes the number of particles that can occupy the same energy level (for spin-1/2 particles,  $g = 2$ ). The connection to thermodynamics is via the relations

$$\frac{1}{T} = \left( \frac{\partial S}{\partial E} \right)_{N,V}, \quad -\frac{\mu}{T} = \left( \frac{\partial S}{\partial N} \right)_{E,V},$$

and

$$TS = PV - \mu N + E;$$

these are derived in standard texts. Let's explore what happens in various limits.

### *Non-degenerate, non-relativistic limit*

First, let's take  $K \equiv \exp(\mu/k_B T) \ll 1$ . (In the literature,  $K$  is called the *fugacity*.) Then in equation (6.9), we see that the exponential term dominates. If our system is isotropic, then  $d^3 p = 4\pi p^2 dp$ , and we'll use this substitution from now on. We then have for the number density

$$n(\mu, T) = \frac{gK}{2\pi^2 \hbar^3} \int_0^\infty \exp\left(-\frac{\varepsilon}{k_B T}\right) p^2 dp. \quad (6.10)$$

To do the integral, notice that since  $2m\varepsilon = p^2$ , we have  $p^2 dp = m(2m\varepsilon)^{1/2} d\varepsilon$ ; making the substitution  $x = \varepsilon/(k_B T)$ , we get

$$n(\mu, T) = \frac{gK}{2\pi^2 \hbar^3} \sqrt{2}(mk_B T)^{3/2} \int_0^\infty x^{1/2} e^{-x} dx. \quad (6.11)$$

You have all struggled with this integral in your past, but to avoid unpleasant flashbacks, I will just tell you that it is  $\sqrt{\pi}/2$ . So, we have our first result (but we still don't know what it means),

$$n(\mu, T) = K \left[ g \left( \frac{mk_B T}{2\pi \hbar^2} \right)^{3/2} \right]. \quad (6.12)$$

Let's forge on a little further, though, and try to get the energy per unit volume  $u$ . Once we have  $u$ , we know we can get the pressure from the relation for a non-relativistic gas,  $P = 2/3 u$ . Using equations (6.2) and (6.9),

$$u(\mu, T) = \frac{gK}{2\pi^2 \hbar^3} \int_0^\infty \exp\left(-\frac{\varepsilon}{k_B T}\right) \varepsilon p^2 dp. \quad (6.13)$$

Let's repeat our trick of changing variables from  $p$  to  $x = \varepsilon(p)/k_B T$ ; we then have

$$u(\mu, T) = \frac{gK}{2\pi^2 \hbar^3} \sqrt{2} k_B T (mk_B T)^{3/2} \int_0^\infty x^{3/2} e^{-x} dx. \quad (6.14)$$

Did you notice that if we integrate by parts,

$$\int_0^\infty x^{3/2} e^{-x} dx = \frac{3}{2} \int_0^\infty x^{1/2} e^{-x} dx = \frac{3}{4} \sqrt{\pi},$$

we get the integral we already solved in equation (6.11)? Putting everything together and using the expression for  $n$  (eq. [6.12]), we have

$$u(\mu, T) = \frac{3}{2} \left\{ K \left[ g \left( \frac{mk_B T}{2\pi \hbar^2} \right)^{3/2} \right] \right\} k_B T = \frac{3}{2} n k_B T. \quad (6.15)$$

which gives us the pressure,

$$P = \frac{2}{3}u = nk_{\text{B}}T. \quad (6.16)$$

Whoo-hoo! We've rediscovered the ideal gas.

Now we have to understand this chemical potential  $\mu$ . We can solve equation (6.12) for  $\mu$ ,

$$\exp\left(\frac{\mu}{k_{\text{B}}T}\right) = K = n \left[ g \left( \frac{mk_{\text{B}}T}{2\pi\hbar^2} \right)^{3/2} \right]^{-1}. \quad (6.17)$$

Now  $K$  is dimensionless, a number, so the thing in [ ] must have dimensions of number density. Let's call it  $n_Q$ . Our chemical potential is then  $\mu = k_{\text{B}}T \ln(n/n_Q)$ . To understand the significance of  $n_Q$ , let's calculate the uncertainty in position of a particle having energy  $k_{\text{B}}T$ ; from Heisenberg, we have

$$\Delta x \approx \frac{\hbar}{\Delta p} \sim \frac{\hbar}{\sqrt{mk_{\text{B}}T}}$$

where I am dropping numerical factors and I've made the substitution  $\Delta p \sim p \approx \sqrt{mk_{\text{B}}T}$ . Now what happens if I pack the particles so that on average there are  $g$  particles per box of volume  $(\Delta x)^3$ ? In that case the density would be  $n = g(\sqrt{mk_{\text{B}}T}/\hbar)^3 \approx n_Q$ . So, what appears in the chemical potential is the ratio of the density to that density at which the particles are packed so closely that the uncertainty in their positions is the same size as the typical inter-particle spacing. In the ideal-gas limit  $K \ll 1$ , which makes sense:  $n \ll n_Q$ , so the particles are very far apart compared to their thermal de Broglie wavelengths, and quantum effects ought to be unimportant.

**EXERCISE 6.3** — Get the first order corrections to the Maxwell-Boltzmann gas. Take the fugacity  $K \ll 1$ , and expand the Fermi-Dirac distribution (eq. [6.9]) to lowest order in  $K \exp[-\varepsilon/(k_{\text{B}}T)]$ . Show that

$$\begin{aligned} n(K, T) &= n_0 \left( 1 - 2^{-3/2}K \right) \\ u(K, T) &= \frac{3}{2}n_0 k_{\text{B}}T \left( 1 - 2^{-5/2}K \right), \end{aligned}$$

where  $n_0$  is the density in the limit  $K \rightarrow 0$ . Then derive the equation of state  $P = P(n, T)$  to lowest order in  $K$ . For a given  $(n, T)$ , is the pressure larger or smaller than that of the ideal Maxwell-Boltzmann limit?

### Degenerate, non-relativistic limit

When  $n \gtrsim n_Q$ , we can no longer use the approximation  $K \ll 1$ , so let's go to the opposite limit, for which  $\mu \gg k_{\text{B}}T$ . In this case, notice

from equation (6.9) that

$$f(p) \approx \frac{g}{(2\pi\hbar)^3} \begin{cases} 1 & \varepsilon < \mu \\ 0 & \varepsilon > \mu \end{cases}. \quad (6.18)$$

We can think of this as inserting  $g$  particles in each energy level, starting with the lowest energy level and continuing until all of the particles are used. The last particle is inserted with energy  $\varepsilon \approx \mu$ . The only levels that will be partially filled will be those lying in a thin band  $\varepsilon \approx \mu \pm k_B T$ . If that is the case, we can make the following approximation. Let's take the limit  $T \rightarrow 0$ , and define the *Fermi energy* by  $\varepsilon_F = \mu|_{T \rightarrow 0}$  and the *Fermi momentum* by  $p_F = \sqrt{2m\varepsilon_F}$ . We can then write equation (6.1) as

$$n(\mu) = \frac{1}{\pi^2\hbar^3} \int_0^{p_F} p^2 dp, \quad (6.19)$$

since the integrand is zero for  $p > p_F$ . The main application is for electrons, which are spin one-half, so we substitute  $g = 2$ . Now this is an easy integral,

$$n(\mu) = \frac{p_F^3}{3\pi^2\hbar^3} = \frac{(2m\varepsilon_F)^{3/2}}{3\pi^2\hbar^3}, \quad (6.20)$$

or  $\mu \approx \varepsilon_F = (3\pi^2 n)^{2/3} \hbar^2 / (2m)$ . Let's get the energy per unit volume and the pressure,

$$u(\mu) = \frac{1}{\pi^2\hbar^3} \int_0^{p_F} \frac{p^2}{2m} p^2 dp = \frac{p_F^5}{5\pi^2\hbar^3}. \quad (6.21)$$

Comparing this with equation (6.20), we have

$$u = \frac{3}{5} n \varepsilon_F, \quad (6.22)$$

$$P = \frac{2}{5} n \varepsilon_F. \quad (6.23)$$

To lowest order, neither  $u$  nor  $P$  depend on  $T$ . Substituting for  $\varepsilon_F$  in equation (6.23) gives us the equation of state,

$$P = \frac{2}{5} \left(3\pi^2\right)^{2/3} \frac{\hbar^2}{2m} n^{5/3}. \quad (6.24)$$

Notice that in equation (6.17),  $n_Q \propto m^{3/2}$ . This means that at any given temperature,  $n_Q$  for electrons is  $1836^{3/2} = 80,000$  times smaller than it is for protons, not to mention helium or heavier nuclei. As a result, the electrons will become degenerate ( $n \gtrsim n_Q$ ) at a much lower mass density than the ions. A common circumstance, then, is to have a mixture of degenerate electrons and ideal ions (we will deal with non-ideal corrections due to electric forces later).

Now, let's estimate the boundary between the non-degenerate and degenerate regimes. At a given temperature, we know in the low-density limit that the electrons obey the ideal gas law (eq. [6.16]) and in the high-density limit the electrons are degenerate (eq. [6.24]). So, let's extrapolate our two limiting expressions for the pressure and see where they meet,

$$n_e k_B T = P_{e,\text{ideal}} \sim P_{e,\text{deg.}} = \frac{2}{5} n_e \varepsilon_F, \quad (6.25)$$

or  $\varepsilon_F \approx k_B T$ . No surprise here. Notice that the ratio

$$\frac{\varepsilon_F}{k_B T} \sim \frac{(3\pi^2)^{2/3} \hbar^2}{2m_e k_B T} n_e^{2/3} \sim \left( \frac{n_e}{n_Q} \right)^{2/3}, \quad (6.26)$$

so marking the onset of degeneracy with  $\varepsilon_F \sim k_B T$  also makes sense from that aspect as well. Our boundary in the density-temperature plane between the non-degenerate and degenerate regimes is then determined by setting  $k_B T = \varepsilon_F$ ,

$$T = \frac{(3\pi^2)^{2/3} \hbar^2}{2m_e k_B} \left( \frac{Y_e \rho}{m_u} \right)^{2/3} = 3.0 \times 10^5 \text{ K} (Y_e \rho)^{2/3}. \quad (6.27)$$

If the temperature falls below this value, the electrons will be degenerate. Here  $Y_e$  is the electron molar fraction, or electron abundance and  $m_u$  is the atomic mass unit; consult §3.2 for details.

**EXERCISE 6.4**— Repeat the derivation of equation (6.22) and (6.23) for a relativistic Fermi gas. What is the expression for the temperature at which the gas becomes degenerate (cf. eq. [6.27]) in this case?

### 6.3 Fermi-Dirac integrals

This condition for the onset of degeneracy, eq. (6.27), is only a rule-of-thumb; in any serious calculation we would want to calculate the electron thermal properties from the exact integrals

$$n(\mu, T) = \frac{\sqrt{2}(mk_B T)^{3/2}}{\pi^2 \hbar^3} \int_0^\infty \frac{x^{1/2} dx}{\exp(x - \psi) + 1} \quad (6.28)$$

$$P(\mu, T) = \frac{(2mk_B T)^{3/2} (k_B T)}{3\pi^2 \hbar^3} \int_0^\infty \frac{x^{3/2} dx}{\exp(x - \psi) + 1}, \quad (6.29)$$

where  $\psi = \mu / (k_B T)$ . These integrals cannot be done analytically, but they occur so frequently that there are many published tables and numerical approximation schemes<sup>1</sup>. Specifically, the *non-relativistic Fermi-Dirac integral of order v* is defined as

$$F_v(\psi) = \int_0^\infty \frac{x^v dx}{\exp(x - \psi) + 1}. \quad (6.30)$$

<sup>1</sup> F. X. Timmes and F. D. Swesty. The Accuracy, Consistency, and Speed of an Electron-Positron Equation of State Based on Table Interpolation of the Helmholtz Free Energy. *ApJS*, 126:501, 2000

One can (numerically) invert equation (6.28) to solve for the chemical potential  $\psi k_B T$ .

The general case for a relativistic Fermi gas is left as an exercise.

#### 6.4 Relativistic photon gas

Photons are *bosons*—they have spin 1. For bosons, the distribution function is similar to that in equation (6.9), but with the +1 replaced by −1 in the denominator. In addition, photon number is not conserved: one can freely create and destroy photons. This implies that their chemical potential is zero. Also,  $g = 2$  for photons: there are two independent polarization modes. Putting all of these together, we can write energy per unit volume as

$$u = \frac{1}{\pi^2 \hbar^3} \int_0^\infty \varepsilon p^2 \left[ \exp \left( \frac{\varepsilon}{k_B T} \right) - 1 \right]^{-1} dp. \quad (6.31)$$

Now, use the fact that  $p = \varepsilon/c$  and change variables to  $x = \varepsilon/(k_B T)$  to get

$$u = \frac{k_B^4 T^4}{\pi^2 c^3 \hbar^3} \int_0^\infty \frac{x^3 dx}{e^x - 1}.$$

This integral is a classic and is equal to  $\pi^4/15$ . Hence the energy per unit volume and the pressure are

$$u = \left( \frac{k_B^4 \pi^2}{15 c^3 \hbar^3} \right) T^4 = a T^4 \quad (6.32)$$

$$P = \frac{1}{3} a T^4. \quad (6.33)$$

In CGS units,  $a = 7.566 \times 10^{-15}$  erg cm<sup>−3</sup> K<sup>−4</sup>. With this energy and pressure, we can compute the other thermodynamical quantities.

---

EXERCISE 6.5— Show that the entropy of a blackbody radiation is

$$S_{\text{rad}} = \frac{4}{3}aT^3V.$$

Now consider a mixture of blackbody radiation and an ideal gas. For the ideal gas, the entropy is just

$$S_{\text{gas}} = N\mu m_u s,$$

where  $s$  is given by Eq. (4.25). Along an adiabat,  $dS = d(S_{\text{rad}} + S_{\text{gas}}) = 0$ ; use this to find an expression for  $\nabla_{\text{ad}}$  in terms of

$$\beta \equiv \frac{P_{\text{gas}}}{P},$$

where  $P = P_{\text{gas}} + P_{\text{rad}}$ .

Show that the expression for  $\nabla_{\text{ad}}$  has the correct limiting values for  $\beta \rightarrow 0$  and  $\beta \rightarrow 1$ .

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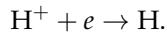
## 6.5 Chemical Equilibrium: The Saha Equation

Consider a reaction,  $A + B + \dots \rightarrow C + D + \dots$ . When this reaction comes into equilibrium, we are at a maximum in entropy, and the condition for equilibrium is that the energy cost, at constant entropy to run the reaction in the forward direction is the same as to run the reaction in reverse. This can be expressed in terms of chemical potentials as

$$\mu_A + \mu_B + \dots \rightarrow \mu_C + \mu_D + \dots \quad (6.34)$$

Note in this formalism that a reaction  $2A \rightarrow B$  would be expressed as  $2\mu_A = \mu_B$ .

As a worked example, we consider the ionization equilibrium of hydrogen,



To use equation (6.34), we need to have both sides on the same energy scale. The reaction in the exothermic direction; that is, heat is evolved if the reaction proceeds as written. This means that the right-hand side is more bound, and its minimum energy is less than that of the left-hand side. To get both sides on the same energy scale, we must subtract the binding energy,  $Q = 13.6 \text{ eV}$ , from the right-hand side:

$$\mu_+ + \mu_- = \mu_0 - Q. \quad (6.35)$$

Another way to see why  $Q$  appears is to add the rest mass for each species to its chemical potential; collecting all terms on the right, we would then have a term  $(m_0 - m_+ - m_-)c^2 = -Q$ . Some ionization potentials, along with the half-ionization temperature for a given electron number density  $n_e$ , are given in Table 6.1.

element	H <sup>-</sup>	Na	H	He
$Q/\text{eV}$	0.75	5.14	13.6	24.6
$T_{1/2}(n_e = 10^{13} \text{ cm}^{-3})/\text{K}$	600	3300	8000	13900
$T_{1/2}(n_e = 10^{16} \text{ cm}^{-3})/\text{K}$	900	5000	11800	20200

Table 6.1: Selected ionization potentials and half-ionization temperatures

For a non-degenerate plasma, we can insert eq. (6.7) into eq. (6.35), divide through by  $k_B T$ , and take the exponential to obtain

$$\frac{n_+ n_-}{n_0} = \frac{g_+ g_-}{g_0} \left( \frac{m_- k_B T}{2\pi\hbar^2} \right)^{3/2} \exp \left( -\frac{Q}{k_B T} \right). \quad (6.36)$$

The number density of all hydrogen in the gas is  $n_0 + n_+ = n_H$ .

Denote the ionized fraction by  $x = n_+/n_H = n_i/n_H$ , so that the left-hand side of equation (6.36) is  $n_H x^2 / (1-x)$ . In the hydrogen atom ground state, the electron spin and proton spin are either aligned or anti-aligned. These states are very nearly degenerate, so that  $g_0 = 2$ . Both the proton and electron have spin 1/2; there are really only two available states, however, because of the freedom in choosing our coordinate system. As a result,  $g_+ g_- = 2$  as well.

Inserting these factors into equation (6.36), and using  $k_B = 8.6173 \times 10^{-5} \text{ eV/K}$ , we obtain

$$\frac{x^2}{1-x} = \frac{2.41 \times 10^{21} \text{ cm}^{-3}}{n_H} \left( \frac{T}{10^4 \text{ K}} \right)^{3/2} \exp \left( -\frac{15.78 \times 10^4 \text{ K}}{T} \right). \quad (6.37)$$

This equation defines a set of points in the  $\rho - T$  plane for which  $x = 1/2$ . We may take this set of points to mark the boundary between neutral and ionized hydrogen. At fixed density, the transition from neutral to fully ionized is very rapid.

#### EXERCISE 6.6—

1. Solve equation (6.37) for a density of  $10^{16} \text{ cm}^{-3}$  (a fiducial value for the solar photosphere) and find the half-ionization temperature, i.e., the temperature at which  $x = 0.5$ . Explain the reason for the discrepancy between the half-ionization temperature and  $k_B T = 13.6 \text{ eV}$ .
2. At this half-ionization temperature, what is the occupancy of the excited levels of the hydrogen atom? Do we need to worry about corrections to the ionization from these excited states?

## 6.6 Coulomb interactions

Under the conditions in a stellar interior, most of the atoms are ionized, and the stellar matter consists of positively charged nuclei and ions and negatively charged electrons. A *plasma* is defined as a gas

of charged particles in which the kinetic energy of a typical particle is much greater than the potential energy due to its nearest neighbors. To make this quantitative, consider a gas with only one species present, with charge  $q$ . Let the mean spacing between particles be  $a$ ; clearly the number density of such particles is  $n = (4\pi a^3/3)^{-1}$ . We may then take the quantity

$$\Gamma \equiv \frac{q^2}{ak_B T} \quad (6.38)$$

as indicating the relative importance of potential to kinetic energy. In a classical plasma,  $\Gamma \ll 1$ . Note, however, that systems with  $\Gamma > 1$  are often (confusingly) called *strongly coupled plasmas*. The meaning is usually clear from context.

**EXERCISE 6.7**— Show that for a non-relativistic plasma, the magnetic interaction between two charged particles is much less than the electrostatic interaction.

### Debye shielding

Imagine a typical charged particle in a plasma. Very close to the particle, we expect the electrostatic potential to be that of an isolated charge  $\Phi = q/r$ . Far from the particle, there will be many other particles surrounding it, and the potential is screened. For example, a positive ion will tend to attract electrons to be somewhat, on average, closer to it than other ions: we say that the ion *polarizes* the plasma. As a result of this polarization, the potential of any particular ion should go to zero much faster than  $1/r$  due to the “screening” from the enhanced density of opposite charges around it.

Let's consider a plasma having many ion species, each with charge  $Z_i$ , and electrons. About any selected ion  $j$ , particles will arrange themselves according to Boltzmann's law,

$$n_i(r) = n_{i0} \exp \left[ -\frac{Z_i e \Phi(r)}{k_B T} \right]. \quad (6.39)$$

Here  $n_{i0}$  is the density of particle  $i$  far from the charge  $j$ , and  $r$  is the distance between particles  $i$  and  $j$ . (A similar equation holds for the electrons, with  $Z$  replaced by  $-1$ .) To solve for the potential, we can use Poisson's equation,

$$\nabla^2 \Phi = -4\pi \sum_i Z_i e n_i(r) + 4\pi e n_e(r). \quad (6.40)$$

Our assumption is that the term in the exponential of equation (6.39) is small, so we may expand it to first order in  $\Phi$  and substitute that

expansion into equation (6.40) to obtain in spherical geometry

$$\frac{1}{r} \frac{\partial^2}{\partial r^2} (r\Phi) = -4\pi e \left[ \sum_i n_{i0} Z_i \left( 1 - \frac{Z_i e \Phi}{k_B T} \right) - n_{e0} \left( 1 + \frac{e \Phi}{k_B T} \right) \right].$$

The overall charge neutrality of the plasma implies that  $n_{e0} = \sum_i Z_i n_{i0}$ ; using this to simplify the above equation gives

$$\frac{1}{r} \frac{\partial^2}{\partial r^2} (r\Phi) = \left[ \frac{4\pi e^2}{k_B T} \sum_i n_{i0} (Z_i^2 + Z_i) \right] \Phi \equiv \lambda_D^{-2} \Phi. \quad (6.41)$$

The quantity in [ ] has dimensions of reciprocal length squared and we define it as  $(1/\lambda_D)^2$  with  $\lambda_D$  being called the *Debye length*.

Multiplying equation (6.41) by  $r$ , integrating twice, and determining the constant of integration from the condition that as  $r \rightarrow 0$ ,  $\Phi \rightarrow Z_j e / r$  gives the self-consistent potential

$$\Phi = \frac{Z_j e}{r} \exp \left( -\frac{r}{\lambda_D} \right). \quad (6.42)$$

The Debye length  $\lambda_D$  determines the size of the screening cloud around the ion.

In order for the above derivation to be valid, we require that  $\lambda_D \gg a$ , where  $a$  is the mean ion spacing; otherwise, there won't be any charges in our cloud to screen the potential! Equivalently, we require the number of particles in a sphere of radius  $\lambda_D$  to be large,

$$\frac{4\pi}{3} \lambda_D^3 \sum_i n_i \gg 1. \quad (6.43)$$

This condition must hold if we are to treat the gas as a plasma.

## EXERCISE 6.8—

1. Show that equation (6.43) is equivalent to  $\Gamma \ll 1$  for a single species plasma.
2. Show that the net charge in the shielding cloud about an ion of charge  $Ze$  is  $-Ze$ ; the shielding cloud cancels out the ion's charge.

*Corrections to the ideal gas EOS*

In a plasma the particles are not independent: shake one particle and other nearby particles will shake as well. In statistical mechanics, this requires introducing *correlation functions* to derive the equation of state. We'll adopt the more intuitive approach of Debye and Hückel to get the lowest-order correction to the ideal gas EOS. First, the total electrostatic energy in a volume  $V$  is

$$E_{\text{Coul}} = \frac{1}{2}V \sum_j Z_j e n_j \Phi_j. \quad (6.44)$$

Here  $\Phi_j$  is the potential at a particle  $j$  due to all the other particles in the plasma. Now, we computed the total potential around a particle (eq. [6.42]); expanding and subtracting off the self-potential of particle  $j$  gives  $\Phi_j = -Z_j e / \lambda_D$ . Inserting this into equation (6.44) and expanding gives

$$E_{\text{Coul}} \approx -V \left( \frac{\pi}{k_B T} \right)^{1/2} e^3 \left[ \sum_i n_{i0} (Z_i^2 + Z_i) \right]^{3/2}. \quad (6.45)$$

This energy is to be added to the kinetic energy of the gas. The effect of the electrostatic interactions is to *decrease* the energy in the gas, that is, to make it more bound.

We can't directly get the pressure from equation (6.45) because the equation isn't in terms of  $S$  and  $V$  (recall that  $P = -(\partial E / \partial V)_S$ ) but rather in terms of  $T$  and  $V$ . In order to get the pressure, we first must find the Helmholtz free energy  $F$ . To do this, we integrate the thermodynamical identity

$$E = -T^2 \left( \frac{\partial}{\partial T} \right)_V \left( \frac{F}{T} \right)$$

and then take  $P = -(\partial F / \partial V)_{T,N}$  to obtain

$$P_{\text{Coul}} \approx -\frac{e^3}{3} \left( \frac{\pi}{k_B T} \right)^{1/2} \left[ \frac{(\langle Z^2 \rangle + \langle Z \rangle) \rho}{\langle A \rangle m_u} \right]^{3/2}. \quad (6.46)$$

The effect of Coulomb interactions is to decrease the pressure below the ideal gas value.

### Coulomb corrections when the electrons are degenerate

The above discussion holds only when both the electrons and ions are non-degenerate. What happens when the electrons are degenerate? In that case the kinetic energy is of order the Fermi energy, not the temperature. We might think to replace  $k_B T$  with  $\epsilon_F$  in equation (6.38). Recalling the formula for  $\epsilon_F$  from § 6.2, we have the condition for the Coulomb interactions to be weak,

$$\frac{e^2}{a\epsilon_F} = \left(\frac{4\pi n}{3}\right)^{1/3} \left(\frac{m_e e^2}{\hbar^2}\right) \frac{2}{(3\pi^2 n)^{2/3}} < 1. \quad (6.47)$$

Here  $m_e$  and  $n$  denote, respectively, the electron mass and number density. Do you recognize the quantity  $\hbar^2/(m_e e^2)$ ? It is the Bohr radius,  $a_B$ . What is  $a_B$  doing in this equation? Well, we are looking for a quantum mechanical system in which the Coulomb interaction is comparable to the non-relativistic kinetic energy. Does that sound like any system you've seen before?

Cleaning up equation (6.47), our condition for the electrons to be weakly interacting when degenerate is

$$\left(\frac{2^{5/3}}{3\pi}\right) \left(na_B^3\right)^{-1/3} < 1, \quad (6.48)$$

or, in terms of mass density  $\rho$  and electron fraction  $Y_e$ ,  $(Y_e\rho) > 0.4 \text{ g cm}^{-3}$ . As the density increases, the electron gas becomes more ideal, that is, the electrostatic interaction matters less and less.

Just to complete the discussion on electrons, what if the electrons are relativistic? In this case,  $\epsilon_F = p_F c = (3\pi^2 n)^{1/3} \hbar c$ , and

$$\frac{e^2}{a\epsilon_F} = \left(\frac{4}{9\pi}\right)^{1/3} \left(\frac{e^2}{\hbar c}\right) = 3.8 \times 10^{-3}. \quad (6.49)$$

In this case  $\epsilon_F \propto n^{1/3}$  so the density dependence cancels. You will note the appearance of the fine structure constant  $\alpha_F = e^2/(\hbar c)$ , as you might have expected when dealing with relativistic electrons and electrostatics.

Under astrophysical conditions, we can almost always regard degenerate electrons as being ideal. What about the ions? They are not usually degenerate under conditions of interest. We can get a simple expression if we go to the opposite limit, in which the electrons are very degenerate. In that case, the electrons are an ideal gas and hence have uniform density. If the temperature is low enough, the ions will have  $Z^2 e^2 / (ak_B T) \gg 1$ ; in this case we might expect the ions will arrange themselves into a lattice that maximizes the inter-ionic spacing.

To get an estimate of the electrostatic energy, let's compute the energy of a charge-neutral sphere centered on a particular ion of charge

$Z_i$ . Because the electrons have a uniform density,  $Y_e \rho / m_u$ , we can find the radius of the sphere  $a$  by requiring it to have  $Z_i$  electrons,

$$\frac{4\pi}{3} a^3 \left( Y_e \frac{\rho}{m_u} \right) = Z_i, \quad (6.50)$$

or  $a = [3Z_i m_u / (4\pi Y_e \rho)]^{1/3}$ . The potential energy of this sphere has two components. The first is due to electron-electron interactions,

$$E_{ee} = \int_0^a \frac{q(r) dq}{r} = \frac{3}{5} \frac{Z_i^2 e^2}{a}, \quad (6.51)$$

where  $q(a) = Z_i e (r/a)^3$  is the charge in a sphere of radius  $r < a$ . The second component of the potential energy is due to the ion-electron interaction,

$$E_{ei} = -Z_i e \int_0^a \frac{dq}{r} = -\frac{3}{2} \frac{Z_i^2 e^2}{a}. \quad (6.52)$$

Combining equations (6.51), (6.52), and (6.50) gives the total electrostatic energy for a single ion-sphere,

$$E = -\frac{9}{10} \frac{Z_i^2 e^2}{a} = -\frac{9}{10} Z_i^{5/3} e^2 \left( \frac{4\pi}{3} \frac{Y_e \rho}{m_u} \right)^{1/3}. \quad (6.53)$$

Multiplying this by  $n_i = Y_i \rho / m_u$ , summing over all ion species  $i$ , and defining  $\langle Z^{5/3} \rangle = n^{-1} \sum Z_i^{5/3} n_i$  where  $n = \sum n_i$ , gives the total Coulomb energy per volume,

$$E_{\text{Coul}} = -\frac{9}{10} n k_B T \left[ \frac{\langle Z^{5/3} \rangle e^2}{k_B T} \left( \frac{4\pi}{3} \frac{Y_e \rho}{m_u} \right)^{1/3} \right]. \quad (6.54)$$

Notice that the quantity in [] reduces to  $\Gamma$  for a single-species plasma. If we therefore define  $\Gamma \equiv []$  for a multi-component plasma, we have  $E_{\text{Coul}} \approx -0.9 n k_B T \Gamma$ ; the pressure is then  $P_{\text{Coul}} = E_{\text{Coul}} / 3 = -0.3 n k_B T \Gamma$ . This holds in the limit  $\Gamma \gg 1$ .

---

EXERCISE 6.9—

1. In the zero-temperature limit (electrons are fully degenerate) use the charge-neutral sphere approximation (p. 68) to calculate the density at which completely ionized  $^{56}\text{Fe}$  has zero pressure.
  2. Estimate the pressure that would be required to compress  $^{56}\text{Fe}$  at the density found in part 1.
  3. Compute the mass-radius relation for a cold object (“a rock”). Write the total pressure as the sum of electron (zero-temperature limit) and Coulomb pressure (use the charge-neutral sphere approximation, eq. [6.54] and following text). To make this easier, use the virial scalings for density and pressure to obtain a relation between mass and radius. Find the mass having the largest radius, and express this mass in terms of fundamental physical constants. How does it compare with the mass of Jupiter? Scale the mass and radii to that of Jupiter, and plot  $R(M)$  for objects composed of pure  $^1\text{H}$ , pure  $^4\text{He}$ , and pure  $^{12}\text{C}$ . Also indicate on this plot the masses and radii of the Jovian planets for comparison.
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# 7

## *Radiation Transport*

The sun is very opaque. Were photons able to stream freely, they would exit in  $\sim R_\odot/c = 2.0\text{ s}$ . Given the luminosity of the sun, however, we derived that the time for the sun to radiate away its stored thermal energy is instead millions of years (see eq. [2.15]). As a result, we can regard the sun as a cavity filled with photons with a very slight leakage. This is the description commonly invoked to describe blackbody radiation, and we expect that in the interior of the sun, the radiation can be described by a photon gas in thermal equilibrium at the ambient temperature.

### 7.1 Description of the Radiation Field

Consider a cavity containing a gas of photons. In general we can describe the mean number of photons in this cavity as

$$N = \int f(\mathbf{p}, \mathbf{x}) d^3x d^3p. \quad (7.1)$$

Here  $f$  is a distribution function, as described in Chapter 6; if we are in thermal equilibrium,  $f$  is of course the Bose-Einstein distribution (cf. eq. [6.31]), but our discussion here will be more general. Consider a small opening on our cavity with area  $dA$  and unit normal  $\hat{\mathbf{n}}$ . The energy incident on this area in a time  $dt$  having propagation vector along  $\hat{\mathbf{n}}$  and propagating into solid angle  $d\Omega$  (see Fig. 7.1) is found by integrating equation (7.1) over a volume  $d^3x = cdt dA$ ,

$$dE = dA cdt \left( p^2 dp d\Omega \right) h\nu f.$$

Since the photon momentum is  $p = h\nu/c$ , we have

$$I_\nu \equiv \frac{dE}{dt dA d\Omega d\nu} = \frac{h^4 \nu^3}{c^2} f. \quad (7.2)$$

This defines the *specific intensity*  $I_\nu$ . It is easy to show that in the absence of interactions with matter,  $I_\nu$  is conserved along a ray (see, e.g., Rybicki & Lightman).

If the photons are in thermal equilibrium, then  $f$  is the Bose-Einstein distribution,  $f = (2/h^3)(\exp[h\nu/k_B T] - 1)^{-1}$ , and the specific intensity becomes

$$B_\nu \equiv \frac{2h\nu^3}{c^2} \left[ \exp\left(\frac{h\nu}{k_B T}\right) - 1 \right]^{-1}. \quad (7.3)$$

Here  $B_\nu$  is called the *Planck function*.

The specific intensity plays the role of the *distribution function* for radiation. We can similarly derive important physical quantities by taking moments of the specific intensity, but in angle rather than momentum space. The *zeroth* moment of the specific intensity is

$$J_\nu = \frac{1}{4\pi} \int I_\nu d\Omega, \quad (7.4)$$

and is called the *mean intensity* as it is essentially an angle average of the specific intensity.

The energy density per frequency  $u_\nu$  can be defined as  $dE/(cdt dA d\nu)$ , that is, the energy per unit frequency that is in a cylinder of length  $cdt$  and cross-sectional area  $dA$ ; comparing  $u_\nu$  with the definition of  $I_\nu$ , we see that

$$u_\nu = \frac{1}{c} \int I_\nu d\Omega = \frac{4\pi}{c} J_\nu. \quad (7.5)$$

For a blackbody,  $I_\nu = B_\nu$  doesn't depend on angle, and we can integrate over  $d\Omega$ ,

$$u_\nu = \frac{8\pi h\nu^3}{c^3} \left[ \exp\left(\frac{h\nu}{k_B T}\right) - 1 \right]^{-1}. \quad (7.6)$$

The total energy density can then be found by integrating over all frequencies, giving

$$u = \left[ \frac{8\pi^5 k_B^4}{15h^3 c^3} \right] T^4 \equiv aT^4$$

in agreement with what we derived from statistical mechanics, §6.4.

The next quantity to define is the *flux* of energy, along direction  $\hat{k}$ , per unit time  $dt$ , per unit area  $dA$ , and per frequency interval  $d\nu$ . We take the *first* moment of the specific intensity by multiplying  $I_\nu$  by a direction vector  $\hat{k}$  and integrating over  $d\Omega$ :

$$F_\nu = \int \hat{k} I_\nu d\Omega. \quad (7.7)$$

Note that  $F_\nu$  is a vector; the net flux along a direction  $\hat{n}$  is

$$F_\nu = \int I_\nu (\hat{n} \cdot \hat{k}) d\Omega.$$

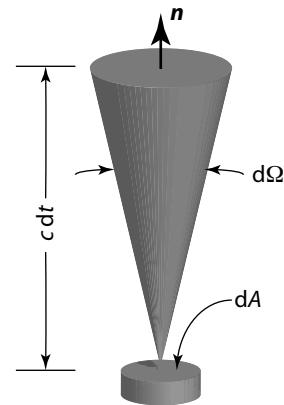


Figure 7.1: Schematic of a pencil of radiation propagating into an angle  $d\Omega$ .

If we take our polar angle with respect to  $\hat{k}$ , then  $(\hat{n} \cdot \hat{k}) d\Omega = \cos \theta \sin \theta d\theta d\phi$ ; defining the direction cosine  $\mu = \cos \theta$ , this becomes

$$F_\nu = \int I_\nu (\hat{n} \cdot \hat{k}) d\Omega = \int_0^{2\pi} \int_{-1}^1 I_\nu \mu d\mu d\phi.$$

Note that if the radiation field is isotropic then  $F_\nu = 0$ : there must be some anisotropy in the radiation field to generate a net flux.

For blackbody radiation, if we only integrate over outgoing directions,  $0 \leq \mu \leq 1$ , as would be the case for thermal radiation emerging from a *hohlraum*,

$$F_\nu = \pi B_\nu.$$

Integrating this  $F_\nu$  over all frequencies, we recover the Stefan-Boltzmann formula,

$$F = \left(\frac{ac}{4}\right) T^4 \equiv \sigma_{\text{SB}} T^4,$$

where  $\sigma_{\text{SB}}$  is the Stefan-Boltzmann constant.

Finally, let's look at the momentum flux along direction  $\hat{j}$  being transported along direction  $\hat{k}$ , per unit time  $dt$ , per unit area  $dA$ , and per unit frequency interval  $d\nu$ . Since for a photon,  $E = pc$ , we divide the energy flux by  $c$ . This is a *tensor*,

$$\mathbf{P}_\nu = \frac{1}{c} \int \hat{j} \hat{k} I_\nu d\Omega. \quad (7.8)$$

This is just the *second* angular moment of the specific intensity. The net (isotropic) momentum flux along a direction  $\hat{n}$  is then

$$P_\nu = \frac{1}{c} \int (\hat{n} \cdot \hat{k})(\hat{n} \cdot \hat{k}) I_\nu d\Omega = \frac{2\pi}{c} \int_{-1}^1 I_\nu \mu^2 d\mu.$$

For blackbody radiation,

$$P_\nu = \frac{4\pi}{3c} B_\nu = \frac{1}{3} u_\nu$$

and we may integrate this over frequency to obtain  $P = u/3$ , the standard result from thermodynamics.

## 7.2 Some simple estimates

We argued in the previous section that the solar interior is quite opaque. Naively, we might imagine some radiative transition, e.g. bremsstrahlung, emitting a photon. The photon speeds away at  $c$ , but it doesn't get very far before being absorbed or scattered by another particle. A new photon, either due to emission or scattering, will be emitted at some random direction, and the whole process repeats. This is just a description of a *random walk*.

For some simple estimate, let's assume that the hop is the same for all photons, regardless of frequency or ambient temperature. If the hop length is  $\ell$ , then we know that the total path length to get from the center to the surface is  $R_\odot(R_\odot/\ell)$  and the time for this to occur is  $R_\odot^2/\ell/c$ . What is a good estimate for  $\ell$ ? Consider a planar electromagnetic wave incident on a collection of scatterers. If these scatterers are uncorrelated, then the probability of scattering is just the number of scatterers times the probability for scattering from a single scatterer. Define the probability of scattering as

$$\mathcal{P} = N \times \left( \frac{\text{energy scattered per unit time by one scatterer}}{\text{energy incident per unit time per unit area}} \right) \frac{1}{\mathcal{A}} \quad (7.9)$$

where  $\mathcal{A}$  is the area normal to the propagation direction  $\hat{k}$  of the volume containing the  $N$  scatterers. The quantity in parenthesis is just the definition of the cross-section  $\sigma$ . Furthermore, if we set  $\mathcal{P} = 1$ , then the total number of scatters is just  $N = n \times \mathcal{A} \times \ell$ , where  $n$  is the number density of scatterers. Thus we define the *mean free path*,

$$\ell = \frac{1}{n\sigma}. \quad (7.10)$$

In stellar work, it is more convenient to use mass density rather than number density. Writing  $n = Y\rho/m_u$ , where  $Y$  is the abundance of scatterers, we have

$$\ell = \rho^{-1} \left( \frac{m_u}{Y\sigma} \right) \equiv (\rho\kappa)^{-1}$$

where  $\kappa$  is the *opacity* and has dimensions  $[\kappa] \sim [\text{cm}^2 \text{ g}^{-1}]$ .

The opacity in the stellar interior is set by a large number of processes (see §7.5): Thomson scattering, free-free absorption, atomic absorption, and photoionization. In general, the cross-section depends on the ambient temperature and density and the frequency of the photon. Over the length of a hop  $\ell$  the temperature and density will only vary slightly. As a result, the conditions are nearly isotropic, so we indeed expect the radiation to come into thermal equilibrium with the ambient material. But the conditions are not perfectly isotropic—otherwise there would be zero net heat flux! It is the small anisotropy that gives rise to the transport of energy. Let's imagine a small cube of material, with the size of this cube being  $\ell$ . Because we are so very nearly isotropic and in thermal equilibrium, the flux through any one face of this cube must be  $(c/6)u$ . Now suppose we have two adjacent cubes, with the common face of the cubes being at  $x = 0$ . The flux across the face has contributions from photons emitted at  $x - \ell$  and  $x + \ell$ , so the net flux is

$$\begin{aligned} F &\approx \frac{c}{6}u(x - \ell) - \frac{c}{6}u(x + \ell) \\ &\approx -\frac{1}{3}c\ell \frac{du}{dx}. \end{aligned} \quad (7.11)$$

This is a diffusion equation with coefficient  $c/(3\rho\kappa)$ . Our derivation is very crude, as it neglects the variation in cross section with the properties of the ambient medium and with the photon frequency. Nonetheless, this is basically the correct scenario; heat diffuses with a coefficient given by some suitably defined average over all sources of opacity.

### 7.3 Equation of Transfer

We're now ready to formalize the crude work in the previous section. In the absence of interactions with matter, the specific intensity  $I_\nu$  is conserved along a ray propagating in direction  $\hat{k}$ :  $dI_\nu/ds = c^{-1}\partial_t I_\nu + \hat{k} \cdot \nabla I_\nu = 0$ . If matter is present, it can do three things to change  $I_\nu$ .

*emit* Matter may spontaneously emit photons and add to the beam:  $dI_\nu/ds = \rho\varepsilon_\nu/(4\pi)$ . Here  $\varepsilon_\nu$  is the energy spontaneously emitted per unit frequency per unit time per unit mass. The factor of  $4\pi$  is to make this term per steradian.

*absorb* Photons have a chance of being absorbed or scattered out of the beam:  $dI_\nu/ds = -\rho\kappa_\nu I_\nu$ . Here the right-hand side is the energy removed from the beam along a path  $ds$  with  $\kappa_\nu = \kappa_\nu^{\text{abs}} + \kappa_\nu^{\text{sca}}$  being the total opacity (absorption plus scattering). The dimensions of opacity are clearly  $[\kappa_\nu] \sim [\text{cm}^2/\text{g}]$ . (If we had stimulated emission, this would be a *negative*  $\kappa_\nu$ .)

*scatter* Photons may be scattered into the beam from other directions:  $dI_\nu/ds = \rho\kappa_\nu^{\text{sca}}\phi_\nu$ . If the scattering is isotropic, then

$$\phi_\nu = \frac{1}{4\pi} \int_0^{2\pi} \int_0^\pi I_\nu \, d\phi \sin\theta \, d\theta \equiv J_\nu, \quad (7.12)$$

where  $J_\nu$  is the mean intensity: the scattering redistributes the energy over all angles.

Putting all these terms together gives us the *equation of transfer*,

$$\frac{1}{c} \partial_t I_\nu + \hat{k} \cdot \nabla I_\nu = \rho \frac{\varepsilon_\nu}{4\pi} - \rho\kappa_\nu I_\nu + \rho\kappa_\nu^{\text{sca}}\phi_\nu \quad (7.13)$$

for the specific intensity  $I_\nu$ .

#### Radiative equilibrium

The emissivity  $\varepsilon_\nu$  and the opacity  $\kappa_\nu$  describe how the radiation interacts with matter. A condition of steady-state is that the gas not gain or lose energy to the radiation. This requires balancing

$$(\text{energy emitted per unit volume}) = \rho \int \frac{\varepsilon_\nu}{4\pi} \, dv \, d\Omega$$

with

$$(\text{energy absorbed per unit volume}) = \rho \int \kappa_\nu^{\text{abs}} I_\nu d\nu d\Omega,$$

or

$$\int_0^\infty \left( \frac{\varepsilon_\nu}{4\pi} - \kappa_\nu^{\text{abs}} I_\nu \right) d\nu = 0. \quad (7.14)$$

Here we assume  $\varepsilon_\nu$  does not depend on angle. We don't include scattering because it doesn't transfer energy between the radiation and the gas.

Now suppose that the level populations of the matter are in thermal equilibrium and can be described by a temperature  $T$ . In that case, detailed balance must hold, so that

$$\frac{\varepsilon_\nu}{4\pi\kappa_\nu^{\text{abs}}} = B_\nu(T), \quad (7.15)$$

where  $B_\nu(T)$  is the Planck function. This defines *local thermodynamic equilibrium (LTE)*. If the radiation field is, in addition, described by a Planck function *at the same temperature* then we would have complete thermodynamic equilibrium.

### Optical depth

Consider a ray directed into a medium in steady-state ( $\partial_t \rightarrow 0$ ). In the absence of emission ( $\varepsilon_\nu = 0$ ) or scattering ( $\kappa_\nu^{\text{sca}} = 0$ ) equation (7.13) takes a particularly simple form:

$$\frac{dI_\nu}{d\tau_\nu} = -I_\nu. \quad (7.16)$$

Here we have set  $\hat{k} \cdot \nabla = (d/ds)$ , where  $ds$  is a infinitesimal along the path of the ray, and further have defined the *optical depth* as

$$\tau_\nu = \int \rho \kappa_\nu ds. \quad (7.17)$$

Note that  $\tau_\nu$  is dimensionless. Taking  $\rho$  and  $\kappa_\nu$  as given, equation (7.16) has a simple solution,

$$I_\nu(\tau_\nu) = I_\nu(0) \exp(-\tau_\nu).$$

Note that  $\rho\kappa = \ell^{-1}$ , so equation (7.17) is just  $\tau = \int ds/\ell$ , i.e., it expresses distance by counting the number of mean free pathlengths traversed.

### Source function

Having defined the optical depth, we can now add the emissivity  $\varepsilon_\nu$  and scattering term (henceforth we will assume isotropic scattering) to equation (7.16) to obtain

$$\frac{dI_\nu}{d\tau_\nu} = S_\nu - I_\nu, \quad (7.18)$$

where

$$S_\nu \equiv \frac{1}{\kappa_\nu} \left( \frac{\varepsilon_\nu}{4\pi} + \kappa_\nu^{\text{sca}} J_\nu \right) \quad (7.19)$$

is the *source function*. In the absence of scattering, so that  $S_\nu = \varepsilon_\nu/(4\pi\kappa_\nu)$  is a known function of  $\tau_\nu$ , we can formally solve equation (7.18):

$$I_\nu(\tau_\nu) = I_\nu(0) \exp(-\tau_\nu) + \int_0^{\tau_\nu} S_\nu(\tau_\nu) \exp(t - \tau_\nu) dt.$$

In the presence of scattering,  $S_\nu$  depends on  $J_\nu = (1/4\pi) \int I_\nu d\Omega$ , so that equation (7.18) is an *integro-differential* equation.

#### 7.4 Diffusion Approximation and the Rosseland Mean Opacity

At large optical depth, such as deep in a stellar interior, the radiation field is in thermal equilibrium, so that  $I_\nu = S_\nu = B_\nu$ . To see this, consider the relative scales of terms in the transfer equation.

$$\begin{array}{ccccc} \frac{1}{c} \frac{\partial I_\nu}{\partial t} & + & \hat{k} \cdot \nabla I_\nu & = & \rho \frac{\varepsilon_\nu}{4\pi} - \rho \kappa_\nu I_\nu + \rho \kappa_\nu^{\text{sca}} J_\nu \\ \text{I} & & \text{II} & & \text{III} \quad \text{IV} \quad \text{V} \end{array}$$

On the left-hand side, term I scales as  $I_\nu/(ct_\odot)$ , where  $t_\odot$  is the evolutionary timescale of the sun (Gyr), and term II scales as  $I_\nu/R_\odot$ . On the right-hand side, term IV scales at  $I_\nu/\ell$ . Hence the ratio of terms of term II to term IV is  $\ell/R_\odot \ll 1$  and that of term I to term IV is  $\ell/\text{Gpc} \ll 1$ . In addition, stellar properties change negligibly on scales of a mean-free path, so conditions are nearly isotropic over much of the interior and  $I_\nu = J_\nu$ . Hence  $I_\nu = J_\nu = S_\nu$ , and inserting the relation between  $\varepsilon_\nu$  and  $\kappa_\nu^{\text{abs}}$  from detailed balance, eq. (7.15), into equation (7.19) implies that  $S_\nu = B_\nu$ .

If the radiation field is perfectly isotropic there is no flux, however, so we must have some small anisotropy. Let's write  $I_\nu$  as a thermal term plus a correction,

$$I_\nu = B_\nu(T) + I_\nu^{(1)}.$$

Substituting this into the steady-state equation of transfer,

$$\frac{1}{\rho\kappa_\nu} \hat{k} \cdot \nabla I_\nu = S_\nu - I_\nu$$

and setting the term  $S_\nu - B_\nu = 0$  on the right-hand side, we obtain

$$I_\nu^{(1)} = -\frac{1}{\rho\kappa_\nu} \hat{k} \cdot \nabla B_\nu = -\frac{1}{\rho\kappa_\nu} \frac{\partial B_\nu}{\partial T} \hat{k} \cdot \nabla T. \quad (7.20)$$

This is anisotropic: the energy transport is largest in the direction “down” the temperature gradient. Let's get the net flux: multiply

equation (7.20) by  $\hat{k}$  to get the flux; and then take the component along a direction  $\hat{n}$  parallel to  $\nabla T$ ; finally replace the two dot products by the angle cosine  $\mu$ , and integrate over  $d\Omega = 2\pi d\mu$  to obtain

$$\mathbf{F}_v = -\frac{4\pi}{3} \frac{1}{\rho} \left[ \frac{1}{\kappa_v} \frac{\partial B_v}{\partial T} \right] \nabla T. \quad (7.21)$$

The quantity in [ ] deserves a closer look. First, suppose  $\kappa_v$  is independent of frequency. Then equation (7.21) means that the energy transport is greatest at the frequency where  $\partial B_v / \partial T$  is maximum, and *not* at the peak of the Planck spectrum.

**EXERCISE 7.1**— Explain on physical grounds why the flux for a grey opacity would be greatest at the frequency for which  $\partial B_v / \partial T$ , rather than  $B_v$ , is maximized.

Let us define the *Rosseland mean opacity* as

$$\kappa_R \equiv \left[ \frac{\int dv \kappa_v^{-1} (\partial B_v / \partial T)}{\int dv (\partial B_v / \partial T)} \right]^{-1}.$$

We can use this to integrate equation (7.21) to obtain the total radiative flux,

$$\mathbf{F} = -\frac{4\pi}{3} \frac{1}{\rho \kappa_R} \nabla \left[ \int dv B_v \right] = -\frac{1}{3} \frac{c}{\rho \kappa_R} \nabla a T^4. \quad (7.22)$$

This is just our formula for radiation diffusion (eq. [7.11]) that we obtained from physical arguments, but now we have an expression for the effective opacity  $\kappa_R$ .

## 7.5 Sources of Opacity

There are several processes that contribute to radiative opacity in stellar interiors. These are well described in standard texts, so we'll just briefly list them here.

### Thomson scattering

Thomson scattering is scattering from non-relativistic electrons when the photon energy is sufficiently low that we can neglect the recoil of the electron. The cross-section for Thomson scattering derived in Jackson and is

$$\sigma_{\text{Th}} = \frac{8\pi}{3} \left( \frac{e^2}{m_e c^2} \right)^2 = 0.665 \times 10^{-24} \text{ cm}^2. \quad (7.23)$$

The opacity for Thomson scattering is then

$$\kappa_{\text{Th}} = \frac{n_e \sigma_{\text{Th}}}{\rho} = (0.4 \text{ cm}^2 \text{ g}^{-1}) Y_e.$$

The factor of  $Y_e$  is because the electrons, which are much lighter than nuclei and therefore easier for an incident wave to shake, do the scattering.

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EXERCISE 7.2— Using Thomson scattering for the dominant opacity, estimate the photon diffusion time for the sun.

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### *Free-free absorption*

Another important one is free-free absorption. This is the inverse of *bremsstrahlung*, which is radiation emitted when an electron is scattered from an ion (see Figure 8.1). The procedure for calculating the opacity is to first compute the emissivity and then use detailed balance (eq. [7.15]) to obtain

$$\kappa_v^{\text{ff}} = \frac{\varepsilon_v}{4\pi B_v(T)}.$$

To calculate the emissivity, we start with the derivation of the momentum gained by an electron (eq. [8.1]). The acceleration leads to an emission of radiation; according to Larmor's formula, the power emitted is

$$P(b) = \frac{2}{3} \frac{e^2}{c^3} |\dot{v}|^2 = \frac{2}{3} \frac{Z^2 e^6}{m_e^2 c^3 b^4}. \quad (7.24)$$

The radiation is distributed over a broad range of frequencies up to a cutoff  $\nu_{\max} \sim v/b$ . Integrating over a range of impact parameters and then over the distribution of electron velocities gives the emissivity, which is (restoring all of the numerical factors)

$$\rho \varepsilon_v = 4\pi \left( \frac{2\pi}{3} \right)^{1/2} Z^2 n_I n_e \hbar c^2 \alpha \sigma_{\text{Th}} \left( \frac{m_e}{k_B T} \right)^{1/2} \exp \left( -\frac{hv}{k_B T} \right) \bar{g}_{\text{ff}}. \quad (7.25)$$

The velocity-averaged Gaunt factor  $\bar{g}_{\text{ff}}$  contains most of the details about the integration. The factor of  $T^{-1/2}$  is because there is a factor of  $v^{-1}$  that appears in the integration (the collision time is  $\sim b/v$ ).

Applying detailed balance, equation (7.15), gives the opacity as a function of frequency,

$$\kappa_v^{\text{ff}} = \pi \left( \frac{2\pi}{3} \right)^{1/2} Z^2 \frac{n_I n_e}{\rho} c^3 \alpha \sigma_{\text{Th}} \left( \frac{m_e c^2}{k_B T} \right)^{1/2} v^{-3} \left[ 1 - \exp \left( -\frac{hv}{k_B T} \right) \right] \bar{g}_{\text{ff}} \quad (7.26)$$

Notice that  $n_I n_e / \rho = \rho / (m_u^2 \mu_I \mu_e)$ , so that the opacity scales with density. Further, note that when taking the Rosseland mean over all frequencies, the factor of  $v^{-3}$  introduces a factor of  $T^{-3}$ , so that  $\langle \kappa_v^{\text{ff}} \rangle \propto \rho T^{-7/2}$ .

**EXERCISE 7.3—** In terms of central density and temperature, under what conditions is free-free opacity more important than Thomson scattering? For what mass range of stars is free-free opacity dominant in the core? What about for Thomson scattering?

### Bound-free and bound-bound absorption; Kramer's opacity

Bound-free and bound-bound transitions have a cross-section with a frequency dependence (at photon energies above threshold) of  $\nu^{-3}$  as well; therefore their Rosseland averages also scale as  $\rho T^{-7/2}$ . An opacities with this form is known as Kramer's opacity,  $\kappa = \kappa_0 \rho T^{-3.5}$ . For conditions in the solar center, a good approximation is

$$\kappa \approx 0.012 \text{ cm}^2 \text{ g}^{-1} \left( \frac{\rho}{1 \text{ g cm}^{-3}} \right) \left( \frac{T}{10^7 \text{ K}} \right)^{-7/2} \times \left[ (1 + X_{\text{H}}) \left( X_{\text{H}} + X_{\text{He}} + \sum_{Z>2} \frac{X_i Z_i^2}{A_i} \right) \right]. \quad (7.27)$$

The expression in [ ] is just an approximation for  $Y_e \langle Z^2 \rangle$  at solar composition.

The total opacity combining all of the above as used in MESA is shown in Figure 7.2.

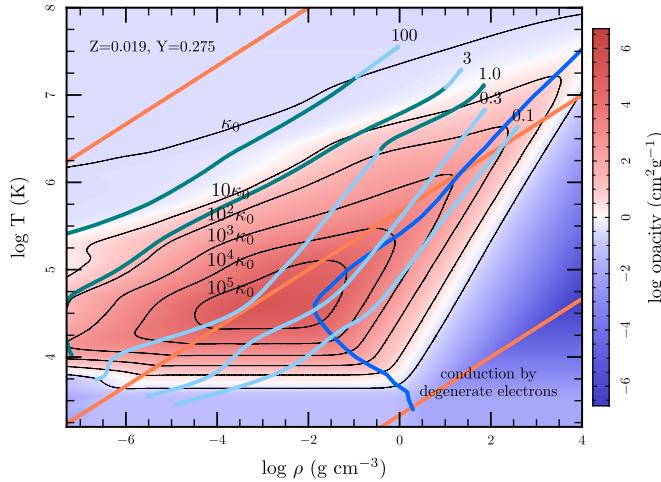


Figure 7.2: Total opacities used in MESA for solar composition. The underlying shades show the value of  $\kappa$ , whereas the contours are in units of the electron scattering opacity,  $\kappa_{\text{Th}} = 0.2(1 + X) \text{ cm}^2 \text{ g}^{-1}$ . Stellar interior profiles for main sequence stars of mass  $M = 0.1, 0.3, 1.0, 3.0$ , and  $100 M_{\odot}$  are shown by the green (radiative regions)–slight blue (convective regions) lines. Electron conduction dominates the opacity to the right of the dark blue line (which is where the radiative opacity equals the conductive opacity). From Paxton et al. [2011].

**EXERCISE 7.4—** Suppose both free-free absorption ( $\kappa_{\nu}^{\text{ff}}$ ) and Thomson scattering ( $\kappa_{\nu}^{\text{Th}}$ ) contribute to the opacity. Denote by  $\langle \cdot \rangle$  the Rosseland averaging of an opacity. Does  $\langle \kappa_{\nu}^{\text{ff}} + \kappa_{\nu}^{\text{Th}} \rangle = \langle \kappa_{\nu}^{\text{ff}} \rangle + \langle \kappa_{\nu}^{\text{Th}} \rangle$ ?

Now suppose that  $\kappa_{\nu}^{\text{ff}}$  is due to free-free absorption on two different ion species (denoted below by subscripts “1” and “2”) with different charge number  $Z$ . In this case does  $\langle \kappa_{\nu,1}^{\text{ff}} + \kappa_{\nu,2}^{\text{ff}} \rangle = \langle \kappa_{\nu,1}^{\text{ff}} \rangle + \langle \kappa_{\nu,2}^{\text{ff}} \rangle$ ?

## 7.6 Eddington Standard Model

Polytropes with index  $n = 3/2$  correspond to fully convective stars ( $P \propto \rho^{5/3}$ , the relation for an adiabat) or for white dwarfs (non-relativistic, degenerate equation of state). Another interesting case, for historical reasons, is the *Eddington Standard Model*, which is a fair approximation to main-sequence stars with  $M \gtrsim M_\odot$ . Suppose we write the equation of state as the sum of ideal gas and radiation pressure,

$$P = \frac{\rho k_B T}{\mu m_u} + \frac{1}{3} a T^4. \quad (7.28)$$

Now make the *ansatz* that

$$\frac{P_{\text{rad}}}{P} = \frac{a T^4}{3P} = 1 - \beta = \text{const.}, \quad (7.29)$$

that is, the radiation pressure is a fixed fraction of the total pressure everywhere. Solving for  $T$  in terms of  $P$  and  $\beta$ ,

$$T = \left[ \frac{3(1-\beta)P}{a} \right]^{1/4},$$

and inserting this into equation (7.28) gives us a simple EOS,

$$P = \left[ \left( \frac{k_B}{\mu m_u} \right)^4 \frac{3}{a} \right]^{1/3} \left[ \frac{1-\beta}{\beta^4} \right]^{1/3} \rho^{4/3} \equiv K \rho^{4/3}. \quad (7.30)$$

This is the equation for a polytrope of index 3.

Using the numerical solutions for an  $n = 3$  polytrope we found in 4.2, we can construct a simple Eddington solar model. In Figure 7.3 such a model is shown and compared to a “Standard Solar Model” as one might construct using, e.g., MESA. The comparison is surprisingly good, given the assumptions of the model. Things go off-the-rails a bit at the edge of the model, but this is to be expected as the assumption of the constant  $\beta$  *must* breakdown there. Furthermore, the Eddington model can be used to derive several important relations for main sequence stars (see below and the Exercises). Incredibly, the Eddington model makes *no* assumptions about the source of the energy generation in the star! Only that the star is in hydrostatic and radiative equilibrium, which is enough to tightly couple the emergent luminosity with the internal luminosity.

**EXERCISE 7.5** — Derive an expression for  $\beta$  in terms of the mass of the star for the Eddington Standard Model.

Why is it at all reasonable to take  $\beta$  as being constant? To explore this, go back to the equation for radiative diffusion

$$F(r) = -\frac{1}{3} \frac{c}{\rho \kappa} \frac{da T^4}{dr}.$$

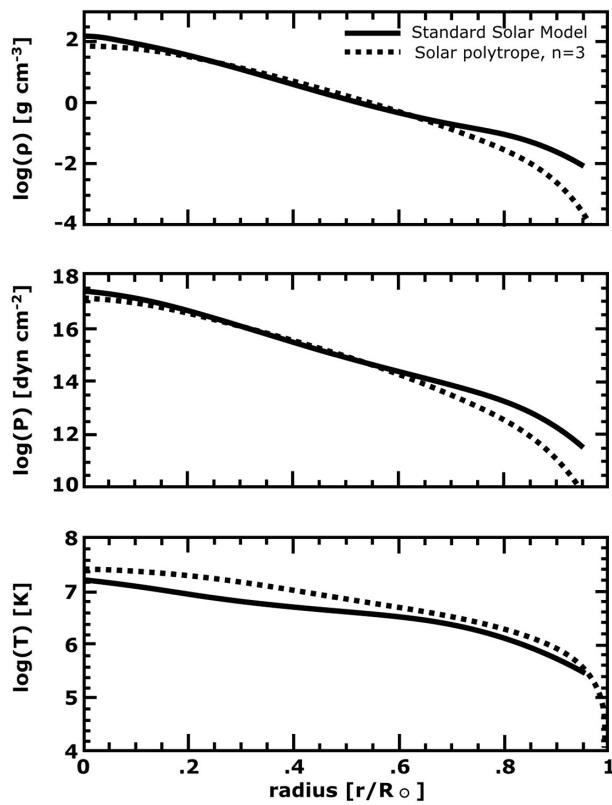


Figure 7.3: The Eddington Model with parameters appropriate for the Sun compared to the Standard Solar Model. (From Lamers & Levesque.)

Write the flux as  $F(r) = L(r)/(4\pi r^2)$ , and since pressure decreases monotonically with radius, write

$$\frac{daT^4}{dr} = \frac{daT^4}{dP} \frac{dP}{dr} = -\rho \frac{Gm(r)}{r^2} \frac{daT^4}{dP}.$$

The equation of radiation transport then becomes

$$L(r) = \frac{4\pi Gm(r)c}{\kappa(r)} \frac{dP_{\text{rad}}}{dP}.$$

Dividing both sides by  $L \cdot M / \kappa_{\text{Th}}$  and rearranging terms,

$$\frac{dP_{\text{rad}}}{dP} = \left[ \frac{L\kappa_{\text{Th}}}{4\pi GMc} \right] \left( \frac{\kappa(r)}{\kappa_{\text{Th}}} \frac{L(r)}{L} \frac{M}{m(r)} \right). \quad (7.31)$$

Here  $L$  is the total luminosity of the star and  $M$  is the total mass. The term in  $[ ]$  is a constant (the Thomson opacity  $\kappa_{\text{Th}}$  doesn't depend on density or temperature) and we define the *Eddington luminosity* as  $L_{\text{Edd}} = 4\pi GMc/\kappa_{\text{Th}}$ . For the sun,  $L_{\text{Edd}} = 1.5 \times 10^{38} \text{ erg s}^{-1} = 3.8 \times 10^4 L_{\odot}$ . For the term  $( )$  on the right-hand side, note the  $L(r)/m(r)$  is basically the average energy generation rate interior to a radius  $r$ . Since nuclear reactions are temperature sensitive, the heating is concentrated toward the stellar center and  $L(r)/m(r)$  decreases with radius. For stars like the sun, free-free opacity is dominant, and since the free-free Rosseland opacity goes as  $T^{-3.5}$ ,  $\kappa(r)$  increases with radius. Thus, if the energy generation rate is not too temperature dependent (the reaction  $p + p \rightarrow {}^2\text{H}$  goes roughly as  $T^{4.5}$  at  $T = 10^7 \text{ K}$ ), then the term in  $( )$  does not vary strongly with radius, and  $dP_{\text{rad}}/dP$  is indeed roughly constant.

**EXERCISE 7.6—** You are now in a position to understand why the luminosity depends strongly on the mass. Cast the flux equation (eq. [7.22]) into dimensionless form. Assume the opacity has the functional form  $\kappa = \kappa_0 \rho^a T^{-b}$ , and scale  $\rho$  and  $T$  in terms of  $M$  and  $R$ .

1. You should be able to find a characteristic scale for the luminosity which depends on the stellar mass  $M$  and radius  $R$ , as well as on the exponents  $a$  and  $b$ . Regard  $\kappa_0$  as a fitting constant, and adjust it so that you get an expression in the form

$$\frac{L}{L_{\odot}} = \left( \frac{M}{M_{\odot}} \right)^{\alpha} \left( \frac{R}{R_{\odot}} \right)^{\beta}.$$

2. If the opacity is dominated by Thomson scattering, what are  $\alpha$  and  $\beta$ ? What about if the opacity is Kramer's (eq. [7.27])?

**MESA** *Radiation pressure and the Eddington luminosity for massive stars*

Construct zero-age main-sequence (ZAMS) stars of masses  $1.0 M_{\odot}$ ,  $3.0 M_{\odot}$ ,  $10.0 M_{\odot}$ , and  $30.0 M_{\odot}$ . You will find the template for the project in the folder `radiation/beta-eddington`.

1. For each star, plot  $\beta \equiv P_{\text{gas}}/P$  as a function of Lagrangian mass coordinate  $m$ . Is  $\beta$  roughly constant, i.e., independent of  $m$ ? For each ZAMS model, assign a “typical” value of  $\beta$  and plot this  $\beta$  as a function of the total stellar mass  $M$ . How well does  $\beta(M)$  agree with what you derived in the warm-up exercise?
2. For each star, plot  $L_{\text{rad}}/L_{\text{Edd}}$  as a function of  $m$ .

The template project files are set up to load a file ‘`plot_radn_variables.inlist`’, which you will write. The file should contain a customized version of ‘`Profile_Panels1`’ that displays the  $\beta$  and  $L_{\text{rad}}/L_{\text{Edd}}$ .

Finally, you may notice that things get interesting near the surface of the star, especially for the more massive stars. Change the independent variable from ‘`mass`’ to ‘`logxq`’ and redo the plots. (You may need to adjust the minimum value of the x-axis and reverse the direction of the x-axis.) Comment on the results.

# 8

## *Transport in a Plasma*

### 8.1 Collisions

Without collisions, a plasma cannot reach thermodynamic equilibrium, and the rate of collisions mediates both the approach to equilibrium and the transport of quantities, such as heat, in a forced system. In this section, we'll make an estimate for the rate of electron-electron ion-ion, and electron-ion collisions.

To begin, let's imagine a light particle (electron) colliding with a much heavier, fixed particle (an ion), as illustrated in Figure 8.1. (This picture also applies to a pseudo particle of reduced mass scattering in a fixed potential.) Let the impact parameter be  $b$ , and the mass of the incident particle is  $\mu$ . For Coulomb interactions, the force on the particle is  $(q_1 q_2 / r^2) \hat{r}$ . The incident momentum is  $p_0$ . Now by assumption, in our plasma most of the interactions are weak (potential energy is much less than kinetic), so let's treat the deflection of the particle as a perturbation. That is, we shall assume that  $p_0 = \text{const}$  and that the effect of the interaction is to produce a perpendicular (to  $p_0$ ) component of the momentum  $p_\perp$ . The total change in  $p_\perp$  is then

$$p_\perp = \int_{-\infty}^{\infty} dt \frac{q_1 q_2}{r^2} \sin \theta,$$

where  $\sin \theta = b/r$  is the angle that the radial vector makes with the horizontal. Substituting  $r = b/\sin \theta$  and  $dt = -\mu b d\theta / p_0 / \sin^2 \theta$ , we have

$$p_\perp = - \int_0^\pi \sin \theta d\theta \frac{\mu}{p_0} \frac{q_1 q_2}{b},$$

leading to the intuitive result

$$\frac{p_0 p_\perp}{2\mu} = \frac{q_1 q_2}{b}. \quad (8.1)$$

Clearly a large angle scattering occurs if  $p_\perp \geq p_0$ , or

$$b \leq b_0 \equiv \frac{2\mu q_1 q_2}{p_0^2}; \quad (8.2)$$

our approach is only valid for  $b \gg b_0$ . Note that  $p_\perp/p_0 = b_0/b$ . What is the rate of large angle scatterings? The cross section for a large angle scattering is  $\sigma_{\text{LA}} = \pi b_0^2$ . Imagine a particle incident on a cylinder of length  $(p_0/\mu)dt$  and cross-sectional area  $\mathcal{A}$ . Within this cylinder there are  $n \times (p_0/\mu)dt\mathcal{A}$  scatterers of cross-section  $\sigma_{\text{LA}}$ . so the probability of the particle interacting per time  $dt$  is

$$\frac{(\sigma_{\text{LA}} \times n \times p_0/\mu)\mathcal{A}}{\mathcal{A}} = n\sigma_{\text{LA}}p_0/\mu.$$

This defines the large-angle collision rate,

$$\nu_{\text{LA}} = n\sigma_{\text{LA}}v_0 = \frac{4\pi\mu(q_1q_2)^2}{p_0^3}. \quad (8.3)$$

Note that it goes as  $p_0^{-3}$ ; fast-moving particles are hard to scatter.

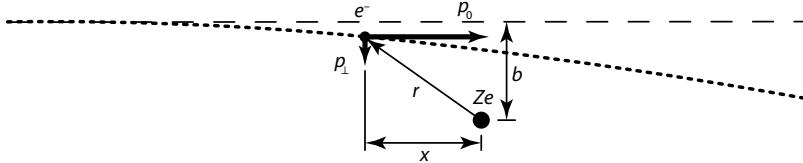


Figure 8.1: Geometry for scattering problem.

As mentioned earlier, we are in a weakly coupled plasma, so we expect large angle scatterings to be a rare occurrence. What happens instead is that the particle suffers a number of small deflections. Let's consider an impact parameter in the range  $(b, b + db)$ . Each deflection will have  $\Delta p_\perp$  in some random direction, so

$$\langle p_\perp \rangle = \sum_{i=1}^N \Delta p_i \approx \mathbf{0}.$$

What is happening is that the component of momentum perpendicular to  $p_0$  is executing a random walk. Indeed, after  $N$  collisions,

$$\begin{aligned} \langle p_\perp \cdot p_\perp \rangle &= \left( \sum_{i=1}^N \Delta p_i \right) \cdot \left( \sum_{i=1}^N \Delta p_i \right) \\ &= \sum_{i=1}^N (\Delta p_i)^2 + 2 \sum_{i \neq j} \Delta p_i \cdot \Delta p_j = N(\Delta p_\perp)^2, \end{aligned}$$

where we have assumed that all  $\Delta p_\perp$  have the same magnitude and are uncorrelated. Now  $N$  is just  $\Delta t \times n \times (2\pi b db) \times (p_0/\mu)$ : the number of particles with impact parameters between  $b$  and  $b + db$  along the length of the particles path over a time  $\Delta t$ . Dividing by  $\Delta t$  and integrating over  $b$ , we have the rate of change of the perpendicular component of the momentum

$$\frac{d\langle p_\perp^2 \rangle}{dt} = \frac{8\pi n\mu(q_1q_2)^2}{p_0} \int_{b_{\min}}^{b_{\max}} \frac{db}{b}. \quad (8.4)$$

What are  $b_{\max}$  and  $b_{\min}$ , the maximum and minimum impact parameters? Clearly, if  $b > \lambda_D$  then the potential will be screened. Since our approximation is only good for  $b > b_0$ , we may take  $b_{\min} = b_0$ . (Our rate of scattering only depends logarithmically on  $b_{\max}/b_{\min}$ , so these estimates are good enough for our purposes). With this substitution,

$$\begin{aligned}\frac{d\langle p_{\perp}^2 \rangle}{dt} &= \frac{8\pi n \mu (q_1 q_2)^2}{p_0} \ln \left( \frac{\lambda_D}{b_0} \right) \\ &\equiv \frac{8\pi n \mu (q_1 q_2)^2}{p_0} \ln \Lambda.\end{aligned}\quad (8.5)$$

In the literature, the quantity  $\ln \Lambda$  is called the Coulomb logarithm; for a plasma such as we are considering it is  $\sim \ln(\lambda_D^3 n)$ , the logarithm of the number of particles in a Debye sphere (see eq. [6.43]). For conditions typical of the solar center (hydrogen plasma,  $\rho \gtrsim 1 \text{ g cm}^{-3}$ ,  $T \approx 10^7 \text{ K}$ ),  $\ln \Lambda \approx (5-10)$ .

For small-angle scattering, the concept of a collision rate is fuzzy: the particle is constantly being bombarded by many tiny collisions. Setting  $d\langle p_{\perp}^2 \rangle/dt = p_0^2 \nu$  allows us to define a deflection rate,

$$\nu \approx \frac{8\pi n \mu (q_1 q_2)^2}{p_0^3} \ln \Lambda \quad (8.6)$$

$$= \frac{8\pi n (q_1 q_2)^2}{(3k_B T)^{3/2} \mu^{1/2}} \ln \Lambda. \quad (8.7)$$

Comparing equations (8.6) and (8.3), we see that many small angle scatterings are more important than single large angle scattering.

Note that the ion-ion collision rate will be about  $\sqrt{m_p/m_e} \approx 43$  times less than the electron-electron collision rate for a given temperature.

One can define a mean free path  $\ell$  from equation (8.6). Consider a particle incident on a cylinder of cross-sectional area  $\mathcal{A}$  and length  $\ell$ , as illustrated in Figure 8.2. We chose  $\ell$  so that the time for the particle to traverse it is  $\nu^{-1}$ , the timescale for deflection. Thus  $\ell = v_0/\nu = p_0/(\mu\nu)$ . Note that for a large angle collision, we can write the probability for scattering as the total cross-section of scatterers per unit area,

$$\mathcal{P} = \frac{N\sigma}{\mathcal{A}} = \frac{n \times (\ell\mathcal{A})\sigma}{\mathcal{A}}, \quad (8.8)$$

so the particle will suffer on average a collision after traversing a distance  $\ell = (n\sigma)^{-1}$ . Comparing equation (8.8) with our expression for  $\ell$  in terms of  $\nu$  allows us to define an effective cross-section for small angle scattering.

## 8.2 Transport coefficients

We now have enough machinery to make estimates of *transport coefficients*, such as the viscosity and the thermal conductivity. Let's begin

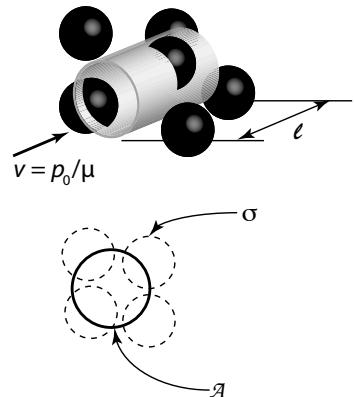


Figure 8.2: Schematic of a particle incident on a cylinder containing  $n \times \ell \times \mathcal{A}$  particles.

with the viscosity. Suppose we have a fluid with a gradient in the velocity, a shear, as depicted in Figure 8.3. Let the mean thermal velocity of a particle be  $v_0$ . In a time  $\Delta t$ , a number of particles will enter the box from the top,  $(1/6)nv_0\Delta t$ , and a similar number will leave the box via the top face. On average, these particles are endowed with the fluid properties of their last scattering, so the *net* momentum carried into the box across the top face is

$$\frac{1}{6}nmv_0\Delta t\Delta A [v_y(z_t + \ell) - v_y(z_t - \ell)] \approx \frac{1}{3}nmv_0\Delta t\Delta A \left. \frac{\partial v_y}{\partial z} \right|_{z_t} \ell. \quad (8.9)$$

Here  $\Delta A$  is the cross-section area of our box in the  $xy$  plane and  $z_t$  is the coordinate of the top face.

A similar process occurs across the bottom face, located at coordinate  $z = z_b$ : the momentum flux across the bottom face is

$$\approx -\frac{1}{3}nmv_0\Delta t\Delta A \left. \frac{\partial v_y}{\partial z} \right|_{z_b} \ell. \quad (8.10)$$

Note the difference in sign: the momentum flux is positive if the  $y$ -velocity is larger below the box. Putting equations (8.9) and (8.10) together, the net change of momentum per time per unit volume  $\Delta A\Delta z$  is

$$\begin{aligned} \frac{m}{\Delta A\Delta z} \frac{\Delta v_y}{\Delta t} &\approx \frac{1}{\Delta z} \frac{1}{3} \left[ \left( nmv_0\ell \frac{\partial v_y}{\partial z} \right)_{z_t} - \left( nmv_0\ell \frac{\partial v_y}{\partial z} \right)_{z_b} \right] \\ &\approx \frac{\partial}{\partial z} \left( \mu \frac{\partial v_y}{\partial z} \right). \end{aligned} \quad (8.11)$$

Here we have defined  $\mu = nmv_0\ell/3$  as the coefficient of dynamic viscosity.

On the left-hand side, the quantity  $m/(\Delta A\Delta z)$  is just the mass density  $\rho$ , and equation (8.11), so the left-hand side is the  $y$ -component of our old friend  $\rho(\partial_t \mathbf{u} + \mathbf{u} \cdot \nabla \mathbf{u})$ . On the right-hand side, we can repeat the above derivation for the transport of momentum in the  $x$ - and  $y$ -directions; if we also add back in forces from gravity and pressure gradients, we transform Euler's equation, eq. (2.5), to the *Navier-Stokes equation*,

$$\partial_t \mathbf{u} + \mathbf{u} \cdot \nabla \mathbf{u} = -\nabla \Phi - \frac{1}{\rho} \nabla P + \frac{1}{\rho} \nabla \cdot (\mu \nabla \mathbf{u}). \quad (8.12)$$

In an isothermal, incompressible fluid, one can pull  $\mu$  outside the divergence operator and the last term becomes

$$\frac{\mu}{\rho} \nabla^2 \mathbf{u} \equiv \nu \nabla^2 \mathbf{u},$$

where  $\nu$  is defined as the *coefficient of kinematic viscosity* (sorry for the overload of notation with the scattering frequency earlier!). Note that

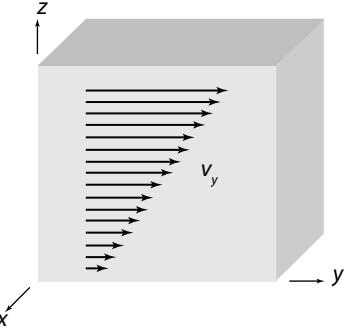


Figure 8.3: An element of fluid with a shear  $\partial v_y / \partial z$ .

in order-of-magnitude

$$\nu \sim \frac{1}{3} v_0 \ell,$$

that is, it is roughly the thermal velocity times the mean free path.

**EXERCISE 8.1**— Estimate the ratio of the pressure to viscous accelerations,

$$\frac{|\rho^{-1} \nabla P|}{|\nu \nabla^2 u|},$$

in equation (8.12). Express your answer in terms of a characteristic lengthscale, Mach number, and mean free path. Under what conditions are viscous effects important?

An identical procedure, but replacing the average momentum of a particle with its average thermal energy, yields an expression for the *thermal conductivity*  $K$ , such that the heat flux is

$$\mathbf{F} = -K \nabla T. \quad (8.13)$$

If one writes the change in energy of a fluid element as being

$$\rho C \partial_t T = \nabla \cdot (K \nabla T),$$

it can be seen that in order of magnitude the thermal diffusivity  $\chi \equiv K/(\rho C)$ , where  $C$  is the specific heat per unit mass, is  $\chi \sim (1/3)v_0\ell$ . (From the form of the equation and dimensional analysis, it has to be like this.) But we need to be careful here: in a plasma with ions and electrons, the ions are responsible for momentum transport, whereas electrons, being more nimble, are more effective at heat transport. Thus the thermal diffusivity is larger than the kinematic viscosity by a factor  $\sim \sqrt{m_p/m_e} \approx 43$ .

**EXERCISE 8.2**— Estimate the plasma thermal conductivity under conditions appropriate to the solar center. How does heat conduction by the electrons compare to that by photons?

Now suppose we wish to write a single equation for heat transport,

$$\mathbf{F} = -\frac{4}{3} \frac{acT^3}{\rho \kappa_{\text{total}}} \nabla T.$$

Derive an expression for  $\kappa_{\text{total}}$  in terms of the electron conductivity  $K$  and the free-free opacity  $\kappa_{\text{ff}}$ .



# 9

## *Convection*

Hot air rises, as a glider pilot or hawk can tell you. The fluid velocities in question are very subsonic, so we have hydrostatic equilibrium to excellent approximation. But the fluid motions make an enormous difference for heat transport! This state of fluid motions induced by a temperature gradient is known as *convection*. You can perform the following demonstration of the onset of convection. Brew tea, and pour the hot tea into a saucepan that is on an unlit burner. Use a straw with your thumb over the top to insert a layer of cold milk under the warm tea in the saucepan. The temperature difference between the tea and milk will inhibit their mixing. Light the burner, and watch for the development of convection—you will know it when you see it.



Figure 9.1: Onset of convection in a tea-milk mixture.

### *9.1 Criteria for onset of convection*

To understand this process, let's consider a fluid in planar geometry and hydrostatic equilibrium,

$$\frac{dP}{dr} = -\rho g. \quad (9.1)$$

Now, imagine moving a blob of fluid upwards from  $r$  to  $r + h$ . We move the blob slowly enough that it is in hydrostatic equilibrium with its new surroundings,  $P_b(r + h) = P(r + h)$ , where the subscript  $b$  refers to “blob.” We do move the blob quickly enough, however, that it does not remain in *thermal* equilibrium with its surroundings; that is, we move the blob *adiabatically*. The entropy of the blob is therefore constant,  $S_b(r + h) = S_b(r) = S(r)$ , and is therefore not, in general, equal to the entropy of the surrounding gas at  $r + h$ :  $S_b(r + h) \neq S(r + h)$ .

As the blob rises, it displaces some of the surrounding fluid. Archimedes tells us that if the displaced fluid is less massive than the blob, then the blob will sink. We can rephrase this in terms of the volume occupied by a unit mass of fluid  $V$ : if the volume occupied by the blob is less than the volume of an equal mass of background, then the blob will sink. Translating this into an equation: if

$$\begin{aligned} V[P(r + h), S(r + h)] - V_b[P_b(r + h), S_b(r + h)] &= \\ V[P(r + h), S(r + h)] - V[P(r + h), S(r)] &> 0 \end{aligned} \quad (9.2)$$

then the blob will sink. If condition (9.2) is violated, the blob will continue to rise, and the system is unstable to convection. Figure 9.2 has a cartoon of this process.

Taking  $h$  to be an infinitesimal displacement and expanding the left-hand side of equation (9.2) gives us a local condition for stability:

$$V[P(r + h), S(r)] + \left( \frac{\partial V}{\partial S} \right)_P \frac{dS}{dr} - V[P(r + h), S(r)] = \left( \frac{\partial V}{\partial S} \right)_P \frac{dS}{dr} > 0. \quad (9.3)$$

Noting that

$$\begin{aligned} \left( \frac{\partial V}{\partial T} \right)_P &= \left( \frac{\partial V}{\partial S} \right)_P \left( \frac{\partial S}{\partial T} \right)_P \\ &= \frac{C_P}{T} \left( \frac{\partial V}{\partial S} \right)_P, \end{aligned}$$

we can rewrite equation (9.3) as

$$\frac{T}{C_P} \left( \frac{\partial V}{\partial T} \right)_P \frac{dS}{dr} > 0.$$

Now,  $(\partial V / \partial T)_P$  is positive (gas expands on being heated), so our condition for stability is simply

$$\frac{dS}{dr} > 0. \quad (9.4)$$

In a convectively stable star, the entropy must increase with radius. if  $dS/dr < 0$ , then convection occurs and carries high-entropy

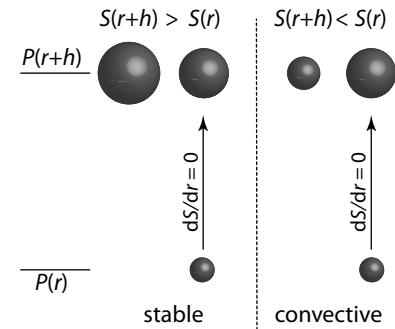


Figure 9.2: Illustration of criteria for convective instability. On the left, raising a blob a distance  $h$  adiabatically and in pressure balance with its surrounding results in a higher density  $V_b < V$ . This is stable: the blob will sink back. On the right, the blob is less dense and hence buoyant: it will continue to rise.

material outward, where it will eventually mix with the ambient medium. As a result, convection drives the entropy gradient toward the marginally stable configuration  $dS/dr = 0$ . If a star is fully convective and mixes efficiently, then the interior of the star lies along an adiabat.

WE CAN DERIVE A CONDITION FOR CONVECTIVE STABILITY in terms of the local gradients of temperature and pressure. Writing  $S = S[P(r), T(r)]$  we expand equation (9.4) to obtain

$$\frac{dS}{dr} = \left( \frac{\partial S}{\partial P} \right)_T \frac{dP}{dr} + \left( \frac{\partial S}{\partial T} \right)_P \frac{dT}{dr}. \quad (9.5)$$

Now,  $P$  is a monotonically decreasing function of  $r$ , which means we can use it as a spatial coordinate and write,

$$\frac{dT}{dr} = \frac{dT}{dP} \Big|_* \frac{dP}{dr}. \quad (9.6)$$

Here  $dT/dP|_*$  is the slope of the  $T(P)$  relation for the stellar interior. In particular, this is *not* a thermodynamic equality. Substituting equation (9.6) into equation (9.5), using hydrostatic equilibrium to eliminate  $dP/dr$ , and recognizing that  $(\partial S/\partial T)_P = C_P/T$ , we obtain

$$\frac{dS}{dr} = -\rho g \left[ \left( \frac{\partial S}{\partial P} \right)_T + \frac{C_P}{T} \frac{dT}{dP} \Big|_* \right]. \quad (9.7)$$

Finally, we can use the identity (see Appendix A.1)

$$\left( \frac{\partial S}{\partial P} \right)_T \left( \frac{\partial T}{\partial S} \right)_P \left( \frac{\partial P}{\partial T} \right)_S = -1 \quad (9.8)$$

to simplify equation (9.7),

$$\begin{aligned} \frac{dS}{dr} &= -\frac{\rho g}{P} C_P \left[ \frac{P}{T} \frac{dT}{dP} \Big|_* - \frac{P}{T} \left( \frac{\partial T}{\partial P} \right)_S \right] \\ &= -\frac{\rho g}{P} C_P [\nabla - \nabla_{ad}]. \end{aligned} \quad (9.9)$$

Here we have introduced the shorthand notation  $\nabla \equiv d \ln T / d \ln P|_*$  and  $\nabla_{ad} \equiv (\partial \ln T / \partial \ln P)_S$ . A mixture of uniform composition is unstable to convection if the local temperature gradient is steeper than an adiabat, i.e., if  $\nabla > \nabla_{ad}$ .

EXERCISE 9.1 — Assuming that  $\nabla \approx \nabla_{ad}$  in a convective region, sketch a plot of temperature as a function of pressure for the following cases.

1. A star with a stable inner layer and a convective outer layer;
2. A star with a convective inner layer and a stable outer layer.

Indicate on both of these plots an adiabat.

## 9.2 A second look at convective instability

Here we'll take a second-look at the convection by imagining we have a background state with velocity  $\mathbf{u} = 0$ ; we then *perturb* this state by displacing fluid elements a distance  $\delta r$ , and obtaining an equation of motion for  $\delta \dot{r}$ . This requires a bit of careful thought on what we are perturbing.

There are two types of perturbations. We may change a fluid quantity  $f$  at a fixed location  $\mathbf{r}$  and time  $t$  (Fig. 9.3):

$$\Delta f \equiv f(\mathbf{r}, t) - f_0(\mathbf{r}, t), \quad (9.10)$$

where the subscript “o” denotes the unperturbed quantity. We call  $\Delta f$  an *Eulerian perturbation*.

We may also change a fluid quantity  $f$  for a given fluid element; the position of this fluid element in the perturbed system is not necessarily at the same position as in the unperturbed case, however (Fig. 9.4):

$$\delta f \equiv f(\mathbf{r}, t) - f_0(\mathbf{r}_0, t). \quad (9.11)$$

We call  $\delta f$  a *Lagrangian perturbation*.

Since the fluid element is displaced  $\delta \mathbf{r} = \mathbf{r} - \mathbf{r}_0$ , we can add and subtract  $f_0(\mathbf{r}, t)$  to eq. (9.11) and expand  $f_0(\mathbf{r}, t)$  to first order in  $\delta \mathbf{r}$  to obtain a relation between the two types of perturbations:

$$\delta f = \Delta f + (\delta \mathbf{r} \cdot \nabla) f_0. \quad (9.12)$$

There are a few useful commutation relations that are easily proved:

$$\partial_t \Delta f = \Delta (\partial_t f), \quad (9.13)$$

$$\nabla \Delta f = \Delta \nabla f, \quad (9.14)$$

$$\frac{D}{Dt} \delta f = \delta \frac{Df}{Dt}. \quad (9.15)$$

And there are operations that do not commute:

$$\partial_t \delta f \neq \delta (\partial_t f), \quad (9.16)$$

$$\nabla \delta f \neq \delta \nabla f, \quad (9.17)$$

$$\frac{D}{Dt} \Delta f \neq \Delta \frac{Df}{Dt}. \quad (9.18)$$

One can further show that  $\delta \mathbf{u} = (D/Dt)\delta \mathbf{r}$ . Finally if the fluid has unperturbed velocity  $\mathbf{u} = 0$ , then  $\Delta \mathbf{u} = \delta \mathbf{u}$ .

ARMED WITH THESE RELATIONS, LET US PERTURB THE MOMENTUM EQUATION by adiabatically displacing a fluid element a distance  $\delta \mathbf{r}$ . We will do this in such a way that the pressure at a fixed location

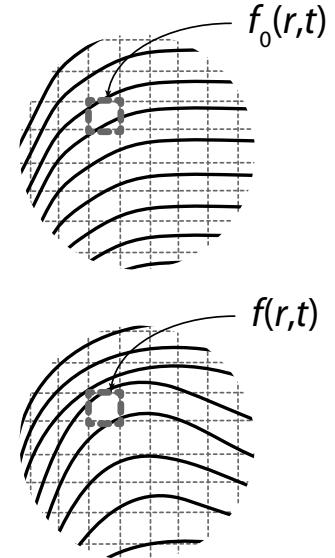


Figure 9.3: An Eulerian perturbation: we compare fluid quantities at corresponding locations.

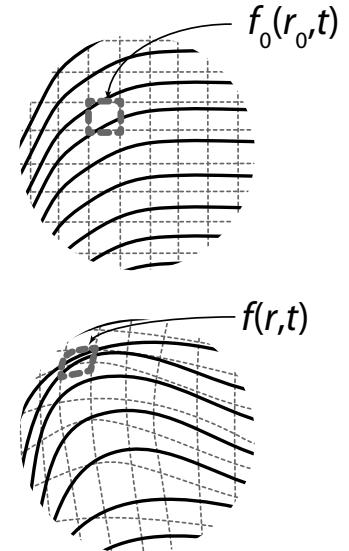


Figure 9.4: A Lagrangian perturbation: we compare fluid quantities for corresponding fluid elements.

does not change, i.e.,  $\Delta P = 0$ . Of course, the pressure and density of a given fluid element will change according to the relation

$$\frac{\delta P}{P} = \Gamma_1 \frac{\delta \rho}{\rho}$$

with  $\Gamma_1 \equiv (\partial \ln P / \partial \ln \rho)_s$ . We'll also assume that the gravitational force does not change,  $\Delta g = 0$ . Our perturbed momentum equation then becomes

$$\frac{D^2 \delta \mathbf{r}}{Dt^2} = -\frac{1}{\rho + \Delta \rho} \nabla P + \mathbf{g}.$$

Since in the unperturbed fluid  $\nabla P = \rho \mathbf{g}$ , this equation simplifies to

$$\frac{D^2 \delta \mathbf{r}}{Dt^2} = \frac{\Delta \rho}{\rho} \mathbf{g}. \quad (9.19)$$

Expanding,

$$\begin{aligned} \frac{\Delta \rho}{\rho} &= \frac{\delta \rho}{\rho} - \frac{1}{\rho} (\delta \mathbf{r} \cdot \nabla) \rho = \frac{1}{\Gamma_1} \frac{\delta P}{P} - \frac{1}{\rho} (\delta \mathbf{r} \cdot \nabla) \rho \\ &= \frac{1}{\Gamma_1} \frac{\Delta P}{P} + (\delta \mathbf{r} \cdot \nabla) \left[ \frac{1}{\Gamma_1} \ln P - \ln \rho \right]. \end{aligned}$$

Since by assumption  $\Delta P = 0$ , the radial component of equation (9.19) becomes

$$\delta \ddot{r} = g \left[ \frac{d \ln \rho}{dr} - \frac{1}{\Gamma_1} \frac{d \ln P}{dr} \right] \delta r \equiv g \mathcal{A} \delta r. \quad (9.20)$$

The quantity  $\mathcal{A}$  is called the *Schwarzschild discriminant*: if  $\mathcal{A} < 0$ , then the motion is oscillatory with frequency  $N = (-g\mathcal{A})^{1/2}$ ;  $N$  is called the *Brunt-Väisälä frequency*. The condition  $\mathcal{A} > 0$  implies that the fluid is convectively unstable; and indeed, one can show that  $\mathcal{A} > 0$  is equivalent to  $dS/dr < 0$ .

The utility of using  $\mathcal{A}$  rather than  $dS/dr$  is that  $\rho$  and  $P$  appear in the equations of stellar structure.

### 9.3 Efficiency of Heat Transport

A superadiabatic temperature gradient,  $\nabla > \nabla_{ad}$ , induces convective motions. A rising blob will be hotter than its surroundings and heat will therefore be conducted from the blob to its surroundings as it rises. The efficiency by which the heat is transported determines by how much convection is able to drive the temperature gradient towards an adiabat. Clearly the gradient must be super-adiabatic to drive the convection in the first place. We shall see, however, that in stars the difference between the gradient and the adiabat are typically exceedingly small. In other words, convection is extraordinarily efficient at transporting heat.

To understand this, let's go back to equation (9.19). Write  $\Delta \rho$  as stemming from differences in temperature between rising and falling

blobs (recall that  $\Delta P = 0$ ). With this substitution, we have

$$(\partial_t \mathbf{u} + \mathbf{u} \cdot \nabla \mathbf{u}) = \frac{\Delta \rho}{\rho} \mathbf{g} = \left( \frac{\partial \ln \rho}{\partial \ln T} \right)_P \frac{\Delta T}{T} \mathbf{g}. \quad (9.21)$$

Our goal is to estimate the velocity of convective motions  $\mathbf{u}$ , the departure of the temperature gradient from an adiabat  $\Delta T$ , and the fraction of the total heat flux carried by convective motions from these equations.

First, the velocity. The left-hand side of equation (9.21) has a characteristic scale  $\sim U^2/L$ , whereas the right-hand side has a scale  $g\Delta T/T$ . (Recall that in an ideal gas,  $(\partial \ln \rho / \partial \ln T)_P = -1$ .) If we take  $L \sim c_s^2/g$ , a pressure scale height, than we get an estimate of the convective velocity,

$$\frac{U}{c_s} \sim \left( \frac{\Delta T}{T} \right)^{1/2}. \quad (9.22)$$

What is the heat flux carried by convection? Hot fluid rises and carries an excess of heat, per gram, of  $c_P \Delta T$ , giving a heat flux  $\approx \rho u c_P \Delta T$ . Thus to carry a given flux  $F$ , we have

$$c_s \rho c_P T \left( \frac{\Delta T}{T} \right)^{3/2} \sim F. \quad (9.23)$$

Note that in order of magnitude,  $c_P T \sim c_s^2$ , so

$$\frac{U}{c_s} \sim \left( \frac{\Delta T}{T} \right)^{1/2} \sim \left( \frac{F}{\rho c_s^3} \right)^{1/3}.$$

For conditions in the solar interior,  $F \ll \rho c_s^3$ , and therefore the convective velocities are very subsonic. Indeed,

$$\begin{aligned} \frac{F}{\rho c_s^3} &\sim \frac{L_\odot}{4\pi R_\odot^2} \frac{4\pi R_\odot^3}{3M_\odot} \left( \frac{R_\odot}{GM_\odot} \right)^{3/2} \\ &\sim \frac{L_\odot}{GM_\odot^2/R_\odot} \left( \frac{R_\odot^3}{GM_\odot} \right)^{1/2} \\ &\sim \frac{t_{\text{dyn}}}{t_{\text{KH}}} \ll 1. \end{aligned}$$

That is, the ratio of the solar flux to what could be carried for near-sonic convective motions is of the order of the dynamical timescale to the Kelvin-Helmholtz timescale. We therefore expect that in a convective region, slow circulation will produce a temperature gradient that is very nearly adiabatic. This argument breaks down near the surface, where the cooling time of a fluid layer (the “local” Kelvin-Helmholtz timescale) can be small.

## 9.4 Turbulence

From the discussion of the previous section, it might seem possible, given the boundary conditions, of solving for the flow planform, that is, the velocity profile  $u(x, t)$ . This is decidedly not the case, however: the flow is turbulent, with intermittent velocity fluctuations seen over a large dynamical range of spatial and temporal scales. Modeling of such flows is a vexing problem in fluid dynamics.

To explore this topic a bit further, we need to introduce the concept of dynamical similarity. Suppose you want to optimize a wing shape for an aircraft, and you wish to test its performance in a wind tunnel. Why should you expect that the behavior of a model wing will have any relation to the full-scale one?

To see how this works, start with the Navier-Stokes equation (for simplicity, we'll keep it in one dimension):

$$(\partial_t + u \cdot \partial_x)u = -\frac{1}{\rho}\partial_x P + \nu\partial_x^2 u. \quad (9.24)$$

Here  $\nu$  is the coefficient of kinematic viscosity, with dimensions  $[\nu] \sim [\text{length}]^2 \cdot [\text{time}]^{-1}$ . Let's recast equation (9.24) into dimensionless form by scaling our variable: let  $L$  and  $U$  represent the characteristic length and velocity scales, and define the dimensionless variables  $\tilde{x} = x/L$  and  $\tilde{u} = u/U$ . This choice then implicitly defines the time variable,  $\tilde{t} = t \cdot U/L$ . Upon changing to the variables  $\tilde{x}$ ,  $\tilde{u}$ , and  $\tilde{t}$ , and writing the equation of state as  $P = c_s^2\rho$  (appropriate for adiabatic flow—we are ignoring heat conduction), we obtain the equation

$$(\partial_{\tilde{t}} + \tilde{u} \cdot \partial_{\tilde{x}})\tilde{u} = -\left\{\frac{c_s^2}{U^2}\right\}\partial_{\tilde{x}} \ln \tilde{\rho} + \left\{\frac{\nu}{UL}\right\}\partial_{\tilde{x}}^2 \tilde{u}. \quad (9.25)$$

Each term in this equation is dimensionless. The physical characteristics of the fluid and the scales involved are described by just two dimensionless parameters:

$$\text{Ma} \equiv \frac{U}{c_s} \quad \text{Mach number (measure of compressibility)}$$

$$\text{Re} \equiv \frac{UL}{\nu} \quad \text{Reynolds number (measure of viscous forces)}$$

So, if we build a model wing at a certain scale and place it in a wind tunnel with a certain velocity, then by adjusting the density and temperature (and hence the sound speed and viscosity) to the desired Ma and Re, the flow pattern in our model will faithfully replicate the flow in the actual system.

For stellar convection,  $\text{Ma} \ll 1$ . What about Re? In typical astrophysical plasmas, the large lengthscales make Re ludicrously large. Terrestrial experiments and simulations cannot approach this

regime. Experimentally, when  $\text{Re} \gtrsim 10^3$ , then flow becomes *turbulent*: the velocity has strong intermittent fluctuations across a wide range of lengthscales and timescales. How to characterize the flow in such a case? It is useful to describe the flow in terms of correlated velocities—an “eddy”—which have some lengthscale.

Suppose we pass water through a pipe that has an embedded mesh screen (Fig. 9.5). For sufficiently large  $\text{Re} = UL/\nu$ , where  $L$  is the mesh spacing, the downstream flow becomes turbulent. The turbulent eddies are damped. Now an eddy of size  $\lambda$  has an effective Reynolds number  $\text{Re}(\lambda) = U(\lambda)\lambda/\nu$ ; if this is very large, then molecular viscosity cannot be the reason for damping fluid motions on that scale. Instead what happens is that an eddy with lengthscale  $\lambda$  and velocity scale  $U(\lambda)$  drives eddies on a smaller scale  $\lambda' < \lambda$ . These in turn drive still smaller eddies, which in turn drive still smaller eddies, and so on, until eventually very tiny eddies are excited, with size  $\lambda_\nu \sim \nu/U(\lambda_\nu)$ ; and these eddies are damped by viscosity!

Kolmogorov argued that in steady-state, intermediate-sized eddies (i.e., those with lengthscales  $\nu/U(\lambda) \ll \lambda \ll L$ ) are neither losing or gaining energy and hence were transferring energy to smaller scales at the same rate as they were being driven; further, this rate at which energy is being transferred to smaller scales is just the net rate of dissipation in the fluid (which is done by the smallest eddies). The huge dynamic range in lengthscales implies that the velocity of the eddy should not depend on either  $L$  or  $\nu$ , and hence  $U(\lambda)$  can only be a function of  $\lambda$  (length) and the rate of energy dissipation per unit mass  $\varepsilon$  (energy/mass/time  $\sim \text{length}^2/\text{time}^3$ ). There is only one way to combine these quantities to form something with a dimension of length/time, and so

$$U(\lambda) \sim \varepsilon^{1/3} \lambda^{1/3}. \quad (9.26)$$

This is seen experimentally: in flows with a large dynamic range of scales, the velocity spectrum follows a power-law with this slope, over an intermediate range of scales, the *inertial range*. A good example is the flow in a tidal channel<sup>1</sup>.

### MESA Convection in a pre-main-sequence star

For this exercise, we'll use the setup from 2.6, namely a contracting, pre-main-sequence star of mass  $1 M_\odot$ . Download the folder `convection/1M-convection` and place it into your projects folder.

For this exercise, we don't need any custom output, so we are just using the standard `'run_star_extras.f'` file. We still need to compile the code, however, so do `'./mk'`.

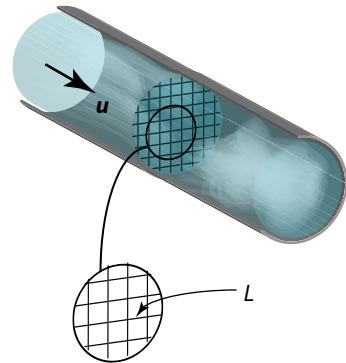


Figure 9.5: A simple mechanism for generating turbulence. A flow of water in a pipe (upstream velocity  $U$ ) flows through a mesh (spacing  $L$ ). If  $\text{Re} = UL/\nu$  is sufficiently large, the downstream flow becomes turbulent.

<sup>1</sup> H. L. Grant, R. W. Stewart, and A. Moilliet. Turbulence spectra from a tidal channel. *Journal of Fluid Mechanics*, 12:241–268, 1962

The first thing we want to plot is the entropy in the star. If we look at the file ‘\$MESA\_DIR/star/defaults/pgstar.defaults’, we notice that `Profile_Panels1` is close to what we want: its first panel plots both  $\log(T/K)$  and  $S/(N_A k_B)$ . We don’t need the second panel, though, so we’ll make the following changes to ‘`basic_plot.inlist`’. We’ll swap out the ‘`TRho`’ plot for ‘`Profile_Panels1`’, and then we’ll set the number of panels to be just 1. The following lines of `basic_plot.inlist` accomplish this:

```
34     Grid1_plot_name(3) = 'Profile_Panels1'
35     ! following are specific settings for that plot
36     Profile_Panels1_title = 'thermal profile'
37     Profile_Panels1_num_panels = 1
```

The defaults for the rest of the settings for ‘`Profile_Panels1`’ are fine, so we don’t need to set them explicitly. Now when you do ‘`./rn`’ the window will have a plot of both entropy and temperature as functions of enclosed mass in the lower left panel.

#### EXERCISE 9.2—

1. Compare the entropy profile when the star is fully convective, and when a radiative region develops as it approaches the main sequence. Does the profile match your expectation given in the warm-up exercise?
2. Look at  $\log[T(m)]$ . In Ch. 1 of the notes, we discuss estimates of the interior temperature, and we often use the central value as a representative value. Is this reasonable? Make this quantitative: at the point when the star joins the main sequence, find the mass at which  $\log[T(m)] = \log[T_c] - 0.5$  and  $\log[T(m)] = \log[T_c] - 1.0$ .

#### *Convective efficiency*

We argued in § 9.3 of the notes that convection was very efficient in the solar interior. By efficient, we mean that a very slow convective velocity  $v_{\text{conv}}$  is sufficient to carry the flux. As a result, the difference between the temperature gradient and the adiabatic one is also expected to be very small:  $\nabla - \nabla_{\text{ad}} \ll 1$ , where  $\nabla \equiv d \ln T / d \ln P$  and  $\nabla_{\text{ad}} = (\partial \ln T / \partial \ln P)_s$ . We also worked through a linear stability analysis in § 9.2 and found (eq. 9.20) the Brunt-Väisälä frequency  $N^2$ . This is the oscillation frequency for a fluid element that is adiabatically displaced in the radial direction;  $N^2 < 0$  in a region with  $\nabla > \nabla_{\text{ad}}$  and therefore unstable to convection.

For this project, you are going to examine  $\nabla_{\text{ad}}$ ,  $\nabla_{\text{ad}} - \nabla$ ,  $N^2$ , and  $v_{\text{conv}}/c_s$ . In the first assignment, we had to customize ‘`run_star_extras.f`’ to output specific quantities, but in this case the information we want

is already computed by MESA. A complete listing of the quantities available for each zone is in ‘\$MESA\_DIR/star/defaults/profile\_columns.list’. Browsing this list, we see that ‘grada’ ( $\equiv \nabla_{\text{ad}}$ ) and ‘conv\\_vel\\_div\\_csound’ ( $\equiv v_{\text{conv}}/c_s$ ) are both written by default to ‘profiledd.data’; the other two variables are computed by MESA, but are not, by default, written to file (in ‘profile\_columns.list’ they are commented out, i.e., preceded by a bang ‘!’). To enable the printing of these two variables, we set, in the ‘star\_job’ section of the ‘1M\_convection.inlist’, the flag

```
25 profile_columns_file = 'convection_variables'
```

The file ‘convection\_variables’ first includes the defaults

```
18 include ''
```

and then sets the remaining two variables  $\nabla_{\text{ad}} - \nabla$  and  $N^2$ ,

```
20 ! now we'll include the difference between the adiabatic gradient
21 ! and the temperature gradient in the star
22 grada_sub_gradT
23
24 ! and then the Brunt frequency, scaled to a dimensionless number
25 brunt_N2_dimensionless ! N2 in units of 3GM/R^3
```

With these four quantities now being written to the profile data, we can plot them. I’ve copied a template for this, which uses the ‘Profile\_Panels6’ variables, into the file ‘plot\_convection\_vars.inlist’. You will need to customize this to produce the desired plot. There are many ways to plot the four variables. Figure 9.6 shows one; see if you can duplicate it! Note that in making this plot you might have to set the minima and maxima for the plots; for some of these quantities the values near the surface are at a very different scale than throughout the bulk of the star. We will look in more detail at what happens near the photosphere in the next exercise.

**Tip:** Unlike the parameters in ‘&controls’ and ‘&star\_job’, the settings in ‘&pgstar’ are reread at each step. This means you can change parameters, such as the minimum and maximum values for axes, while the program is running! To pause the program execution during an interactive run, type ‘ctrl-z’; you should then get a message such as

```
1 ^Z
2 [1]+ Stopped ./rn
```

To restart, type ‘%1’ (where the ‘1’ is whatever number precedes the word ‘Stopped’).

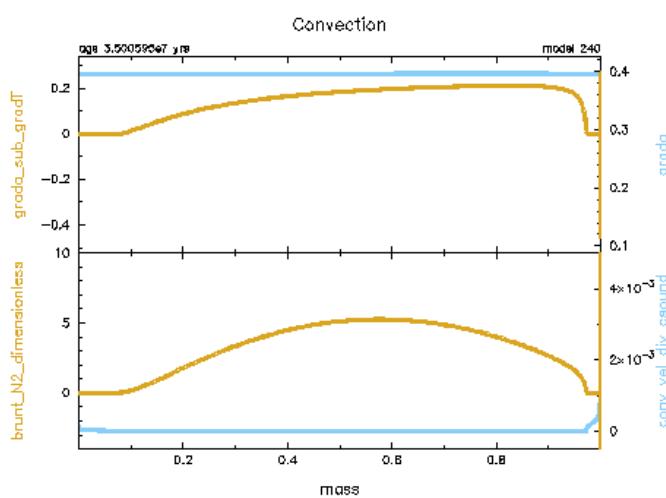


Figure 9.6: Snapshot of convective variables throughout the star.

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#### EXERCISE 9.3—

1. How does the value of  $\nabla_{ab} \equiv (\partial \ln T / \partial \ln P)_s$  compare with the value expected for an ideal gas?
  2. The Brunt frequency  $N^2$  that is plotted is actually in units of  $3GM/R^3$ ; explain from the derivation in the notes why the factor  $GM/R^3$  is sensible. What is a characteristic value of  $N$  (when the star has approached the main sequence and most of the star is not convective)?
- 

#### *Approaching the photosphere*

As you will notice from these plots, there is a lot happening for  $m(r) > 0.99 M_\odot$ , but all of the action is compressed against the right-hand side of the plot. Wouldn't it be nice to rescale the x-axis to zoom in on this region? Fortunately, there is an easy way to do this. MESA defines a variable  $q = m(r)/M$ , and one of the variables written out by default is ‘logxg’, defined as  $\log(1 - q)$ . Since  $q$  ranges from 0 at the stellar center to 1 at the surface, ‘logxq’ is 0 at the center and goes to a very large, negative number near the surface. This means that the surface will be at the left-hand edge of the plot and the center at the right, so the orientation is opposite to plots using ‘mass’ or ‘radius’ on the x-axis. You can switch the orientation of the x-axis by setting ‘Profile\_Panels6\_xaxis\_reversed = .true.’.

## EXERCISE 9.4—

1. Re-run the evolution, this time setting the x-axis to ‘`logxq`’ for the plot of  $N^2$ ,  $v_{\text{conv}}/c_s$ ,  $\nabla_{\text{ad}}$ , and  $\nabla_{\text{ad}} - \nabla$ . Setting

```
1 Profile_Panels6_xmin = -12
2 Profile_Panels6_xaxis_reversed = .true.
```

is a good choice here. You will notice that  $\nabla_{\text{ad}}$  behaves in an interesting fashion in the outer regions of the star. Do  $\rho$  and  $T$  take on any characteristic values at the location of the feature? Hypothesize about what may be happening there.

---

*What to turn in*

Make plots for models when the center of the star first becomes radiatively stable, and when the star reaches the main sequence. To have the plots written to a file, set ‘`Grid1_file_flag = .true.`’ and ‘`Profile_Panels6_file_flag = .true.`’. Make sure that the flags ‘`Grid1_file_dir`’ and ‘`Profile_Panels6_file_dir`’ are both set to ‘`frames`’. There is no need to generate a plot for each timestep; you can set the interval between plots via ‘`Grid1_file_interval`’ and ‘`Profile_Panels6_file_interval`’ (I set them to be `10`).

## 10

# Nuclear Physics

The lives (and deaths) of stars are largely concerned with the history of the nuclear burning they undergo. Following the gravitational collapse of a protostar and the subsequent contraction to the main sequence (see Chapter 11), nuclear reactions involving the transmutation of H into He will provide sufficient energy input to balance the radiative luminosity from the stellar surface. The star will reach radiative and hydrostatic equilibrium. Or at least, nearly so. This phase, during which a star processes H into He, will be the longest of any phase of stellar evolution, lasting for the Sun some 10 Billion years. This time scale is quite accelerated for more massive stars as they consume their H fuel at incredibly rapid rates, as we shall see in subsequent chapters.

In order understand how nuclear burning plays a key role in stellar evolution, we must develop some basic knowledge of nuclear physics and nuclear reactions. This then will provide the final key piece of the puzzle that is the equations of stellar structure (Chapter 6), the elusive local source term due to nuclear reactions  $q$ . We shall first describe the basic terminology and definition used to describe nuclear reactions. Then we will discuss the nuclear landscape and binding energies. We will then proceed to computing non-resonant reaction cross sections and rates. After that, we will discuss all important resonant reactions. Following a brief presentation of computing inverse reaction rates, we will proceed to how reaction rates need be modified to account for important plasma effects. Finally, we will discuss how all of these rates are folded into systems of differential equations describing the evolution of composition and attendant energy release.

Many excellent and “classic” references are available on these topics that go into far greater detail than we will be able. Perhaps the most “classic” (yet still useful...) reference is <sup>1</sup>. A good exposition of nuclear reactions and networks, and their role in stellar evolution, can be found in <sup>2</sup>. More modern, and slightly experimental, takes can

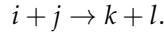
<sup>1</sup> Donald D. Clayton. *Principles of Stellar Evolution and Nucleosynthesis*. University of Chicago Press, 1983

<sup>2</sup> D. Arnett. *Supernovae and Nucleosynthesis: An Investigation of the History of Matter from the Big Bang to the Present*. 1996

be found in <sup>3</sup>. And the chapter on nuclear energy generation in <sup>4</sup> is also a useful summary of the key aspects.

### 10.1 Nuclear Nomenclature

Consider a reaction that transmutes two bodies into some other two bodies,



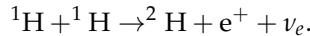
For nuclear reactions we have two key constraints: charge conservation  $Z_i + Z_j = Z_l + Z_k$ , and baryon number conservation  $A_i + A_j = A_l + A_k$ . Reactions must also conserve lepton number, so we need to keep track of electrons, positrons, neutrinos, etc. We define the *reaction rate*  $r_{ijkl}$  as the number of reactions per second per gram. The change in energy for the system due to a single such reaction is given by the *reaction Q-value*,

$$Q_{ijkl} = (m_i + m_j - m_k - m_l)c^2, \quad (10.1)$$

where the term in parentheses is the relativistic change in mass, or *mass defect*, of the reaction. This corresponds to a change in binding energy amongst the nuclei involved. Finally, the total *energy generation rate* for a given reaction is

$$\epsilon_{ijkl} = r_{ijkl}Q_{ijkl}. \quad (10.2)$$

In nuclear astrophysics, we always consider any complex reaction to be the result of chains of two-body interactions. Three- and more-body interactions are...complicated. It is often the case that a give reaction will generate a short-lived, highly-excited compound nucleus that decays so quickly, we typically elide over it's very existence. For instance, the first step in the proton-proton chain, the primary nuclear fusion going on in the core of Sun, the first reaction is



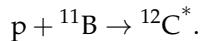
But in fact, the two protons must first make a  ${}^2\text{He}$  nucleus. But this is an incredibly unstable nucleus and so must decay almostly instantly, either back into protons or to deuterium,  ${}^2\text{H}$ , by the emission of the a positron (and electron neutrino in order to conserve lepton number).

It is common to write nuclear reactions using a shorthand notation,  $X(\alpha, \beta)Y$ , where the “target” nucleus is  $X$ , the “projectile” impacting the target is  $\alpha$ , and the final nuclear products are  $Y$  and  $\beta$ . Thus, we would write the proton-proton interaction above as  ${}^1\text{H}({}^1\text{H}, \text{e}^+ + \nu_e){}^2\text{He}$ .

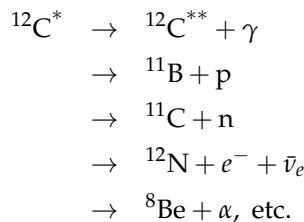
<sup>3</sup> C. Iliadis. *Nuclear Physics of Stars*. Wiley-VCH, 2007

<sup>4</sup> Carl J. Hansen, Steven D. Kawaler, and Virginia Trimble. *Stellar Interiors*. Springer-Verlag, 2d edition, 2004

Let's consider a more complex reaction, the creation of an excited  $^{12}\text{C}$ nucleus. The first reaction is proton capture onto  $^{11}\text{B}$



This first reaction leading to the excited nucleus is typically called the "entrance channel" for the reaction. This, then, is followed by the decay of the excited carbon nucleus via one of a number of reactions, or "exit channels:"



The first reaction is simply the generation of some different excited state via the emission of a gamma-ray photon. Not all exit channels are equally likely. The probability that a given decay will occur can be computed by comparing the mean lifetimes for the given states,  $\tau_i$ .

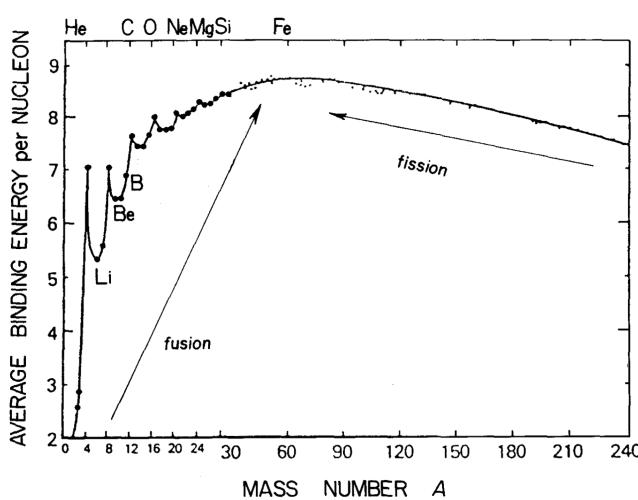
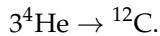


Figure 10.1: Binding energy per nucleon.

Computing the  $Q$  value of a reaction involves comparison the change in binding energy or the the reactants or, equivalently, the mass defect of the reaction. Figure 10.1 shows the classic, experimentally derived binding energy per nucleon as a function of charge number  $Z$ . For example, consider the creation of  $^{12}\text{C}$ via the *triple- $\alpha$*  reaction,



The binding energy per nucleon of the  $^4\text{He}$ nucleus is 7.074 MeV and that for  $^{12}\text{C}$ is 7.680 MeV. We can compute the reaction  $Q$  value then

as

$$\begin{aligned} Q_{3\alpha} &= 3 \times [4 \times (-7.1 \text{ MeV})] - 12 \times (-7.7 \text{ MeV}) \\ &= -85.2 \text{ MeV} + 92.4 \text{ MeV} = 7.2 \text{ MeV}. \end{aligned}$$

## 10.2 The Nuclear Landscape

From nucleon-nucleon scattering, we find that the nuclear force operates over a range of  $\lesssim 2 \times 10^{-13} \text{ cm} \equiv 2 \text{ fm}$ , where fm denotes the unit of length known as a “fermi.” At low energies the nuclear force can be described as the exchange of spin-zero pions; recall from quantum field theory that the exchange of a spin-zero particle produces an attractive potential of the form  $e^{-r/\lambda_\pi}/r$ , where  $\lambda_\pi = \hbar/(m_\pi c)$  is the Compton wavelength of the force carrier. The pion rest mass is  $\approx 140 \text{ MeV}/c^2$ , so  $\lambda_\pi \approx 1.4 \text{ fm}$ . The nuclear force is indeed attractive over this range, but it becomes repulsive at distances  $< 1 \text{ fm}$  where higher order terms in the interaction become important.

This interaction—short range and attractive, but with a repulsive core—resembles the interaction between molecules in a fluid. Indeed, for heavy nuclei, the nucleons form a “nuclear fluid” with a characteristic density of  $0.16 \text{ fm}^{-3}$ . The **binding energy** is defined as

$$B(N, Z) = [Zm_p + Nm_n - M(N, Z)] c^2, \quad (10.3)$$

where  $M$  is the total mass of a nucleus with  $N$  neutrons and  $Z$  protons. The total number of nucleons is  $A = N + Z$ , the proton rest mass is  $m_p = 938.272 \text{ MeV}/c^2$ , and the neutron rest mass is  $m_n = 939.565 \text{ MeV}/c^2$ . From the form of this definition  $B > 0$  for bound nuclei and is  $\approx 8 \text{ MeV}$  for most nuclei. This binding of the nucleons into a “nuclear fluid” implies that one can make a crude model of the nucleus as a liquid drop, and use this model to fit the binding energy. Such a model, due to Weizäcker, is

$$-B(N, Z) = a_V A + a_S A^{2/3} + a_A \frac{(N - Z)^2}{A} + a_C \frac{Z^2}{A^{1/3}} + a_p \delta A^{-1/2}, \quad (10.4)$$

and has five terms: a bulk energy term  $a_V A$  that scales with the total number of nucleons; a surface term  $a_S A^{2/3}$  that corrects for the weaker binding near the surface; an asymmetry term  $a_A (N - Z)^2 / A$  that accounts for the energy cost to have an imbalance in the number of neutrons and protons; and a Coulomb term  $a_C Z^2 / A^{1/3}$  that accounts for the repulsion between other protons. The final term

accounts for the pairing between like nucleons, with

$$\delta = \begin{cases} +1, & N, Z \text{ even;} \\ -1, & N, Z \text{ odd;} \\ 0, & N \text{ even, } Z \text{ odd, or } N \text{ odd, } Z \text{ even} \end{cases} .$$

A sample fit for the coefficients is listed in Table 10.1.

coefficient	$a_V$	$a_S$	$a_A$	$a_C$	$a_p$
5-parameter fit (MeV)	-15.67	17.04	23.09	0.71	-14.55
4-parameter fit (MeV)	-15.5	16.6	22.7	0.71	—

Table 10.1: Coefficients for the Weizäcker mass formula.

**EXERCISE 10.1—** In the r-process, a heavy seed nucleus captures a large number of neutrons and then decays back to stability. Suppose we start with  $^{56}\text{Fe}$  in a bath of free neutrons, so that the iron nucleus captures 152 neutrons (with  $\beta$ -decays occurring as necessary to keep the nucleus bound) until it reaches the stable nucleus  $^{208}\text{Pb}$ . Is this process exothermic or endothermic? Explain your answer.

Equation (10.4) gives a good, if somewhat crude, description of the nuclear landscape. As a first example, let's look at how the binding energy per nucleon,  $-B/A$  trends with  $A$ . For simplicity, we'll ignore the pairing term. Dividing equation (10.4) by  $A$ , and denoting the neutron asymmetry by  $\eta \equiv (N - Z)/(N + Z)$ , we obtain

$$-\frac{B}{A} = a_V + a_S A^{-1/3} + a_A \eta^2 + \frac{a_C}{4} (1 - \eta)^2 A^{2/3}. \quad (10.5)$$

Minimizing this expression for small  $A$ , we see that the most bound nuclei (smallest  $-B/A$ ) have  $\eta \rightarrow 0$ , i.e., equal numbers of neutrons and protons. As  $A$  increases, however, the Coulomb term becomes important. Expanding the last two terms and combining gives the sum of the asymmetry and Coulomb terms,

$$a_A \eta^2 + \frac{a_C}{4} (1 - \eta)^2 A^{2/3} = \frac{a_C}{4} A^{2/3} - \frac{a_C}{2} A^{2/3} \eta + \left[ \frac{a_C}{4} A^{2/3} + a_A \right] \eta^2.$$

For large  $A$ , this expression is minimized (although it cannot be made to vanish) for  $\eta > 0$ : that is,  $N > Z$  and the most bound massive nuclei are neutron-rich. The nuclei for which  $-B/A$  is a minimum for a fixed  $A$  define the *valley of stability*. How does the binding energy change with  $A$  along this valley? At small  $A$ , the asymmetry term dominates and  $\eta^2 \ll 1$ . As  $A$  increases, the surface term  $a_S A^{-1/3}$  becomes smaller and  $B/A$  increases. At large  $A$ , however, the sum of the Coulomb and asymmetry terms decreases  $B/A$ . As a result, there is a peak in  $B/A$ , which is around  $A = 56$ .

Next, let's find the boundaries of our nuclear landscape: the most neutron-rich and proton-rich nuclei that are still bound. Define the

**neutron separation energy**  $S_n$  as the energy needed to remove a neutron from a nucleus,

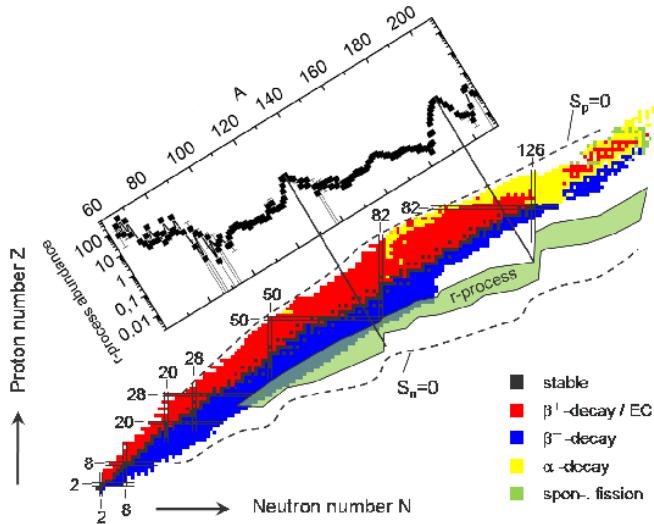
$$\begin{aligned} S_n(N, Z) &\equiv c^2 \{ [M(N-1, Z) + m_n] - M(N, Z) \} \\ &= B(N, Z) - B(N-1, Z). \end{aligned} \quad (10.6)$$

Likewise, define the **proton separation energy** as

$$\begin{aligned} S_p(N, Z) &\equiv c^2 \{ [M(N, Z-1) + m_p] - M(N, Z) \} \\ &= B(N, Z) - B(N, Z-1). \end{aligned} \quad (10.7)$$

If we take a nucleus  $(N, Z)$  in the valley of stability and add protons keeping  $N$  fixed, we will eventually reach a nucleus for which  $S_p = 0$ , that is, it costs no energy to add or to remove a proton. This defines the **proton-drip line**: nuclei more proton-rich are unstable to proton emission. Likewise, on the neutron-rich side there is the **neutron-drip line**, for which  $S_n = 0$ . Note that because of the pairing term there is an odd-even staggering in the  $S_n$  and  $S_p$ ; it is therefore useful, sometimes, to define the two-neutron and two-proton separations energies  $S_{2n}$  and  $S_{2p}$ . Figure 10.2 shows an annotated nuclear landscape in neutron number and proton number space.

Figure 10.2: The Nuclear Landscape



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EXERCISE 10.2—

1. (a) For a fixed  $A$ , find  $Z_*(A)$  such that the binding energy per nucleon,  $f = B(N = A - Z_*, Z_*)/A$  is maximized.
  - (b) Plot  $Z_*$  vs  $N$ .
  - (c) Using this  $Z_*$ , plot  $Y_e = Z_*/A$  for  $20 < A < 200$  and explain qualitatively any trends.
  - (d) Now substitute the value of  $Z_*$  into the expression for  $B(N, Z)$  and plot  $B(N, Z_*)/A$  as a function of  $A = N + Z_*$ . Explain qualitatively any trends.
  2. (a) For each  $10 \leq Z \leq 82$ , find the maximum value of  $N$  such that  $S_{2n}(N, Z) > 0$ . Plot the values  $(N, Z)$  you find.
  - (b) For each  $10 \leq N \leq 120$ , find the maximum value of  $Z$  such that  $S_{2p}(N, Z) > 0$ . Plot the values  $(N, Z)$  you find.
  3. Compare the plots of problems 1b, 2a, and 2b to a chart of the nuclides.
- 

**10.3 Non-resonant nuclear reactions**

The situation of interest is the reaction between two nuclei,  $(A_1, Z_1)$  and  $(A_2, Z_2)$ . The nuclear radius is  $r_N \approx A^{1/3}$  fm, and the Coulomb energy at this distance is

$$\frac{Z_1 Z_2 e^2}{r_N} = \frac{Z_1 Z_2 \alpha \hbar c}{r_N} \approx 1.4 Z_1 Z_2 A^{-1/3} \text{ MeV} \gg kT. \quad (10.8)$$

For nuclear reactions, typical energy scales are  $\sim$  MeV and typical length scales are  $\sim$  fm. In these units,  $\hbar c = 197$  MeV fm. In the first equality in eq. (10.8), we also introduce the fine-structure constant  $\alpha = e^2/(\hbar c) = 1/137$ . In “nuclear units,”  $e^2 = \alpha \hbar c = (197 \text{ MeV fm})/137 = 1.44 \text{ MeV fm}$ . Remember these numbers!

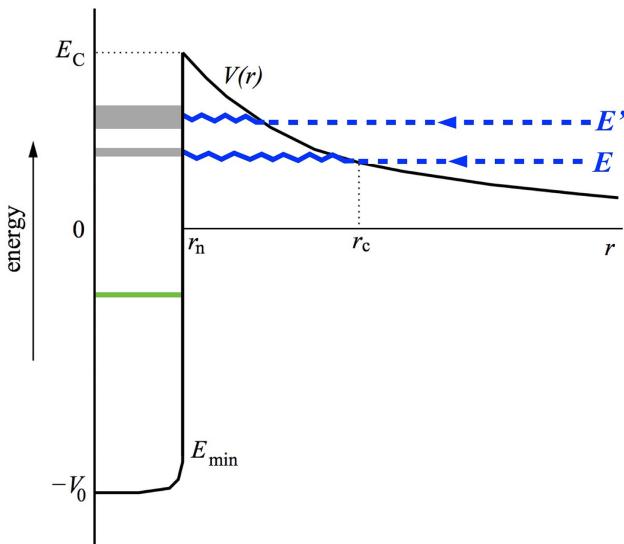
If we scatter two nuclei together, the closest approach (cf. §8.1) is  $\sim e^2/(k_B T) \sim 1440$  fm at typical stellar energies  $k_B T \sim 1$  keV.

Clearly the cross-section for a reaction between our pair of particles is controlled by the probability of tunneling through the Coulomb potential. The situation is illustrated in Figure 10.3.

For a two-body system, it is convenient to transform into a center-of-mass frame. Our problem then reduces to a one-body problem with reduced mass  $m = Am_u$ , with  $A = A_1 A_2 / (A_1 + A_2)$  and incident energy  $E = mv^2/2$ , where  $v$  is the relative velocity of the two particles. For now, we'll neglect angular momentum ( $\ell = 0$ ) so our scattering is s-wave. At low energies, we can form a “geometrical” cross-section from the particle wavenumber  $k = p/\hbar$ , with

$$\pi k^{-2} = \pi \frac{\hbar^2}{(2mE)} = 660 \text{ b} \frac{1}{A} \left( \frac{\text{keV}}{E} \right) \quad (10.9)$$

Figure 10.3: The Coulomb barrier.



Here the cross-section is in units of *barns*, with  $1\text{ b} = 10^{-24}\text{ cm}^2$ . This is the first part of our nuclear cross-section  $\sigma(E)$ .

The second portion of the nuclear cross-section is the probability of tunneling through the Coulomb barrier. First, let's get the classical turning point  $r_E$  from

$$\begin{aligned} \frac{Z_1 Z_2 e^2}{r_E} &= E, \\ r_E &= 1440 \text{ fm } Z_1 Z_2 \left( \frac{\text{keV}}{E} \right). \end{aligned} \quad (10.10)$$

Now the wavelength is  $k^{-1} = \hbar(2Am_u E)^{-1/2} = \hbar c(2Am_u c^2 E)^{-1/2}$  and since  $m_u c^2 = 932 \text{ MeV}$  the wavelength  $k^{-1} = 145 \text{ fm (keV/E)}^{1/2}$ . The important point is that since  $k^{-1} \ll r_E$ , we can solve the Schrödinger equation using the WKB approximation.

The WKB approximation is standard, so let me just remind you that the probability of tunneling through the barrier depends on the *action*,

$$\mathcal{P} \propto \exp \left\{ \frac{2}{\hbar} \int_{r_E}^{r_N} \left[ 2m \left( \frac{Z_1 Z_2 e^2}{r} - E \right) \right]^{1/2} dr \right\}. \quad (10.11)$$

To do this integral, note that  $r_E \gg r_N$ , so we can make the approximation  $r_N \rightarrow 0$  in the integral's lower limit; with the substitution

$$\sin \phi = \left[ 2m \left( \frac{Z_1 Z_2 e^2}{r} - E \right) \right]^{1/2} \left( \frac{r}{2m Z_1 Z_2 e^2} \right)^{1/2}$$

we tame the integral and obtain

$$\mathcal{P} \propto \exp \left\{ -\frac{8mZ_1Z_2e^2}{\hbar(2mE)^{1/2}} \int_0^{\pi/2} \sin^2 \phi \, d\phi \right\} = \exp \left[ -\left( \frac{E_G}{E} \right)^{1/2} \right], \quad (10.12)$$

where

$$E_G \equiv 2\pi^2 A m_u c^2 \alpha^2 (Z_1 Z_2)^2 = 979 \text{ keV } A (Z_1 Z_2)^2 \quad (10.13)$$

is the *Gamow energy*. Note the strong dependence on  $Z_1 Z_2$ :  $E_G$  determines which reactions can occur at a given temperature. If you stare at the factor multiplying the integral in equation (10.12), you will see that  $\mathcal{P} \propto \exp(-r_E/\lambda)$ , the exponential of ratio of the width of the forbidden region to the wavelength of the incident particle. This makes intuitive sense.

Now we have the second part of our cross-section, the probability of getting through the Coulomb barrier. This third part depends on the nuclear interactions. For non-resonant reactions, this third part does not depend strongly on energy, so it is common to define the *astrophysical S-factor* by writing the cross section as the product (geometrical)  $\times$  (tunneling)  $\times$  (nuclear),

$$\sigma(E) = \frac{1}{E} \exp \left[ -\left( \frac{E_G}{E} \right)^{1/2} \right] S(E). \quad (10.14)$$

It is easier to extrapolate the slowly varying  $S(E)$  from lab energies of  $> 100 \text{ keV}$  down to center-of-mass energies of  $\sim \text{keV}$  than it would be to fit the rapidly varying cross-section.

Now each nucleus has a Maxwellian velocity distribution,

$$n_1(v_1) d^3v = n_1 \left( \frac{m_1}{2\pi kT} \right)^{3/2} \exp \left( -\frac{mv^2}{2kT} \right) d^3v, \quad (10.15)$$

and similarly for particle 2. Let's call a particular nucleus 1 (having velocity  $v_1$ ) the target. By definition the cross section is

$$\frac{\text{number of reactions/target/time}}{\text{number of incident particles/area/time}},$$

so to get the number of reactions per target per time we need to multiply  $\sigma(E)$  by the number of incident particles per unit area per unit time. The incident flux is just  $n_2(v_2)|v| d^3v_2$  where  $v = v_2 - v_1$ . Hence the reaction rate per unit volume per unit time between a pair of particles having velocities in volumes  $d^3v_1$  and  $d^3v_2$  about  $v_1$  and  $v_2$  is just

$$\frac{1}{1 + \delta_{12}} n_1(v_1) n_2(v_2) \sigma(E) |v| d^3v_1 d^3v_2.$$

The factor  $(1 + \delta_{12})^{-1}$  is equal to  $1/2$  if particles 1 and 2 are identical, and is there to avoid double-counting in that case. To get the total reaction rate per unit time, we need to integrate over the joint velocity distribution  $d^3v_1 d^3v_2$ ,

$$r_{12} = \frac{n_1 n_2}{1 + \delta_{12}} \left[ \frac{m_1 m_2}{(2\pi kT)^2} \right]^{3/2} \times \int \sigma(E) v \exp \left( -\frac{m_1 v_1^2}{2kT} - \frac{m_2 v_2^2}{2kT} \right) d^3v_1 d^3v_2. \quad (10.16)$$

Now  $E$  and  $v$  are the relative energies and velocity in the center-of-mass frame. We can change variable using the relations

$$\begin{aligned} v_1 &= V - \frac{m_2}{m_1 + m_2} v \\ v_2 &= V + \frac{m_1}{m_1 + m_2} v. \end{aligned}$$

where  $V$  is the center-of-mass velocity. It is straightforward to show that  $dV_{1,x} dv_{2,x} = dV_x dv_x$ , and likewise for the  $y, z$  directions. Furthermore,  $m_1 v_1^2 + m_2 v_2^2 = (m_1 + m_2)V^2 + mv^2$ , and multiplying and dividing the integral in equation (10.16) by  $m_1 + m_2$  allows us to write

$$r_{12} = \frac{n_1 n_2}{1 + \delta_{12}} \left( \frac{m_1 + m_2}{2kT} \right)^{3/2} \left( \frac{m}{2kT} \right)^{3/2} \times \int d^3V \int d^3v \sigma(E) v \exp \left[ -\frac{mv^2}{2kT} \right] \exp \left[ -\frac{(m_1 + m_2)V^2}{2kT} \right].$$

The integral over  $d^3V$  can be factored out and is normalized to unity. Hence we have for the reaction rate between a pair of particles 1 and 2,

$$\begin{aligned} r_{12} &= \frac{1}{1 + \delta_{12}} n_1 n_2 \left\{ \left( \frac{m}{2\pi kT} \right)^{3/2} \int_0^\infty \sigma(E) v \exp \left( -\frac{mv^2}{2kT} \right) 4\pi v^2 dv \right\}. \\ &\equiv \frac{1}{1 + \delta_{12}} n_1 n_2 \langle \sigma v \rangle. \end{aligned} \quad (10.17)$$

The term in  $\{\}$  is the averaging over the joint distribution of the cross-section times the velocity, and is usually denoted as  $\langle \sigma v \rangle$ .

Changing variables to  $E = mv^2/2$  in equation (10.17) and inserting the formula for the cross-section, equation (10.14), gives

$$\langle \sigma v \rangle = \left( \frac{8}{\pi m} \right)^{1/2} \left( \frac{1}{kT} \right)^{3/2} \int_0^\infty S(E) \exp \left[ - \left( \frac{E_G}{E} \right)^{1/2} - \frac{E}{kT} \right] dE. \quad (10.18)$$

Now, we've assumed that  $S(E)$  varies slowly; but look at the argument of the exponential. This is a competition between a rapidly rising term  $\exp[-(E_G/E)^{1/2}]$  and a rapidly falling term  $\exp(-E/kT)$ .

As a result, the exponential will have a strong peak, and we can expand the integrand in a Taylor series about the maximum. Let

$$f(E) = -\left(\frac{E_G}{E}\right)^{1/2} - \frac{E}{kT}.$$

Then we can write

$$\begin{aligned} & \int_0^\infty S(E) \exp \left[ -\left(\frac{E_G}{E}\right)^{1/2} - \frac{E}{kT} \right] dE \\ & \approx \int_0^\infty S(E_{\text{pk}}) \exp \left[ f(E_{\text{pk}}) + \frac{1}{2} \left. \frac{d^2 f}{dE^2} \right|_{E=E_{\text{pk}}} (E - E_{\text{pk}})^2 \right]. \end{aligned}$$

Here  $E_{\text{pk}}$  is found by solving  $(df/dE)|_{E=E_{\text{pk}}} = 0$ . This trick allows us to turn the integral into a Gaussian! (Before the internet, all there was to do for fun were integrals.)

Solving for  $E_{\text{pk}}$ , we get

$$E_{\text{pk}} = \frac{E_G^{1/3} (kT)^{2/3}}{2^{2/3}},$$

and

$$\exp [f(E_{\text{pk}})] = \exp \left[ -3 \left( \frac{E_G}{4kT} \right)^{1/3} \right].$$

Further,

$$\frac{1}{2} \left. \frac{d^2 f}{dE^2} \right|_{E=E_{\text{pk}}} = -\frac{3}{2(2E_G)^{1/3} (kT)^{5/3}} = -\frac{3}{4E_{\text{pk}} kT}.$$

Defining a variable  $\Delta = 4(E_{\text{pk}} kT / 3)^{1/2}$ , our integral becomes

$$\begin{aligned} \langle \sigma v \rangle &= \left( \frac{8}{\pi m} \right)^{1/2} \left( \frac{1}{kT} \right)^{3/2} \\ &\times S(E_{\text{pk}}) \exp \left[ -3 \left( \frac{E_G}{4kT} \right)^{1/3} \right] \int_0^\infty \exp \left[ -\frac{(E - E_{\text{pk}})^2}{(\Delta/2)^2} \right] dE. \end{aligned} \quad (10.19)$$

How well does this approximation do? Figure 10.4 shows the integrand (*solid line*) and the approximation by a Gaussian (*dashed line*). Although the integrand is skewed to the right, the area is approximately the same. We could correct for this by taking more terms in our expansion. Consult Clayton for details.

Another simplification can be made because both the Gaussian and the original integrand go to zero as  $E \rightarrow 0$ . As a result, we can extend the lower bound of our integral (eq. [10.19]) to  $-\infty$ , and obtain

$$\begin{aligned} \langle \sigma v \rangle &\approx \left( \frac{8}{\pi m} \right)^{1/2} \left( \frac{1}{kT} \right)^{3/2} S(E_{\text{pk}}) \exp \left[ -3 \left( \frac{E_G}{4kT} \right)^{1/3} \right] \frac{\Delta}{2} \\ &= \frac{2^{13/6}}{\sqrt{3m}} \frac{E_G^{1/6}}{(kT)^{2/3}} \exp \left[ -3 \left( \frac{E_G}{4kT} \right)^{1/3} \right] S(E_{\text{pk}}). \end{aligned} \quad (10.20)$$

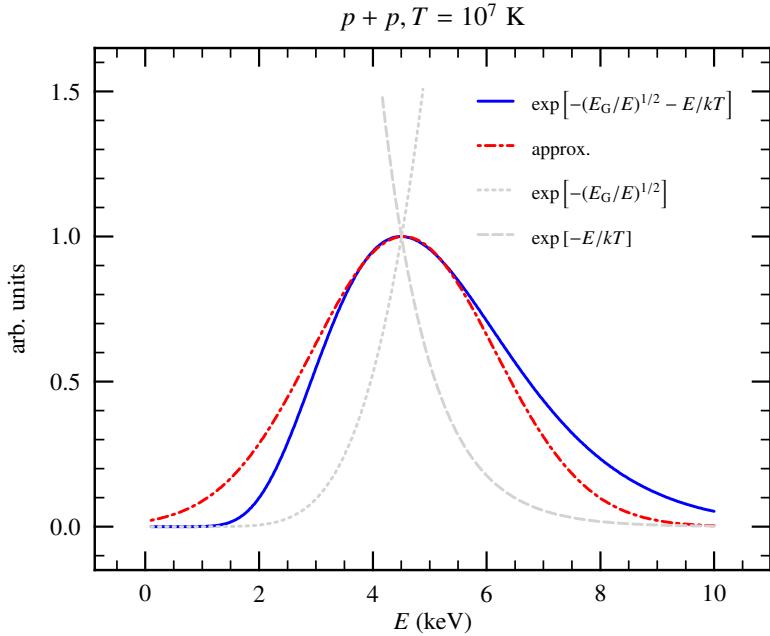


Figure 10.4: Integrand of eq. (10.18) (solid line) and the Gaussian (dot-dashed line) constructed by expanding to second order the argument of the exponential. The parameters for  $E_G$  were taken from the  $p + p$  reactions ( $Z_1 Z_2 = 1$ ,  $A = 1/2$ ), and the temperature is  $10^7$  K. Note that the grey curves, showing the two terms of the exponential, have been rescaled to fit on the same plot.

On to some numbers. Table 10.2 lists quantities for some common reactions. A couple of notes. First,  $\Delta/E_{\text{pk}}$  indicates how well our Gaussian approximation works—you will see it is less than 1 in all cases. We evaluated  $\Delta/E_{\text{pk}}$ , which decreases with temperature as  $T^{-1/6}$ , at  $T = 10^7$  K. Second, the quantity  $n(T)$  is the exponent if we want to approximate the reaction rate as a power-law,  $r \propto T^n$ . We compute this as

$$n(T) = \frac{d \ln r}{d \ln T} = -\frac{2}{3} + \left( \frac{E_G}{4kT} \right)^{1/3}, \quad (10.21)$$

as you can easily verify for yourself. In the table, the exponent is evaluated at  $T = 10^7$  K; obviously  $n$  depends on temperature. Finally, note the size of  $E_G/(4k)$ . This makes the argument of the exponential in equation (10.20) large in absolute value, and sets the temperature scale at which a given reaction comes into play.

#### 10.4 Resonances

This section contains my condensed notes on resonances following Blatt and Weisskopf's excellent text<sup>5</sup>. Other treatments of the subject, such as that in Iliadis<sup>6</sup> and in Clayton<sup>7</sup>, mostly follow their approach. In this section, we shall often make use of the notation  $X(a, b)Y$  to mean the reaction  $X + a \rightarrow b + Y$ .

<sup>5</sup> John M. Blatt and Victor F. Weisskopf. *Theoretical Nuclear Physics*. Springer-Verlag, 1979

<sup>6</sup> C. Iliadis. *Nuclear Physics of Stars*. Wiley-VCH, 2007

<sup>7</sup> Donald D. Clayton. *Principles of Stellar Evolution and Nucleosynthesis*. University of Chicago Press, 1983

Reaction	$p + p$	$p + {}^3\text{He}$	${}^3\text{He} + {}^3\text{He}$	$p + {}^7\text{Li}$	$p + {}^{12}\text{C}$
$A$	1/2	3/4	3/2	0.88	0.92
$Z_1 Z_2$	1	2	4	3	6
$E_G$ (MeV)	0.489	2.94	23.5	7.70	32.5
$E_G / (4k)$ (GK)	1.4	8.5	68.0	22.0	94.0
$E_{\text{pk}} _{T=10^7 \text{ K}}$ (keV)	4.5	8.2	16.3	11.3	18.2
$\Delta / E_{\text{pk}} _{T=10^7 \text{ K}}$	1.0	0.75	0.53	0.64	0.50
$n(T = 10^7 \text{ K})$	4.6	8.8	18.3	12.4	20.5

Table 10.2: Parameters for non-resonant reactions

### Orbitals

Nuclei exhibit shell effects: one can often treat the nucleons as independent particles occupying orbitals determined by a mean force. Unlike in the atomic case, the spin-orbit term in the Hamiltonian,  $-a\mathbf{L} \cdot \mathbf{S}$ , is quite strong. Since the total angular momentum is  $\mathbf{J} = \mathbf{L} + \mathbf{S}$ , we have

$$\mathbf{L} \cdot \mathbf{S} = \frac{1}{2} (\mathbf{J} \cdot \mathbf{J} - \mathbf{L} \cdot \mathbf{L} - \mathbf{S} \cdot \mathbf{S}),$$

and hence states with larger  $\mathbf{J}$  have a lower energy. The strong  $\mathbf{L} \cdot \mathbf{S}$  coupling leads to the presence of “gaps” in the energy spectra; nuclei that have filled (either neutrons or protons) shells up to this gap are unusually bound and the nucleon number (either neutron or proton) is termed a *magic number*. The magic numbers are 2, 8, 20, 28, 50, 82, and 126. For example,  ${}^{16}\text{O}$  (8 protons, 8 neutrons) and  ${}^{40}\text{Ca}$  (20 protons, 20 neutrons) are doubly magic and hence more strongly bound than other nuclides of similar mass.

We label the orbitals as  $n\ell_j$ , where  $n$  is the radial quantum number,  $\ell$  is the orbital angular momentum ( $s, p, d, f, \dots$ ), and  $j$  is the total angular momentum. The first few orbitals are listed in Table 10.3. Each orbital has  $2j + 1$  nucleons, and a fully occupied orbital has  $J = 0$ . For example, we would expect that the ground state of  ${}^{13}\text{C}$  (6 protons, 7 neutrons) to have closed  $1s_{1/2}$  and  $1p_{3/2}$  shells for both neutrons and protons, and the remaining neutron would then occupy the  $1p_{1/2}$  shell.

The remaining quantum number is parity ( $\pi$ ), which is conserved under the strong force. The parity of a nucleon orbital is  $(-1)^\ell$ , where the angular momentum number  $\ell$  must be summed over all nucleons. Since a closed shell has an even number of nucleons, it must have positive parity. For example, the ground state of  ${}^{17}\text{O}$  has 8 protons filling the  $1s_{1/2}$ ,  $1p_{3/2}$ , and  $1p_{1/2}$  shells 8 neutrons in closed shells, so the remaining neutron must be in the  $1d_{5/2}$  shell ( $\ell = 2$ ): the angular momentum and parity of the ground state must there-

Orbital	Number, $2j + 1$ , in orbital	Total number
$1d_{3/2}$	4	20
$2s_{1/2}$	2	16
$1d_{5/2}$	6	14
$1p_{1/2}$	2	8
$1p_{3/2}$	4	6
$1s_{1/2}$	2	2

Table 10.3: Neutron orbitals

fore be  $j^\pi = \frac{5}{2}^+$ . As a second example,  $^{14}\text{N}$  has 6 protons in closed shells and 6 neutrons in closed shells, with the remaining proton and neutron both in  $1p_{1/2}$  orbitals. Hence the angular momentum of the ground state could either be 0 or 1; it turns out that the symmetric state has lower energy, so  $j = 1$ . The parity of the ground state is  $(-1)^{1+1} = 1$ , so  $J^\pi(^{14}\text{N}) = 1^+$ .

The angular momentum matters because it sets the possible range of relative angular momenta that the incoming particles can have. Writing the wave function as  $\psi(r) = (u_\ell(r)/r)Y_{\ell 0}$  and substituting into Schrödinger's equation,

$$-\frac{\hbar^2}{2m}\nabla^2\psi + V(r)\psi = E\psi,$$

gives the following equation for  $u_\ell$ ,

$$\frac{d^2}{dr^2}u_\ell + \left\{ k^2 - \frac{\ell(\ell+1)}{r^2} - \frac{2mZ_aZ_Xe^2}{\hbar^2} \right\} u_\ell = 0.$$

Here  $k$  is the wavenumber for the particle at  $r \rightarrow \infty$ ,  $E = \hbar^2k^2/2m$ . Note the presence of the "centrifugal barrier,"

$$\frac{\hbar^2}{2mr^2}\ell(\ell+1) = 20.9 \text{ MeV} \frac{m_u}{m} \left( \frac{\text{fm}}{r} \right)^2 \ell(\ell+1) :$$

there is a price to pay if the particles must have a high relative  $\ell$ .

### *Formation of a compound nucleus*

Even orbitals with positive energy can be long-lived; suppose we excite a proton in  $^{11}\text{C}$  to an orbital just above the threshold for decay into  $\text{p} + ^{10}\text{B}$ . Although this state has positive energy and can decay, the particle has to tunnel through the coulomb barrier, and potentially an angular momentum barrier if  $s$ -wave emission is forbidden. We saw that the probability of getting through the Coulomb barrier (for  $s$ -wave) is given by eq. (10.12). If this is very small, then we can imagine a classical particle oscillating back and forth in the well, which is does many times because the probability of escaping each time it approaches the barrier is so small. Thus, if there is a

substantial energy barrier impeding escape, the classical oscillation period  $P$  is much less than the lifetime of the state  $\tau$ . Now the classical oscillation period depends inversely on the spacing  $D$  between energy levels,  $P \sim \hbar/D$ , as you can verify for an infinite square well potential. Hence if the probability of tunneling is sufficiently small,

$$\frac{P}{\tau} \sim \frac{\hbar}{D\tau} = \frac{\Gamma}{D} \ll 1,$$

where the width of the state to particle emission is  $\Gamma = \hbar/\tau$ . Hence for reactions at low excitation energies involving light nuclei with widely separated levels, we can look at captures into discrete levels in the compound nucleus.

### *Derivation of resonant cross-section*

We're now ready to do some heavy lifting and derive the cross-section for a resonant reaction. The nuclear force is short-ranged, so we can define a "channel radius"  $R$  exterior to which our potential is purely Coulomb. Our strategy is just like doing transmission resonances in quantum mechanics: we'll solve for the wave function exterior to  $R$  and match it to the wave function inside  $R$ . Since we don't completely know the form of the potential inside  $R$ , our expression will have terms that must be experimentally constrained.

The reaction is  $a + X$ ; the coulomb potential is  $Z_a Z_X e^2/r$  and the relative angular momentum of  $a$  and  $X$  is  $\ell$ . Let  $m$  be the reduced mass of  $a$  and  $X$ . Our wave function is then  $\psi(r) = (u_\ell(r)/r)Y_{\ell 0}$ ; substituting this into the Schrödinger equation,

$$-\frac{\hbar^2}{2m}\nabla^2\psi + V(r)\psi = E\psi,$$

gives the following equation for  $u_\ell$ ,

$$\frac{d^2}{dr^2}u_\ell + \left\{ k^2 - \frac{\ell(\ell+1)}{r^2} - \frac{2m}{\hbar^2} \frac{Z_a Z_X e^2}{r} \right\} u_\ell = 0.$$

Here  $k$  is the wavenumber for the particle at  $r \rightarrow \infty$ ,  $E = \hbar^2 k^2 / 2m$ . There are two solutions to this differential equation; the solutions are known as Coulomb wave functions,  $F_\ell$  and  $G_\ell$ . The regular solution  $F_\ell$  vanishes as  $r \rightarrow 0$ , and  $G_\ell$  blows up at the origin. For  $kr \gg 1$ , the Coulomb wave functions go over to

$$F_\ell(r) \simeq \sin \left[ kr - \frac{1}{2}\ell\pi - \gamma \ln(2kr) + \sigma_\ell \right] \quad (10.22)$$

$$G_\ell(r) \simeq \cos \left[ kr - \frac{1}{2}\ell\pi - \gamma \ln(2kr) + \sigma_\ell \right]. \quad (10.23)$$

Here the parameter  $\sigma_\ell$  is a Coulomb phase shift, and the parameter

$$\gamma = \frac{1}{2\pi} \left( \frac{E_G}{E} \right)^{1/2}$$

contains the Gamow energy. This shouldn't be too surprising: since  $F_\ell$  and  $G_\ell$  are the exact solutions to the motion of a particle in a Coulomb potential, they must behave like what we found using a WKB approximation in some limit.

For  $r \geq R$ , we can write our solution in terms of outgoing waves,  $u_\ell^+$ , and incoming waves,  $u_\ell^-$ ; here the outgoing and incoming waves are defined as

$$\begin{aligned} u_\ell^+ &= e^{-i\sigma_\ell} [G_\ell + iF_\ell] \\ u_\ell^- &= e^{i\sigma_\ell} [G_\ell - iF_\ell]. \end{aligned}$$

At large distances these go over to plane waves, and so we can determine the coefficients,

$$u_\ell = \frac{\sqrt{\pi}}{k} r^{2\ell+1} (2\ell+1)^{1/2} [u_\ell^- - \eta_\ell u_\ell^+]. \quad (10.24)$$

The effect of the nucleus is to affect the outgoing waves via the coefficient  $\eta_\ell$ . Furthermore, if we integrate the current over a large sphere we obtain the cross-section for the reaction (if the current is zero, then every particle that enters the sphere leaves it, so there is no reaction),

$$\sigma_\ell = \frac{\pi}{k^2} (2\ell+1) [1 - |\eta_\ell|^2]. \quad (10.25)$$

To determine  $\eta_\ell$ , we need to find the logarithmic derivative

$$\alpha_\ell \equiv R \left. \frac{u'_\ell}{u_\ell} \right|_{r=R}$$

where  $u'_\ell = du_\ell/dr$ . For now  $\alpha_\ell$  is undetermined, since it depends on the wave function inside the nucleus. It is useful to define the following quantities:

$$\frac{u_\ell^-}{u_\ell^+} = \frac{G_\ell - iF_\ell}{G_\ell + iF_\ell} e^{2i\sigma_\ell} \equiv e^{2i\xi} \quad (10.26)$$

$$R \left. \frac{u'_\ell}{u_\ell^+} \right|_{r=R} = R(G'_\ell G_\ell + F'_\ell F_\ell)v_\ell + ikRv_\ell \equiv \Delta_\ell + is_\ell. \quad (10.27)$$

In the last equation, we made use of the fact that  $G_\ell F'_\ell - F_\ell G'_\ell = k$  for all  $r$  and defined

$$v_\ell = \left( G_\ell^2 + F_\ell^2 \right)_{r=R}^{-1}.$$

To see the significance of  $v_\ell$ , note that the ratio of an outgoing wave at  $r \rightarrow \infty$  and that at  $r = R$  is

$$\frac{|u_\ell(\infty)|^2}{|u_\ell(R)|^2} = \frac{1}{G_\ell^2(R) + F_\ell^2(R)},$$

where the numerator follows from the asymptotic forms, eqns. (10.22) and (10.23). This ratio, however, is just the probability of the wave

transmitting through the potential barrier, and for  $\ell = 0$  is approximately what we found earlier (eq. [10.12]) using the WKB approximation.

Evaluating  $\alpha_\ell$  for the solution defined at  $r \geq R$ , eq. (10.24), and using our definitions, eq. (10.26) and (10.27), we determine the phase shift,

$$\eta_\ell = \frac{\alpha_\ell - \Delta_\ell + is_\ell}{\alpha_\ell - \Delta_\ell - is_\ell} e^{2i\zeta}.$$

Using this to evaluate the reaction cross-section, eq. (10.25), we have

$$\sigma_\ell = \frac{\pi}{k^2} (2\ell + 1) \frac{-4s_\ell \Im \alpha_\ell}{(\Re \alpha_\ell - \Delta_\ell)^2 + (\Im \alpha_\ell - s_\ell)^2} \quad (10.28)$$

Notice that the reaction cross-section vanishes if  $\alpha_\ell \in \mathbb{R}$ , i.e.,  $\Im \alpha_\ell = 0$ . So far, our efforts may just look like we are reshuffling terms for no apparent reason, but there is a method to the algebraic madness. We see explicitly, for example, that the cross-section is proportional to the penetration through the term  $s_\ell$  in the numerator.

To make further progress, we have to evaluate  $\alpha_\ell$ , and this requires making some constraints on the form of the wavefunction at  $r < R$ . Although the precise form of the potential is unknown, we do know that it is a rather deep well. Just inside the surface  $R$ , we expect the radial wavefunction to be composed of spherical waves,

$$u_\ell(r < R) = A \left\{ \exp(-iKr) + e^{2i\zeta} e^{-2q} \exp(iKr) \right\}.$$

Our reasoning for this form is as follows. In general the state will be a standing wave, which we can write as a sum of incoming and outgoing waves. The jump in the potential at  $r = R$  introduces a phase shift, which we parameterize by  $e^{2i\zeta}$ . If the state decays by another channel than simply re-emitting the particle, i.e. if a reaction occurs, then the amplitude of the outgoing wave will be less than that of the incoming wave; we parameterize this by the factor  $e^{-2q}$ . We expect that  $q \ll 1$ , because otherwise the state wouldn't be long-lasting and have a well-defined energy.

We can factor  $u_\ell(r < R)$ ,

$$\begin{aligned} u_\ell(R) &= \left( 2Ae^{i\zeta} e^{-q} \right) \left\{ \frac{\exp(-iKR - i\zeta + q) + \exp(iKR + i\zeta - q)}{2} \right\} \\ &= C \cos(KR + \zeta + iq). \end{aligned}$$

Using this to evaluate the logarithmic derivative,

$$\alpha_\ell = -KR \tan(KR + \zeta + iq). \quad (10.29)$$

Now, recall that  $K \gg k$ ; the wavenumber inside the nuclear potential well is much larger than that outside. The only way to smoothly join two waves with such discrepant wavenumbers is to have them match

where  $u'_\ell = 0$ , i.e., where  $\alpha_\ell = 0$ . Let's expand  $\alpha_\ell$  about such a point with energy  $\varepsilon_r$  and  $q = 0$ :

$$\alpha_\ell \approx \frac{\partial \alpha_\ell}{\partial \varepsilon} \Big|_{q=0, \varepsilon=\varepsilon_r} (\varepsilon - \varepsilon_r) + \frac{\partial \alpha_\ell}{\partial q} \Big|_{q=0, \varepsilon=\varepsilon_r} q.$$

Substituting this into equation (10.28) gives

$$\sigma = \frac{\pi}{k^2} (2\ell + 1) \left\{ \frac{4KRqs_\ell}{[(\partial_\varepsilon \alpha_\ell)(\varepsilon - \varepsilon_r) - \Delta_\ell]^2 + [-KRq - s_\ell]^2} \right\}. \quad (10.30)$$

We expect that near a resonance, the cross-section will have a Lorentzian profile, with a total width

$$\Gamma = \sum_i \Gamma_i + \Gamma_\gamma$$

that is the sum of particle decay widths  $\Gamma_i$  and widths for radiative transitions  $\Gamma_\gamma$ . The entrance channel width will be proportional to  $s_\ell = kRv_\ell$ . Furthermore, a reaction will take place if the nucleus does something other than decay in the entrance channel, and this will occur with probability  $\Gamma - \Gamma_a$ .

Motivated by these considerations, we note that if we define the width for decay in the entrance channel as

$$\Gamma_a = -\frac{2s_\ell}{\partial_\varepsilon \alpha_\ell} = -\frac{2kRv_\ell}{\partial_\varepsilon \alpha_\ell}, \quad (10.31)$$

the reaction width as (it can be shown that in general  $\partial_\varepsilon \alpha_\ell < 0$ )

$$\Gamma_r = \Gamma - \Gamma_a = -\frac{2KRq}{\partial_\varepsilon \alpha_\ell}, \quad (10.32)$$

and the "observed" resonance energy as

$$\varepsilon_{r,\text{obs}} = \varepsilon_r + \frac{\Delta_\ell}{\partial_\varepsilon \alpha_\ell},$$

then the cross-section becomes

$$\sigma = \frac{\pi}{k^2} (2\ell + 1) \frac{\Gamma_a \Gamma_r}{(\varepsilon - \varepsilon_{r,\text{obs}})^2 + (\Gamma/2)^2}. \quad (10.33)$$

This is the cross-section for a compound nucleus to form in channel  $a$  and decay by any other channel. To get the cross-section for a specific exit channel  $b$ , we must multiply by the branching ratio  $\Gamma_b/\Gamma_r$ . Finally, for an unpolarized incident beam, we need to multiply the cross section by a statistical factor

$$\omega = \frac{2J+1}{(2J_a+1)(2J_X+1)} \quad (10.34)$$

to account for the fraction of angular momentum states in the target  $X$  and beam  $a$  that can enter the level with the appropriate angular momentum  $\ell$ .

You might be troubled by our identification of the particle width, eq. (10.31), and reaction width, eq. (10.32). Notice that if we were to solve the problem of a quasi-bound state leaking out of its well, we would encounter similar equations, and this would lead to the identification of the level width. We can make the argument more plausible by the following consideration. Let's consider  $\partial_\epsilon \alpha_\ell$ , evaluated about the point where  $\alpha_\ell = 0$  and  $q = 0$ . The general form of  $\alpha_\ell$  is a tangent function (eq. [10.29]), so if we increase the energy by the spacing between levels  $D$ , the phase of the tangent must increase by  $\pi$ . Hence,  $\partial_\epsilon \alpha_\ell|_{KR+\zeta=n\pi,q=0} \sim -KR\pi/D$ . Substituting this into eq. (10.31) and rearranging,

$$\Gamma_a \sim \hbar \left( \frac{4k}{K} \right) \left( \frac{D}{2\pi\hbar} \right) v_\ell.$$

Now, when a plane wave is incident on a step in the potential, the transmission coefficient across the step is roughly  $k/K$ , as you can verify by solving a one-dimensional Schrödinger equation. We saw earlier that classical oscillation period for a particle in a well is  $2\pi\hbar/D$ . Hence

$$\begin{aligned} \Gamma_a &\sim \hbar \times (\text{oscillation frequency}) \\ &\times (\text{transmission across potential step at nuclear surface}) \\ &\times (\text{probability of penetrating coulomb, centrifugal barrier}). \end{aligned}$$

But this is precisely what we would write down for  $\hbar \times (\text{rate of decay in channel } a)$ —the particle has a small probability on each oscillation to penetrate the potential jump and barrier.

### *A worked example*

Consider the reaction  ${}^{10}\text{B}(p, \alpha){}^7\text{Be}$ . We can think of this reaction proceeding first via the formation of a compound nucleus in an excited state,  ${}^{11}\text{C}^*$ . It can be shown that the cross-section for the formation of  ${}^{11}\text{C}^*$  via proton capture is proportional to the width for the decay of that state via proton emission:

$$\sigma(p + {}^{10}\text{B} \rightarrow {}^{11}\text{C}^*) \propto \Gamma_p. \quad (10.35)$$

The second stage of the reaction is the decay of the  ${}^{11}\text{C}^*$  into  $\alpha + {}^7\text{Be}$ . Because of the short timescale  $\sim 10^{-23}$  s for nuclear interactions—roughly the crossing time for a nucleon in a 20 MeV by 2 fm well—we make the assumption that the decay of an excited state does not depend on how the state was formed. If  $\Gamma_\alpha$  represents the decay  ${}^{11}\text{C}^* \rightarrow \alpha + {}^7\text{Be}$ , then

$$\sigma(p + {}^{10}\text{B} \rightarrow \alpha + {}^7\text{Be}) \propto \Gamma_p \Gamma_\alpha.$$

The full expression for the cross section is

$$\sigma(p + {}^{10}\text{B} \rightarrow \alpha + {}^7\text{Be}) = \frac{\pi}{k^2} (2\ell + 1) \omega \frac{\Gamma_p \Gamma_\alpha}{(\varepsilon - \varepsilon_{r,\text{obs}})^2 + (\Gamma/2)^2}. \quad (10.36)$$

Here the first term  $\pi(2\ell + 1)k^{-2}$  is the geometrical cross-section and  $\Gamma$  represents the total decay width of the excited state in the compound nucleus. The factor  $\omega$  is the statistical factor

$$\omega = \frac{2J + 1}{(2J_p + 1)(2J_{10} + 1)}$$

that accounts for the fraction of angular momentum states that can enter the level with the appropriate angular momentum  $\ell$ . For our example, the reaction  ${}^{10}\text{B}(p, \alpha){}^7\text{Be}$  can proceed via the  $J^\pi = 5/2^+$  level at 8.70 MeV in  ${}^{11}\text{C}$ . The angular momentum and parity of  ${}^{10}\text{B}$  and the proton are  $J^\pi({}^{10}\text{B}) = 3^+$ ,  $J^\pi(p) = 1/2^+$ . These two spins can add to either  $7/2^+$  (multiplicity of 8) or  $5/2^+$  (multiplicity of 6). As a result, they can enter into the resonance with  $\ell = 0$ , and they will do so with 6/14 probability. The cross section for this reaction is then

$$\sigma(p + {}^{10}\text{B} \rightarrow {}^7\text{Be} + \alpha) = \frac{6}{14} \frac{\pi}{k^2} \frac{\Gamma_p \Gamma_\alpha}{(\varepsilon - \varepsilon_{8.70})^2 + (\Gamma/2)^2}.$$

This is multiplied by the entrance channel velocity  $v$  and integrated over the thermal distribution to produce the reaction rate.

### 10.5 Inverse Rates

Consider a photodisintegration reaction  $Y(\gamma, a)X$ . We could compute the rate for this by first computing the excitation  $Y \rightarrow Y^*$  by absorption of a photon followed by decay through the  $a$ -channel. There is another, easier, way to compute the thermally averaged rate, however, if we already have an expression for the forward reaction  $X(a, \gamma)Y$ . Suppose we allow our plasma, consisting of  $a$ ,  $X$ , and  $Y$  to come into thermal balance. In such a plasma, the composition does not change with time, so our forward and inverse rates must balance:

$$n_X n_a \langle \sigma v \rangle_{Xa} = n_Y \lambda, \quad (10.37)$$

where  $\lambda$  is the photodissociation rate. Since we are in thermal equilibrium, however, we also have a relation between the chemical potentials,

$$\mu_X + \mu_a = \mu_Y.$$

For an ideal gas, the chemical potentials are given by eq. (6.7). Rearranging terms gives and taking the exponential gives us the equation

$$\frac{n_X n_a}{n_Y} = \frac{n_X^Q n_a^Q}{n_Y^Q} \exp\left(-\frac{Q}{k_B T}\right). \quad (10.38)$$

In this expression we have substituted  $Q = (m_a + m_X - m_Y)c^2$ , and

$$n_i^Q = g_i \left( \frac{m_i k_B T}{2\pi\hbar^2} \right)^{3/2}.$$

Substituting eq. (10.38) into eq. (10.37) gives us an expression for the photodissociation rate  $\lambda$  in terms of the forward rate,

$$\begin{aligned} \lambda &= \frac{n_X n_a}{n_Y} \langle \sigma v \rangle_{Xa} \\ &= \frac{g_X g_a}{g_Y} \left( \frac{m_u k_B T}{2\pi\hbar^2} \right)^{3/2} \left( \frac{A_X A_a}{A_Y} \right)^{3/2} \exp \left( -\frac{Q}{k_B T} \right) \langle \sigma v \rangle_{Xa} \end{aligned} \quad (10.39)$$

This expression is specific to this type of reaction, but similar formula can be generated for other types of reactions, such as  $X(a, b)Y$ .

## 10.6 Plasma corrections to the reaction rate

In stars, the nuclear reactions do not occur in isolation, but rather in the midst of a plasma. The effect of these ambient charges is to screen the long-range Coulomb interaction. This in turn perturbs the penetration factors and hence the reaction rates. We now derive the lowest order correction<sup>8</sup>.

The first thing to consider are the typical scales involved. In a plasma, the mean interparticle spacing is

$$a = \left( \frac{3Am_u}{4\pi\rho} \right)^{1/3} = 1.6 \times 10^4 \text{ fm} \times A^{1/3} \left( \frac{100 \text{ g cm}^{-3}}{\rho} \right)^{1/3},$$

and the Debye length (cf. eq. [6.41]) is

$$\lambda_D = 2.8 \times 10^4 \text{ fm} \times \left( \frac{T}{10^7 \text{ K}} \right)^{1/2} \left( \frac{100 \text{ g cm}^{-3}}{\rho} \right)^{1/2} \left( \frac{A}{\langle Z^2 + Z \rangle} \right)^{1/2}.$$

Both of these lengths, for typical conditions in a stellar plasma, are much larger than the size of the classically forbidden region through which the particle must tunnel: at  $E_{\text{pk}} = 10 \text{ keV}$ , this length is  $r_E = 144 Z_1 Z_2 \text{ fm} \ll a < \lambda_D$ . The nuclear scale is  $\sim \text{ fm}$  and is much smaller than all of these.

The penetration factor (eq. [10.11]) depends on the potential in the barrier; given that  $r_E \ll \lambda_D$ , we may expand the potential, eq. (6.42), and write

$$\mathcal{P} \propto \exp \left\{ \frac{2}{\hbar} \int_{r_E}^{r_N} \left[ 2m \left( \frac{Z_1 Z_2 e^2}{r} - \frac{Z_1 Z_2 e^2}{\lambda_D} - E \right) \right]^{1/2} dr \right\}. \quad (10.40)$$

Since  $\lambda_D$  doesn't depend on either  $E$  or  $r$ , the effect of the screening potential is just to change the zero point of the energy scale. Assuming the S-factor doesn't depend sensitively on energy, the effect

<sup>8</sup> E. E. Salpeter. Electrons Screening and Thermonuclear Reactions. *Australian Journal of Physics*, 7:373–+, September 1954

of screening on the rate is just to multiply the integrand in equation (10.17) by

$$f_{\text{scr}} = \exp\left(\frac{Z_1 Z_2 e^2}{\lambda_D k_B T}\right). \quad (10.41)$$

Since this factor doesn't depend on  $E$ , it simply multiplies the rate  $\langle\sigma v\rangle$ . Note that since  $\lambda_D > a$  and since  $e^2/(ak_B T) \ll 1$  in a plasma, this screening factor is a small correction to the rate.

## 10.7 Equations for Chemical Evolution

Now that we have a reaction cross-section, we can write down the equations describing the chemical evolution of the star, and the nuclear heating. To make this concrete, let's look at the reaction  $p + p \rightarrow e^+ \nu_e + {}^2\text{H}$ . The reaction rate is (cf. eqs. [10.17] and [10.41])

$$r_{pp} = \frac{1}{2} n_H^2 f_{pp}^{\text{scr}} \langle\sigma v\rangle_{pp}.$$

Each reaction destroys 2 protons, so we can write down our equation for the change in  $n_H$ ,

$$\partial_t n_H + \nabla \cdot (n_H \mathbf{u}) = -n_H^2 f_{pp}^{\text{scr}} \langle\sigma v\rangle_{pp}. \quad (10.42)$$

Likewise, the  $pp$  reaction produces deuterium, which is in turn destroyed by  $D + p \rightarrow {}^3\text{He}$ :

$$\partial_t n_D + \nabla \cdot (n_D \mathbf{u}) = \frac{1}{2} n_H^2 f_{pp}^{\text{scr}} \langle\sigma v\rangle_{pp} - n_D n_H f_{Dp}^{\text{scr}} \langle\sigma v\rangle_{Dp}. \quad (10.43)$$

Let's write each number density in terms of its abundance:  $n_H = Y_H N_A \rho$ . We can then use the equation of mass continuity (eq. [3.1]) to simplify these equations to

$$(\partial_t + \mathbf{u} \cdot \nabla) Y_H = -Y_H^2 \rho f_{pp}^{\text{scr}} [N_A \langle\sigma v\rangle]_{pp} \quad (10.44)$$

$$\begin{aligned} (\partial_t + \mathbf{u} \cdot \nabla) Y_D &= \frac{1}{2} Y_H^2 \rho f_{pp}^{\text{scr}} [N_A \langle\sigma v\rangle]_{pp} \\ &\quad - Y_D Y_H \rho f_{Dp}^{\text{scr}} [N_A \langle\sigma v\rangle]_{Dp}. \end{aligned} \quad (10.45)$$

The left-hand side of these equations are just the Lagrangian time derivatives. On the right-hand sides, the quantities in  $[ ]$  are just functions of temperature and are compiled into rate libraries, such as REACLIB. We will need equations like (10.44) and (10.45) for each species in our star. A collection of such equations is known as a **reaction network**.

Each reaction is specified by a  $Q$ -value, which is just the energy deposited into the gas by the reactions. If there is no neutrino released, this will just be the change in nuclear binding energy, but for reaction like  $p + p$  one has to account for the energy carried off by

the neutrino. For the two reactions we consider here, the total heating rate, per unit mass, is

$$q = \frac{1}{\rho} \left( \frac{Q_{pp}}{2} n_H^2 f_{pp}^{\text{scr}} \langle \sigma v \rangle_{pp} + Q_{DP} n_D n_H f_{Dp}^{\text{scr}} \langle \sigma v \rangle_{Dp} \right) \quad (10.46)$$

$$\begin{aligned} &= \frac{Q_{pp} N_A}{2} Y_H^2 \rho f_{pp}^{\text{scr}} [N_A \langle \sigma v \rangle]_{pp} \\ &\quad + Q_{Dp} N_A Y_D Y_H \rho f_{Dp}^{\text{scr}} [N_A \langle \sigma v \rangle]_{Dp}. \end{aligned} \quad (10.47)$$

The equations for the change in chemical composition and the heating rate close our system of equations describing the structure and evolution of the star. We are now ready to discuss the life cycle of stars in detail.



# 11

## *Star Formation*

Stars form in clouds of mostly molecular hydrogen embedded in the interstellar medium (ISM). These molecular clouds have typical temperatures of 10–20 K and number densities of  $10^3$ – $10^4$  g cm $^{-3}$ . This is far denser than typical phases of the ISM ( $\lesssim 50$  g cm $^{-3}$ ). Turbulence in these clouds results in density perturbations which can become gravitationally unstable and collapse to ultimately form stars.

In this chapter, we will present the basics behind star formation, from the conditions for collapse of cold molecular cores to the formation of radiative pre-main sequence stars.

### *11.1 The Jeans' criterion*

Local overdensities in turbulent molecular clouds in the ISM must reach a minimum size in order to become gravitationally unstable. We can derive this minimum core size by reviewing a classic piece of stability analysis. Let's consider the collapse of a homogeneous, isotropic fluid. Our equations are conservation of mass and momentum, plus Poisson's equation for the gravitational potential:

$$\partial_t \rho + \nabla \cdot (\rho \mathbf{u}) = 0 \quad (11.1)$$

$$\partial_t (\rho \mathbf{u}) + \nabla \cdot (\rho \mathbf{u} \mathbf{u}) = -\rho \nabla \Phi - \nabla P \quad (11.2)$$

$$\nabla^2 \Phi = 4\pi G\rho. \quad (11.3)$$

Right away we run into a snag: if our system is homogenous and isotropic, then  $\nabla \Phi = 0$ , since there is no preferred direction for the vector to point. In this case the left-hand side of Poisson's equation vanishes, which is inconsistent with our having a background density. This, of course, is the central point to cosmology, and if we want to do this calculation correctly we need general relativity and an expanding background universe.

Instead, we shall take an alternate route: following Jeans' lead, we simply assert that there is a background state of uniform density  $\rho_0$ .

This is known as the Jeans Swindle [c.f., Binney and Tremaine, 1987] and, while not quite consistent, still provides insight. Forging ahead, we write the density, velocity, and potential as a background piece (subscript “0”) plus a perturbation (subscript “1”):

$$\begin{aligned}\rho(\mathbf{x}, t) &= \rho_0 + \rho_1 \exp(i\mathbf{k} \cdot \mathbf{x} - i\omega t) \\ \mathbf{u}(\mathbf{x}, t) &= \mathbf{u}_1 \exp(i\mathbf{k} \cdot \mathbf{x} - i\omega t) \\ \Phi(\mathbf{x}, t) &= \Phi_1 \exp(i\mathbf{k} \cdot \mathbf{x} - i\omega t).\end{aligned}$$

Here,  $k$  is the wavenumber of the perturbations and has the usual relation to wavelength,  $k = 2\pi/\lambda$ . We take the perturbations to be adiabatic:  $dP = c_s^2 d\rho$ .

Re-writing Eq. (11.1) in perturbed form gives,

$$\partial_t(\rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}) + \nabla \cdot (\rho_0 \mathbf{u}_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} + \rho_1 \mathbf{u}_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}) \mathbf{x}.$$

We shall drop all terms higher than first order in the perturbation and so the second term in the divergence is zero. Furthermore, we shall pick a direction by dotting  $\mathbf{k}$  into the perturbed equations, yielding

$$-i\omega \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} + ik \rho_0 u_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} = 0,$$

or

$$u_1 = \frac{\omega \rho_1}{k \rho_0}. \quad (11.4)$$

Similarly, for Poisson’s equation (11.3) we have

$$\nabla^2(\Phi_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}) = 4\pi G \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t},$$

where we remember the Swindle and drop the dependence on  $\rho_0$ .

Dotting by  $\mathbf{k}$  and taking a couple spatial derivatives gives

$$-k^2 \Phi_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} = 4\pi G \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t},$$

or

$$\Phi_1 = -\frac{4\pi G}{k^2} \rho_1. \quad (11.5)$$

Finally, for the momentum equation (11.2)

$$\begin{aligned}\partial_t(\rho_0 \mathbf{u}_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} + \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} \mathbf{u}_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}) &= \\ -\nabla(c_s^2 \rho_0 + c_s^2 \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}) - (\rho_0 + \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}) \nabla(\Phi_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t}),\end{aligned}$$

where we have already dropped all terms in the divergence since they are all second order or higher and we have used the adiabatic relation for  $P$ . Dotting by  $\mathbf{k}$  and taking some derivatives, we have

$$\begin{aligned}-i\omega \rho_0 u_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} &= -ik c_s^2 \rho_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} - i\rho_0 k \Phi_1 e^{i\mathbf{k} \cdot \mathbf{x} - i\omega t} \\ \omega \rho_0 u_1 &= kc_s^2 \rho_1 + \rho_0 k \Phi_1.\end{aligned}$$

Using Eqs. (11.4) and (11.5) to eliminate  $u_1$  and  $\Phi_1$  gives us

$$\begin{aligned}\omega^2 \rho_1 &= k^2 c_s^2 \rho_1 - 4\pi G \rho_0 \rho_1 \\ \omega^2 &= k^2 c_s^2 - 4\pi G \rho_0.\end{aligned}\quad (11.6)$$

Equation (11.6) is the *dispersion relation* for the Jeans instability. For sufficiently small  $k$  (long wavelengths)  $\omega^2 < 0$  and the perturbations will grow exponentially. Setting  $\omega = 0$  defines the **Jeans' length**,

$$\lambda_J = c_s \sqrt{\frac{\pi}{G \rho_0}}. \quad (11.7)$$

We recognize the right-hand side as just being  $\sim c_s \tau_{\text{dyn}}$ : equation (11.7) simply says that regions where the sound-crossing time are longer than the dynamical timescale are subject to collapse. In other words, the *pressure* of the system cannot respond quickly enough to reestablish hydrostatic balance against gravity. The mass contained in a box of size  $\lambda_J$  is

$$M_J = c_s^3 \sqrt{\frac{\pi^3}{G^3 \rho}} = \sqrt{\frac{\pi^3 k_B^3}{G^3 m_u^4}} \sqrt{\frac{T^3}{\mu^4 n}} \approx 100 M_\odot \sqrt{\frac{T^3}{\mu^4 n}}, \quad (11.8)$$

where<sup>1</sup>  $c_s^2 = P/\rho$ ,  $P = nk_B T$ , and  $\rho = \mu m_u n$ . The last approximate equality holds for conditions appropriate to the dense cores of molecular clouds—temperatures  $\sim 10$  K, H<sub>2</sub> densities  $\sim 10^3$  cm<sup>-3</sup>.<sup>2</sup> Thus, the typical mass of gravitationally collapsing cores in molecular clouds is about  $100 M_\odot$ . Typical giant molecular clouds in the disk of the Milky Way have masses of  $10^6$  to  $10^7 M_\odot$ .

## 11.2 Collapse of Dense Cores: the Free-fall Phase

Once a cloud exceeds the Jeans' mass, collapse ensues on roughly the free-fall time scale,

$$\tau_{\text{ff}} = \sqrt{\frac{3\pi}{32G\rho}} \approx \frac{5 \times 10^7 \text{ yr}}{\sqrt{\mu n}}.$$

For conditions appropriate to molecular clouds ( $\mu = 2$ ,  $n = 10^3$  g cm<sup>-3</sup>) the free-fall time is about  $10^6$  yr. If the cloud contracts adiabatically, the temperature will increase with the density. For the adiabatic contraction of an ideal gas,  $P \sim \rho^{5/3}$  and  $T \sim P/\rho$ , so  $T \sim \rho^{2/3}$ . Plugging this scaling into the Jeans mass gives  $M_J \sim \rho^{1/2}$ . At some point, the cloud will no longer exceed its Jeans' mass, i.e., it will reach hydrostatic equilibrium. This would be a problem as no stars would ever form, and we wouldn't be here to ask why not. Fortunately, the collapse is *not* adiabatic. The collapsing clouds cool radiatively as they contract.

<sup>1</sup> Molecular clouds in the ISM are approximately isothermal.

<sup>2</sup> Yeah, we kinda ignored the fact that  $\mu = 2$  for H<sub>2</sub> but, hey, we also ignored a factor of  $4\pi/3$  by pretending that it was a collapsing *cube* rather than *sphere*, so...

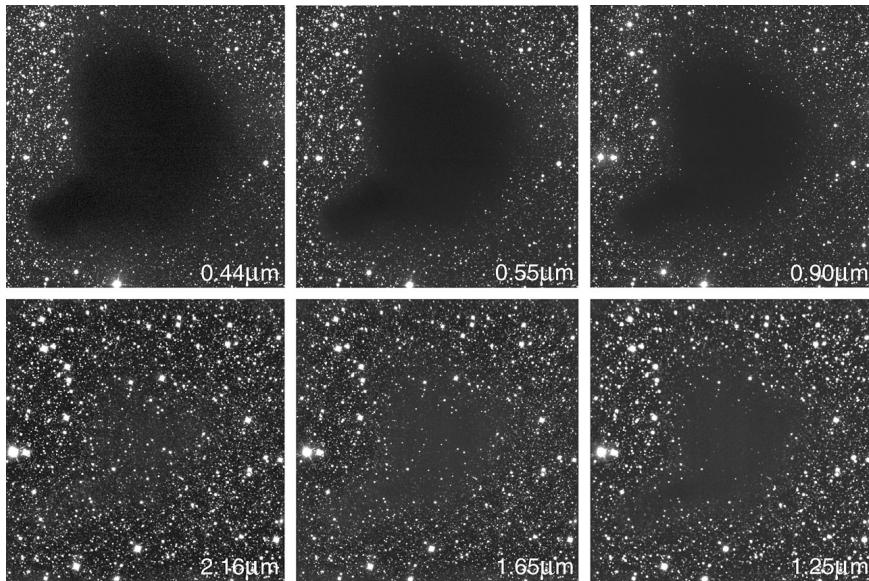


Figure 11.1: Barnard 68, a classic “Bok globule.” This turns out to be a stellar nursery, with the young stars embedded in an opaque (at optical wavelengths) shroud. William Herschel was the first astronomer to discover such “nebulae” as this one, upon which he exclaimed, “My god! There’s a hole in the sky!”

Molecular clouds are optically thick to visible and UV photons. Figure 11.1 shows images of a molecular cloud at various different wavelengths. At optical wavelengths, the cloud is clearly opaque, but at longer wavelengths into the IR, the cloud becomes transparent, revealing young stars embedded within. The cooling of such clouds is thus dominated by emission in the IR where the clouds are transparent. The main drivers of this cooling are emission by molecules and by dust. The cooling is sufficiently efficient that the collapse is essentially *isothermal*, occurring at roughly constant temperature throughout. This overcomes the increasing Jeans’ mass catastrophe mentioned above, allowing the continued collapse of unstable clouds.

FRAGMENTATION of the molecular clouds occurs as they collapse. This happens since, for isothermal collapse, the Jeans’ mass decreases as  $\rho^{-1/2}$  and so smaller overdensities become Jeans unstable and collapse within the overall collapsing cloud. That is, the Jeans *length* decreases faster than the overall size of the cloud. Smaller and smaller fragments collapse faster and faster leading to a hierarchical breakup of the overall cloud and a distribution of so-called *clump* masses.

These clumps will become protostars and, eventually, stars proper. Fragmentation of molecular clouds occurs down to a minimum mass of about  $0.01 M_{\odot}$ .

Collapse clumps eventually become stars, but this process is rather inefficient. Massive stars form first (because more massive clumps have shorter free-fall times, and well, every other time scale is also short!). The winds and ultimate explosions of these stars drive gas

out of star forming regions. Roughly, only about 10% of the original mass of a giant molecular cloud makes it into newborn stars.

**EXERCISE 11.1**— In this exercise you will show that the minimum mass of a star collapsing via the Jeans instability is

$$M_{\min} \approx 10^{-2} M_{\odot}. \quad (11.9)$$

As discussed above, in order for a clump to continue collapsing, it must be able to radiate away the liberated gravitational binding energy, at a minimum. In other words,

$$L_{\text{clump}} > -\frac{dE_{\text{grav}}}{dt}.$$

The cooling will become inefficient was the clump is sufficiently dense to be optically thick to IR photons, at which point it will radiate like a black body (that's a *hint*). Finally, since the clump being considered did in fact collapse, you may assume it is of at least the Jeans' mass (Eq. 11.8). Given all the above, show that the minimum clump mass that will gravitationally collapse is indeed Eq. (11.9).

### 11.3 The End of Free-fall and the Birth of a Protostar

During the free-fall phase, a contracting clump remains transparent to the IR photons that dominate the cooling. Once the clump becomes sufficiently dense, however, it becomes opaque to the molecular cooling photons and the isothermal contraction turns to adiabatic contraction, at least for the inner core of the clump. This drives the temperature up ( $T \sim \rho^{2/3}$ ). Eventually, the high temperatures will dissociate the H<sub>2</sub> molecules, robbing the core of its principle coolant, further exacerbating the increase in temperature. Ultimately, the atomic hydrogen is ionized. The temperature, and hence pressure, will rise enough to bring the clump into hydrostatic equilibrium, at which point we may consider it a *protostar*.

Consider the free-fall collapse of a  $1 M_{\odot}$  clump. To find our initial conditions for this, we need to compute when the Jeans' mass is  $1 M_{\odot}$  (at which point such clumps will fragment out from the larger, collapsing giant molecular cloud). Assuming isothermal collapse up to this point, Eq. (11.8) tells us that  $1 M_{\odot}$  clumps will collapse when the density reaches about  $10^6 \text{ cm}^{-3}$ . The free-fall time for such a clump is only about 35,000 yrs! In short order, then, this  $1 M_{\odot}$  clump will become dense enough to be opaque to the cooling IR photons. This happens at about a density of  $10^{-13} \text{ g cm}^{-3} \sim 3 \times 10^{10} \text{ cm}^{-3}$ . At which point, the clump has a radius  $R_0 \sim 24,000 R_{\odot}$ . From this point on, the collapse will proceed somewhat adiabatically, and the

temperature will increase. Eventually, the temperature will be high enough that the clump will reach hydrostatic equilibrium and the free-fall collapse will stop. The end of this phase is marked by core temperatures sufficiently high to dissociate H<sub>2</sub> molecules and fully ionize H atoms.

We can estimate the conditions at the end of the free-fall phase by considering the energetics of H<sub>2</sub> dissociation and H ionization. Dissociating a H<sub>2</sub> molecule requires an energy  $x_{\text{H}_2} = 4.5 \text{ eV}$  while the ionization potential is  $x_{\text{H}} = 13.6 \text{ eV}$ .<sup>3</sup> So, assuming a hydrogen fraction  $X = 0.7$  and a helium fraction  $Y = 0.3$ , the total energy to dissociate and ionize all the hydrogen in a clump is

$$\begin{aligned} E_{\text{dis,ion}} &= \frac{MX}{m_u}(0.5x_{\text{H}_2} + x_{\text{H}}) \\ &= 1.3 \times 10^{58} \frac{M}{M_\odot} \text{ eV} \approx 2 \times 10^{46} \frac{M}{M_\odot} \text{ erg}. \end{aligned}$$

This energy is provided by released gravitational binding energy due to the collapse, which for an  $n = 2$  polytrope<sup>4</sup> is

$$\Delta E_{\text{grav}} \approx GM^2 \left( \frac{1}{R} - \frac{1}{R_0} \right) \approx \frac{GM^2}{R} \sim 4 \times 10^{48} \frac{(M/M_\odot)^2}{R/R_\odot} \text{ erg},$$

where we have assumed  $R_0 \sim \infty$ .

We can now solve for the final radius  $R$  at which the clump will be hot enough to fully dissociate and ionize the H, but we must account for the fact that not all of the liberated gravitational binding energy will go into heating the gas. A large fraction will be radiated away. If we assume the clump contracts slowly enough from this point on to be in rough virial equilibrium, then this fraction will be 0.5. Thus, we have

$$\begin{aligned} 0.5\Delta E_{\text{grav}} &= E_{\text{dis,ion}} \\ 2 \times 10^{48} \frac{(M/M_\odot)^2}{R/R_\odot} \text{ erg} &= 2 \times 10^{46} \frac{M}{M_\odot} \text{ erg} \\ \Rightarrow \frac{R}{R_\odot} &\sim 100 \frac{M}{M_\odot}. \end{aligned} \quad (11.10)$$

So our  $1 M_\odot$  protostar has quite a long way to go yet before it looks like the Sun!

To estimate the temperature of the protostar we will once again appeal to an  $n = 2$  polytrope. From Eq. (4.18), with the appropriate numerical values from Table 16.1, we have

$$T_c \approx \left( \frac{\mu}{N_A k_B} \right) \left( \frac{GM}{2R} \right) \approx 5 \times 10^6 \text{ K} \frac{M/M_\odot}{R/R_\odot}. \quad (11.11)$$

For our  $1 M_\odot$  clump with  $R \sim 100 R_\odot$ , this gives  $T_c \sim 5 \times 10^4 \text{ K}$ . Not hot enough for fusion of any kind! The mean density of a  $1$

<sup>3</sup> Recall, 1 ev =  $1.602 \times 10^{-12} \text{ erg}$ .

<sup>4</sup> Simply for convenience.

$M_{\odot}$  protostar at the end of the free-fall phase is  $\sim 10^{-6} \text{ g cm}^{-3} \sim 6 \times 10^{17} \text{ g cm}^{-3}$ . A good estimate for the total time spent in the free-fall phase is to assume the free-fall time with the geometric mean of the initial critical density ( $10^3 \text{ g cm}^{-3}$  for the GMC) and the final critical density ( $10^6 \text{ g cm}^{-3}$  for a  $1 M_{\odot}$  clump). This yields

$$\tau_{\text{ff}} \approx 200,000 \text{ yr}$$

#### 11.4 The Hayashi Track

Under the conditions at the end of the free-fall phase ( $\rho \sim 10^{-6} \text{ g cm}^{-3}$  and  $T \sim 5 \times 10^4 \text{ K}$ ), the total opacity will be  $\kappa \sim 10^3 \text{ cm}^2 \text{ g}^{-1}$  (see Figure 7.2). This is opaque enough that the protostar is fully convective. Since the protostar is in rough hydrostatic equilibrium, it will evolve on the Kelvin-Helmholtz timescale,  $\tau_{\text{KH}} \sim E_{\text{grav}}/L$ . Now, let's estimate the luminosity of the protostar, for which will need the effective temperature.

For a fully convective star, we can use the fact that the entropy per unit mass is the same at the photosphere as at the center to relate the surface temperature to the mass and radius of the star. Namely,

$$T_{\text{eff}} = T_c \left( \frac{P_{\text{ph}}}{P_c} \right)^{2/5}, \quad (11.12)$$

with

$$T_c = 0.5 \frac{GM\mu m_u}{k_B R} \quad (11.13)$$

$$P_c = 0.8 \frac{GM^2}{R^4} \quad (11.14)$$

being the central density and pressure of a polytrope of index  $n = 3/2$  (Eqs. 4.18 and 4.17). We can estimate the pressure at the photosphere, defined as the location where the optical depth  $\tau \approx 1$ , by rewriting the definition of optical depth (Eq. 7.17) as a differential equation,

$$\frac{d\tau}{dz} = -\rho\kappa, \quad (11.15)$$

with  $\kappa$  being the total opacity. Using the equation of hydrostatic equilibrium to eliminate  $\rho$

$$\frac{d\tau}{dz} = \frac{\kappa dP}{gdz}. \quad (11.16)$$

This is pretty easy to integrate. Doing so, and setting  $\tau = 1$ , we find an approximate expression for the pressure at the photosphere

$$P_{\text{ph}} = \frac{g}{\kappa}. \quad (11.17)$$

We write a general expression for a complicated total opacity that depends on both density and temperature as

$$\kappa = \kappa_0 \rho^r T^s. \quad (11.18)$$

Substituting this into our expression for  $P_{\text{ph}}$ , and eliminating the density by assuming an ideal gas law, we find a transcendental equation for the pressure at the photosphere

$$\begin{aligned} P_{\text{ph}} &\approx \frac{GM}{R^2} \frac{1}{\kappa_0 (\mu m_u / k_B)^r P_{\text{ph}}^r T_{\text{eff}}^{s-r}} \\ P_{\text{ph}} &\approx [GMR^{-2} \kappa_0^{-1} (\mu m_u / k_B)^{-r} T_{\text{eff}}^{r-s}]^{\frac{1}{r+1}}. \end{aligned} \quad (11.19)$$

Now for some crazy fractions: insert equations (11.13), (11.14), and (11.19) into equation (11.12) and solve for  $T_{\text{eff}}$  to find

$$T_{\text{eff}}^{5+3r+2s} = 0.55^{5(1+r)} \kappa_0^{-2} \left( \frac{G \mu m_u}{k} \right)^{5+3r} M^{3+r} R^{3r-1}. \quad (11.20)$$

What does this say for Thomson scattering ( $\kappa_0 \approx 0.4$ ,  $r = s = 0$ )?

Inserting these values into eq. (11.20) gives

$$T_{\text{eff}} \approx 250 \text{ K} \left( \frac{M}{M_\odot} \right)^{3/5} \left( \frac{R}{R_\odot} \right)^{-1/5}, \quad (11.21)$$

a ridiculous value. Let's try it with H<sup>-</sup> opacity ( $\kappa_0 \approx 2.5 \times 10^{-31}$ ,  $r = 1/2$ ,  $s = 9$ ). In this case,

$$T_{\text{eff}} \approx 2200 \text{ K} \left( \frac{M}{M_\odot} \right)^{1/7} \left( \frac{R}{R_\odot} \right)^{1/49}. \quad (11.22)$$

Note the extremely weak dependence on  $R$ . This is a consequence that  $R$  is very sensitive to the entropy at the photosphere. Writing  $L = 4\pi R^2 \sigma_{\text{SB}} T_{\text{eff}}^4$ , we can solve for  $R$  and insert it into equation (11.22) to get the effective temperature in terms of mass and luminosity,

$$T_{\text{eff}} \approx 2300 \text{ K} \left( \frac{M}{M_\odot} \right)^{7/51} \left( \frac{L}{L_\odot} \right)^{1/102}. \quad (11.23)$$

This is a crude estimate, so don't take these numbers too seriously. What this exercise illustrates, however, is that the effective temperature for fully convective, low-mass stars is essentially independent of luminosity. On an HR diagram, these stars follow a vertical track, known as the *Hayashi track*, as they contract to the main sequence.

The HR diagram in Figure 11.2 shows the pre-main sequence evolution for stars of various masses. The tracks begin in the upper right of the diagram. The Hayashi phase is the roughly vertical portion.

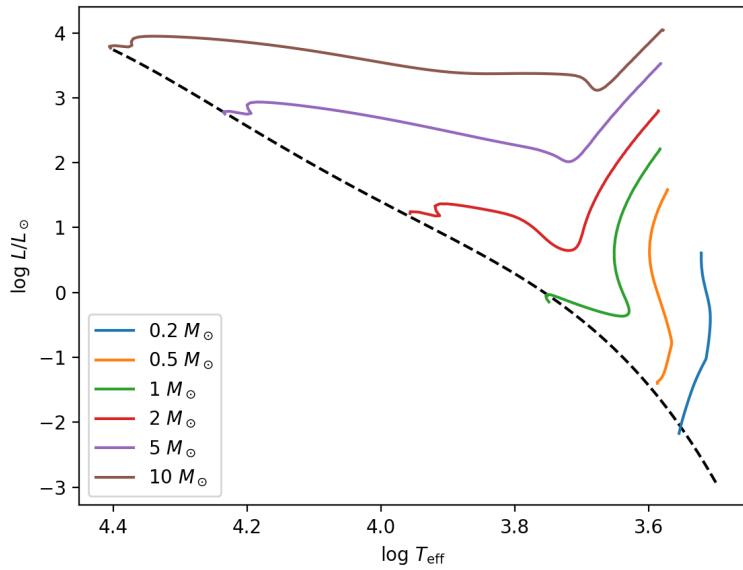


Figure 11.2: Pre-main sequence tracks for stars of several different masses. The main sequence is represented by the dashed line. The Hayashi tracks are the roughly vertical, constant temperature lines on the right while the radiative, protostellar phases are the horizontal, constant luminosity lines.

The time a protostar will spend on the Hayashi track is given by the Kelvin-Helmholtz time. Now that we know the effective temperature and the radius (Eq. 11.10), we can find the luminosity

$$\begin{aligned} L &= 4\pi R^2 \sigma_{\text{SB}} T_{\text{eff}}^4 \\ &= 10^{36} \text{ergs}^{-1} (M/M_{\odot})^{130/49} \approx 250 L_{\odot} (M/M_{\odot})^{130/49}. \end{aligned}$$

This is the luminosity at the *beginning* of the Hayashi phase. The radius will decrease as the star cools, dropping the luminosity drastically. This phase will end when the star is no longer fully convective. That is, when the opacity has dropped sufficiently that the energy transport can be accomplished by radiation. This should be the case when the mean temperature  $\langle T \rangle$  has risen to about  $3 \times 10^6$  K (see Fig. 7.2). Since the protostar is in approximate hydrostatic balance, we can apply the virial theorem to find the radius at which the mean temperature has thus,

$$\begin{aligned} \frac{3}{2} \frac{M}{\mu m_u} k_B \langle T \rangle &= \frac{GM^2}{R} \\ \Rightarrow R &\approx \frac{\mu m_u}{k_B} \frac{2GM}{3\langle T \rangle}. \end{aligned} \quad (11.24)$$

So for  $\langle T \rangle = 3 \times 10^6$  K,

$$R_{\text{PMS}} \sim 2R_{\odot} \frac{M}{M_{\odot}}, \quad (11.25)$$

where we have labeled this as the “pre-main sequence” radius.

The mean luminosity of a protostar on the Hayashi track can be

approximated using the geometric mean radius during the contraction, keeping the effective temperature fixed. So

$$\langle R \rangle \sim \sqrt{100 R_\odot \times 4 R_\odot} \frac{M}{M_\odot} \sim 14 R_\odot \frac{M}{M_\odot}.$$

Rounding up a bit and taking  $T_{\text{eff}} \sim 3000$  K, the mean luminosity is

$$\langle L \rangle \approx 4\pi \langle R \rangle^2 \sigma_{\text{SB}} T_{\text{eff}}^4 \approx 5 \times 10^{34} \text{ erg s}^{-1} \left( \frac{M}{M_\odot} \right)^2 \approx 13L_\odot \left( \frac{M}{M_\odot} \right)^2.$$

Thus the Kelvin-Helmholtz time for a protostar on the Hayashi track is

$$\begin{aligned} \tau_{\text{KH}} &= \frac{-0.5\Delta E_{\text{grav}}}{L} \approx \frac{GM^2}{R\langle L \rangle} \\ &\approx \frac{GM^2}{100 R_\odot (M/M_\odot) \times 13 L_\odot \text{ erg s}^{-1} (M/M_\odot)^2} \\ &\approx 10^6 \text{ yr} \left( \frac{M}{M_\odot} \right)^{-1}. \end{aligned} \quad (11.26)$$

We rounded up a bit on the gravitational energy as the protostar becomes more and more centrally concentrated, but this estimate gives good agreement with the observed values.

### 11.5 Formation of a radiative core

For low-mass stars, the opacity in the core has a Kramers'-like form,  $\kappa \propto \rho T^{-7/2}$ . From our scalings, we see that  $\kappa \propto M^{-5/2} R^{1/2}$ . As a result, as the star contracts, the central opacity decreases. If it decreases enough, then photons can carry the flux and a radiative core will develop. This marks the end of the Hayashi phase.

To find out when a radiative core forms, let's start with our equation for the flux,

$$F(r) = -\frac{1}{3} \frac{c}{\rho \kappa} \frac{daT^4}{dr},$$

and use hydrostatic balance to rewrite this as

$$L(r) = 4\pi r^2 F(r) = \frac{16\pi}{3} \frac{acT^4}{\kappa} \frac{Gm(r)}{P} \frac{d \ln T}{d \ln P}.$$

Now imagine we consider a tiny amount of matter  $\delta m$  about the center of the star. The luminosity coming out of this sphere is  $\delta L(r)$ , and since in a convectively stable atmosphere  $d \ln T / d \ln P < (\Gamma_2 - 1)/\Gamma_2$ , the maximum amount of energy that can be generated in the sphere  $\delta m$  and transported away in the absence of convection is

$$\frac{\delta L(r)}{\delta m} < \frac{16\pi}{3} \frac{Gac}{\kappa} \frac{T_c^4}{P_c} \frac{\Gamma_2 - 1}{\Gamma_2}. \quad (11.27)$$

Flashback! Do you remember doing problem 3.3 of chapter 3? This gave us an expression, eq. (3.60), for  $\partial L/\partial m$ , which we can equate with the LHS of equation (11.27):

$$q - \frac{P_c}{\rho_c(\Gamma_3 - 1)} \frac{D}{Dt} \ln \left( \frac{P_c}{\rho_c^{\Gamma_1}} \right) < \frac{16\pi Gac}{3} \frac{T_c^4}{P_c} \frac{\Gamma_2 - 1}{\Gamma_2}.$$

Let's do the time warp again: in problem 4.2 of chapter 4, you showed that  $P/\rho^{\Gamma_1} \propto R$  for an ideal gas (i.e., the polytropic constant  $K$ ); making this substitution in the derivative gives the condition for the formation of a radiative core,

$$q - \frac{N_A k_B T_c}{\mu(\Gamma_3 - 1)} \frac{D \ln R}{Dt} < \frac{16\pi Gac}{3} \frac{T_c^4}{P_c} \frac{\Gamma_2 - 1}{\Gamma_2}. \quad (11.28)$$

Stars with  $M \gtrsim 0.3M_\odot$  form a radiative core while contracting to the main-sequence; in contrast, lower-mass stars remain fully convective throughout their main-sequence lives.

**EXERCISE 11.2** — We showed that low-mass pre-main-sequence stars (including brown dwarfs) are fully convective and have nearly constant effective temperatures. Use these facts to model their pre-main sequence contraction. Assume  $T_{\text{eff}} = \text{const.}$  so that the luminosity is  $L = 4\pi R^2 \sigma_{\text{SB}} T_{\text{eff}}^4$ .

1. Use the appropriate polytropic relation for the energy of the protostar and assume that the luminosity is entirely powered by contraction, i.e., the star is not yet approaching the main-sequence. Derive an equation for  $R(t)$ . What is the characteristic timescale for a low-mass star to contract? Scale your answer to  $T_{\text{eff}} = 3000\text{ K}$  and  $M = 0.1 M_\odot$  (i.e., get an analytical solution in terms of the variables  $\tilde{T} = [T_{\text{eff}}/3000\text{ K}]$  and  $\tilde{M} = [M/0.1 M_\odot]$ ).
2. Compare your findings with more elaborate calculations. You will find a review in “[Theory of Low-Mass Stars and Substellar Objects](#),” G. Chabrier and I. Baraffe, *Ann. Rev. Astron. Astrophys.* **38:** 337 (2000).

During the radiative phase, the central temperatures are still not high enough to start H fusion. Using a similar estimate as above for the central temperature we find  $T_c \sim 6 \times 10^6\text{ K}$ . Close but not quite (H fusion kicks in around  $10^7\text{ K}$ ). Thus, the stars luminosity is still powered by gravitational contraction. And, since the star is fully radiative, it will evolve with roughly constant luminosity as it contracts. This is shown as the roughly horizontal lines in Figure 11.2.

We can estimate the time spent in the radiative pre-main sequence (PMS) phase by again using the Kelvin-Helmholtz time. The final radius will be the main sequence radius, which we will show in a later chapter scales with mass like

$$R_{\text{MS}} \approx R_\odot \left( \frac{M}{M_\odot} \right)^{0.7}.$$

Assuming a constant luminosity fixed at the *final* luminosity from the Hayashi phase, we have

$$\begin{aligned}\tau_{\text{PMS}} &\approx \frac{-0.5\Delta E_{\text{grav}}}{L_{\text{PMS}}} \approx \frac{GM^2}{4\pi R_{\text{PMS}}^2 \sigma_{\text{SB}} T_{\text{eff,PMS}}^4} \left( \frac{1}{R_{\text{MS}}} - \frac{1}{R_{\text{PMS}}} \right) \\ &\approx \frac{GM^2}{10^{33} \text{ erg s}^{-1} (M/M_{\odot})^2} R_{\odot}^{-1} \left[ \left( \frac{M}{M_{\odot}} \right)^{-0.7} - 0.5 \left( \frac{M}{M_{\odot}} \right)^{-1} \right] \\ &\approx 10^8 \text{ yr} \left[ \left( \frac{M}{M_{\odot}} \right)^{-0.7} - 0.5 \left( \frac{M}{M_{\odot}} \right)^{-1} \right].\end{aligned}\quad (11.29)$$

This scaling is really only very accurate around  $1 M_{\odot}$ . It breaks down, in particular, for massive stars which spend far less time in the PMS phase than this scaling implies ( $10^5$  vs.  $10^7$  years).

OBSERVED PRE-MAIN SEQUENCE STARS are identified as two different types of objects: *T Tauri* stars for low-mass objects ( $M \lesssim 2M_{\odot}$ ) and *Herbig Ae-Be stars* for high-mass objects ( $M \gtrsim 2M_{\odot}$ ). Herbig Ae-Be stars will end up as A and B stars on the main sequence and are characterized by the presence of strong emission lines coming from the accretion disk around the protostar. O stars evolve so quickly through the PMS phase they are rarely observed.

*T Tauri* stars end up as F, G, and K stars on the main sequence. These stars are further separated into different classes based on their age:

- Class o:  $\lesssim 10^4$  yr. Roughly spherical still. Far IR emission.
- Class I:  $\sim 10^5$  yr. Embedded in thick disk with jets. Hayashi track.
- Class II:  $\sim 10^6$  yr. Light from central “star” becomes visible as the disk becomes less opaque. On the constant-luminosity PMS track.
- Class III: Weak emission lines. Emission stronger in the optical. Thin accretion disk. Near the main sequence.

HYDROGEN IGNITION occurs when the central temperature exceeds about  $6 \times 10^6$  K and the star joins the main sequence. Protostars with  $M \lesssim 0.08 M_{\odot}$  will never ignite H because these cores will become degenerate, halting collapse, prior to these temperatures being reached. These objects are called *brown dwarfs* and shine primarily in the IR, powered by the release of gravitational binding energy.

## 11.6 The Initial Mass Function

As originally formulated<sup>5</sup>, the initial mass function (IMF) is the num-

<sup>5</sup> E. E. Salpeter. The Luminosity Function and Stellar Evolution. *ApJ*, 121: 161–+, January 1955

ber of stars, per unit volume, that have formed per logarithmic (base 10) mass interval:

$$\xi(\lg m) \equiv \frac{d(N_*/V)}{d \lg m}. \quad (11.30)$$

Here  $m \equiv M/M_\odot$ . The IMF is derived from an observed *luminosity function*—the amount of starlight in a given waveband emitted per unit mass—and stellar models. As might be expected, this function is not well-constrained, but it is roughly a power law for  $m > 1.0$ , and at lower masses it flattens out. One such formulation<sup>6</sup> for the solar neighborhood is

$$\xi(\lg m) = \begin{cases} A_0 \exp\left[-\frac{(\lg m - \lg m_c)^2}{2\sigma^2}\right], & m < 1.0 \\ A_1 m^b & m > 1.0 \end{cases}. \quad (11.31)$$

Here  $A_0 = 0.158$ ,  $A_1 = 4.43 \times 10^{-2}$ ,  $m_c = 0.079$ ,  $\sigma = 0.69$ , and  $b = -1.3 \pm 0.3$ . The coefficients  $A_0$ ,  $A_1$  are in  $\lg M_\odot^{-1} \text{ pc}^{-3}$ .

<sup>6</sup> Gilles Chabrier. Galactic stellar and substellar initial mass function. *Publ. Astron. Soc. Pac.*, 115:763–796, 2003

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**EXERCISE 11.3**— For the IMF given in eq. (11.31), what fraction of the stars formed with  $m > 1.0$  will end their lives as core-collapse supernovae, if the mass threshold for forming a core-collapse SNe is  $8.0 M_\odot$ ? What is the fraction if the mass threshold is  $12.0 M_\odot$ ?

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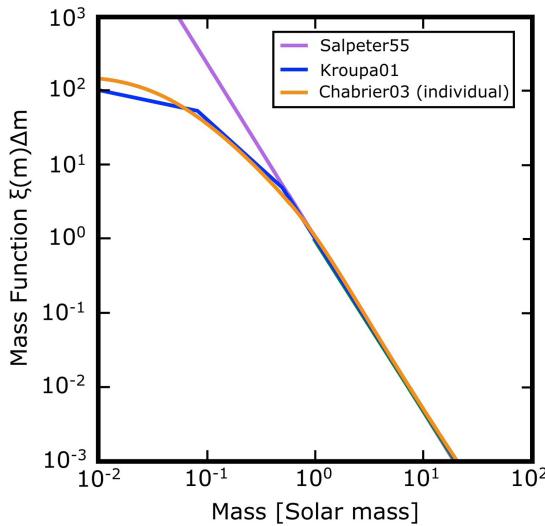


Figure 11.3: title

The main-sequence lifetime is a rapidly decreasing function of mass: for  $m > 1.0$ , it goes roughly like  $\tau_{\text{MS}} \approx 10.0 \text{ Gyr } m^{-2.5}$ . For stars with lifetimes comparable to or longer than the age of the galactic disc  $\tau_G$ , all stars that were ever formed are still on the main sequence, so that the IMF is identical to the *present day mass function* (PDMF)  $\phi$ . For more massive stars, however, we only see those that were formed a time  $\tau_{\text{MS}}(m)$  ago.

Let's define a birthrate  $B(t)$  as the number of stars per unit volume formed per interval of time. If we make the *ansatz* that the IMF doesn't depend on time, then we can define a creation function,

$$C(\lg m, t) \equiv \xi(\lg m) \frac{B(t)}{\int_0^{\tau_G} B(t) dt}. \quad (11.32)$$

Here the birthrate is normalized to the total number of stars formed over the age of the galactic disk. The present-day mass function is then

$$\phi(\lg m) = \int_{\max(0, \tau_G - \tau_{\text{MS}})}^{\tau_G} C(\lg m, t) dt.$$

For a constant birthrate over the age of the disk, the integral is trivial and

$$\phi(\lg m) = \begin{cases} \xi(\lg m) \frac{\tau_{\text{MS}}(m)}{\tau_G} & \tau_{\text{MS}} < \tau_G \\ \xi(\lg m) & \tau_{\text{MS}} > \tau_G \end{cases}$$

As the galaxy ages, the stellar population becomes increasingly dominated by long-lived, low-mass stars. Empirically, the Milky Way birthrate has in fact been more or less constant (deviations less than a factor of 2) over the life of the galactic disk. The timescale for converting the present supply of gas into stars is  $\sim (1\text{--}5)\text{Gyr}$ .

For an IMF, we can define, for each generation of stars, a *lock-up fraction*, which is the amount of gas that is not eventually returned to the interstellar medium. Clearly this will include all stars with  $\tau_{\text{MS}}(m) > \tau_G$ , as well as the remnant mass,  $m_{\text{rem}}(m)$ , of the remaining stars. For stars with mass  $< 8 M_\odot$ , the mass of the white dwarf as a function of the progenitor's mass is fairly well known; more massive stars leave behind either a neutron star, for which observed masses (in binaries!) are  $1\text{--}2 M_\odot$ ; for black holes the remnant mass is more uncertain.

**EXERCISE 11.4—** Suppose the IMF is simply a power-law,  $\xi(\lg m) \propto m^{-1.3}$ , for  $0.1 < m < 120$ . On average, how many stars are formed out of one solar mass of gas?

### 11.7 Application: The delay time of type Ia supernovae

Type Ia supernovae are observed in elliptical galaxies, which typically have an old stellar population and no ongoing star formation. There must be a substantial delay, then, between the time the progenitor was born and the supernovae. To see how this works, let's define a *delay time distribution*  $\mathcal{D}(\tau)$  and a *realization probability*  $A_{\text{Ia}}(t)$ . These are defined as follows: if  $N_*(t)$  is the total number of stars formed at time  $t$ , then define  $N_*(t)A_{\text{Ia}}(t)$  as the total number of SNe Ia that

will ever result from this generation of stars. The SNe Ia rate at the current time,  $t = \tau_G$ , is then

$$\Gamma_{\text{Ia}}(t = \tau_G) = \int_{\tau_{\min}}^{\min(\tau_G, \tau_{\max})} \int_{\lg m_{\min}}^{\lg m_{\max}} C(\lg m, t - \tau) A_{\text{Ia}}(t - \tau) \mathcal{D}(\tau) d\lg m d\tau. \quad (11.33)$$

In this definition, the delay time distribution is normalized to unity.

As an example, suppose all the stars are born in a burst of star formation at  $t = 0$ , so that  $B(t) = \delta(t)$ , then the SNe Ia rate at late times is

$$\Gamma_{\text{Ia}}(t = \tau_G) = \int_{\tau_{\min}}^{\min(\tau_G, \tau_{\max})} N_{\star} \delta(t - \tau) A_{\text{Ia}}(t - \tau) \mathcal{D}(\tau) d\tau = N_{\star} A_{\text{Ia}} \mathcal{D}(t).$$

Notice that if  $\mathcal{D}(t)$  is a very broad function of time, then the SNe Ia rate is proportional, for this case of a rapid burst of star formation at early times, to the total number of stars in the galaxy.

**EXERCISE 11.5**— Consider the SNe Ia rate, eq. (11.33), following a burst of star formation,  $B(t) = \delta(t)$ , but now suppose that the delay time for each mass is just the main-sequence lifetime, and that  $A_{\text{Ia}}$  is independent of mass. That is, for stars with mass  $m = M/M_{\odot} < 8$ , we assume that some fraction  $A_{\text{Ia}}$  will become SNe Ia, and that the time for a particular mass star to evolve to explosion is just its main-sequence lifetime  $\tau(m) = \tau_{\text{MS}}(m)$ . Show that, for  $\xi(\lg m) \propto m^{-1.3}$  and  $\tau_{\text{MS}} = 10.0 \text{ Gyr } m^{-2.5}$ , the Ia rate is  $\Gamma_{\text{Ia}} \propto t^{-0.5}$ , for  $t > \tau_{\text{MS}}(m = 8)$ .



## 12

# *Hydrogen Burning and the Main Sequence*

### *12.1 Hydrogen burning via pp reactions: the lower main sequence*

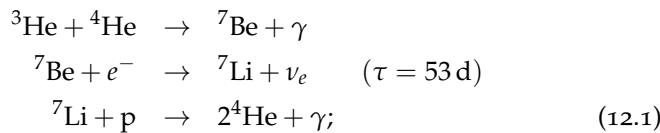
In a contracting pre-main sequence star, the reaction  ${}^2\text{H}(p, \gamma){}^3\text{He}$  proceeds rapidly owing to the small Coulomb barrier; in fact, this reaction can occur in objects as small as  $\approx 12 M_{\text{Jupiter}}$ . The small primordial abundance of deuterium, however, prevents this reaction from doing anything more than slowing contraction slightly. The reaction  $p + p$  is much slower, because there is no bound nucleus  ${}^2\text{He}$ ; the only possible way to form a nucleus is to have a weak interaction as well, giving the reaction  $p(p, e^+ \nu_e){}^2\text{H}$ .

The weak cross section goes roughly as  $\sigma_{\text{weak}} \sim 10^{-20} \text{ b } (E/\text{keV})$ , so that

$$\frac{\sigma_{\text{weak}}}{\sigma_{\text{nuc}}} \sim 10^{-23} \left( \frac{E}{\text{keV}} \right).$$

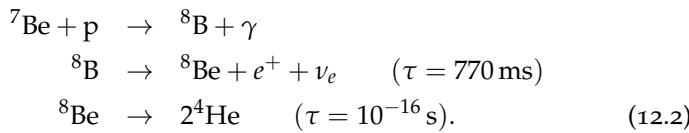
The S-factor for the  $p + p$  reaction is very small, and as a result the characteristic temperature for this reaction to occur is  $\approx 1.5 \times 10^7 \text{ K}$ ; at this temperature, the lifetime of a proton to forming deuterium via capture of another proton is about 6 Gyr. Once a deuterium nucleus is formed, it is immediately destroyed via  ${}^2\text{H}(p, \gamma){}^3\text{He}$ . The nucleus  ${}^4\text{Li}$  is unbound with a lifetime of  $10^{-22} \text{ s}$ ; the nucleus  ${}^6\text{Be}$  is likewise unbound ( $\tau \sim 5 \times 10^{-21} \text{ s}$ ). As a result, the next reaction that can occur is  ${}^3\text{He}({}^3\text{He}, 2p){}^4\text{He}$ . Despite having a much greater Gamow energy than  $p + p$  (see Table 10.2), this reaction still is much faster than  $p + p$  owing to the small weak cross-section.

In addition to capturing another  ${}^3\text{He}$ , it is also possible that



furthermore, at slightly higher temperatures  ${}^7\text{Be}$  can capture a proton

instead of an electron, giving the third branch



The end result of these chains is the conversion of hydrogen to helium, although the amount of energy carried away by neutrinos differs from one chain to the next.

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**EXERCISE 12.1** — Compute the mass of H, in units of solar masses, that must be converted into  ${}^4\text{He}$  in order to supply the solar luminosity over  $10^{10}$  yr.

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## 12.2 Hydrogen burning via the CNO cycle: the upper main sequence

As we saw in the previous section, the smallness of the  $\text{p} + \text{p}$  cross-section means that captures onto heavier nuclei can be competitive at stellar temperatures. Let's get a rough estimate of how charged a nucleus can be before the Coulomb barrier makes the reaction slower than  $\text{p} + \text{p}$ . Assuming  $A = 2Z$ , and taking the  $S$ -factor for  $\text{p} + \text{p}$  to be  $10^{-22}$  times smaller than that for  $\text{p} + {}^A\text{Z}$  gives us the rough equation

$$10^{-22} \exp\left(-\frac{33.81}{T_6^{1/3}}\right) \approx \exp\left(-\frac{41.47Z^{2/3}}{T_6^{1/3}}\right),$$

where the factors in the exponentials come from the peak energy for the reaction (see eq. [10.21]), and  $T_6 \equiv (T/10^6 \text{ K})$ . Solving for  $Z$ , we see that at  $T_6 = 10$ , proton captures onto  ${}^{12}\text{C}$  have a comparable cross-section to  $\text{p} + \text{p}$ ; at  $T_6 = 20$ , proton captures onto  ${}^{16}\text{O}$  have a comparable cross-section.

Thus at temperatures slightly greater than that in the solar center, the following catalytic cycle becomes possible.

reaction	$\log[(\tau/\text{yr}) \times (\rho X_H / 100 \text{ g cm}^{-3})]$
${}^{12}\text{C}(\underline{\text{p}}, \gamma) {}^{13}\text{N}$	3.82
${}^{13}\text{N}(e^+, \nu_e) {}^{13}\text{C}$	$\tau = 870 \text{ s}$
${}^{13}\text{C}(\underline{\text{p}}, \gamma) {}^{14}\text{N}$	3.21
${}^{14}\text{N}(\underline{\text{p}}, \gamma) {}^{15}\text{O}$	5.89
${}^{15}\text{O}(e^+, \nu_e) {}^{15}\text{N}$	$\tau = 178 \text{ s}$
${}^{15}\text{N}(\underline{\text{p}}, {}^4\text{He}) {}^{12}\text{C}$	1.50

As indicated by the underlined symbols, this cycle takes in 4 protons and releases 1 helium nucleus. The reaction timescales are evaluated at a temperature of 20 MK. The reaction  ${}^{14}\text{N}(\text{p}, \gamma) {}^{15}\text{O}$  is by far

the slowest step in the cycle; as a result, all of the CNO elements are quickly converted into  $^{14}\text{N}$  in the stellar core, and this reaction controls the rate of heating. At  $T = 2 \times 10^7 \text{ K}$ ,  $d \ln \varepsilon_{\text{CNO}} / d \ln T = 18$ ; in contrast the  $\text{p} + \text{p}$  reaction has a temperature exponent of only 4.5.

The strong temperature sensitivity of the CNO cycle has a profound effect on the properties of the upper main sequence. Dividing equation (3.55) by equation (3.54), we have

$$\frac{dT}{dP} = \frac{3}{16\pi Gm} \frac{\kappa}{acT^3} L_r. \quad (12.3)$$

For stars with masses  $\gtrsim M_\odot$ , the structure roughly follows a polytrope of index  $n = 3$ . We can insert the relations  $T \sim M/R$  and  $P \sim M^2/R^4$  into equation (12.3) and scale to the solar luminosity to obtain

$$L \approx L_\odot (M/M_\odot)^3. \quad (12.4)$$

On the other hand, we can integrate our equation for the heating rate per unit mass,  $\varepsilon \approx \varepsilon_0 \rho T^n$ , over the star; inserting the scalings for  $T$  and  $\rho$  and normalizing to solar values, we obtain  $L \approx L_\odot (M/M_\odot)^{2+n} (R/R_\odot)^{-3-n}$ . Equating this with  $L$  from eq. (12.4), we find that on the upper main sequence,

$$\frac{R}{R_\odot} \approx \left( \frac{M}{M_\odot} \right)^{(n-1)/(n+3)}. \quad (12.5)$$

For  $n = 18$ , this gives  $R \sim M^{0.81}$ . Since  $L = L_\odot (M/M_\odot)^3 = (R/R_\odot)^2 (T_{\text{eff}}/T_{\text{eff},\odot})^4$ , we can obtain a relation between  $T_{\text{eff}}$  and  $L$  on the upper main sequence,

$$\left( \frac{T_{\text{eff}}}{T_{\text{eff},\odot}} \right) = \left( \frac{L}{L_\odot} \right)^{0.12}. \quad (12.6)$$

The fact that  $T_{\text{eff}}$  is so insensitive to  $L$  is a consequence that the radius increases with mass, which follows from the central temperature being roughly constant. The strong temperature sensitivity of the CNO cycle ensures that the central temperature varies only slightly over a large range of luminosity.

**EXERCISE 12.2—** Suppose that there were no CNO cycle, and hydrogen could only be consumed via the PP chains. Estimate the effective temperature-luminosity relation for the upper main-sequence in this case. Would it be observationally distinguishable from the CNO-dominated upper MS?

A second effect on the stellar structure is that the luminosity is generated in a very small region concentrated about the center. Inserting  $P_{\text{rad}} = (a/3)T^4$  and  $L_{\text{Edd}} = 4\pi GMc/\kappa$  into equation (12.3)

and solving for  $\nabla = d \ln T / d \ln P$ , we obtain

$$\nabla = \frac{1}{4} \frac{P}{P_{\text{rad}}} \frac{L}{L_{\text{Edd}}} \left( \frac{L_r}{L} \right) \left( \frac{M}{M(r)} \right). \quad (12.7)$$

For the Eddington standard model,  $P/P_{\text{rad}} \approx 2600(M/M_{\odot})^{-2}$ , and  $L/L_{\text{Edd}} \approx 2.7 \times 10^{-5}(M/M_{\odot})^2$ . Inserting these factors and using the criteria for convective stability,  $\nabla < (\partial \ln T / \partial \ln P)_S = 2/5$ , we see that if

$$\frac{L_r}{L} > 23 \frac{M(r)}{M}$$

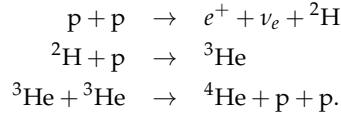
we have convective instability. Thus, if the luminosity is produced in the innermost 4% (by mass) of the star, the core will be convective. The strong temperature sensitivity of the CNO reactions ensure that this is the case, and so the cores of upper main sequence stars have convective zones.

This convective zone changes the structure of the star, so that  $L \sim M^{3.5}$  rather than the  $M^3$  scaling used above. It also means the star can burn more of the hydrogen in its interior. The hotter  $T_{\text{eff}}$  means, however, that the  $H^-$  opacity is not important and the surface layers of upper main sequence stars are radiative. Table 12.1 gives a summary of the properties of main sequence stars.

characteristic	lower ( $M \lesssim M_{\odot}$ )	upper ( $M \gtrsim M_{\odot}$ )
hydrogen burning	pp	CNO
opacity	Kramers	Thomson
core	radiative, $\approx 0.1M$	convective, $\approx 0.2M$
envelope	convective	radiative

Table 12.1: Characteristics of main-sequence stars

**EXERCISE 12.3—** In this exercise, we are going to examine how the core of a solar-mass star approaches a steady conversion of 4 hydrogen nuclei to 1 helium nucleus via the PPI chain



Denote the abundances of protons, deuterium,  ${}^3\text{He}$ , and  ${}^4\text{He}$  by  $Y_p$ ,  $Y_d$ ,  $Y_3$ , and  $Y_4$ , respectively. Furthermore, evaluate the rates at a fiducial central temperature and density (obtained with the MESA stellar evolution code)  $T_{c,\odot} = 1.35 \times 10^7 \text{ K}$ ,  $\rho_{c,\odot} = 83.2 \text{ g cm}^{-3}$ :

$$\begin{aligned} \lambda_{pp} \equiv \rho N_A \langle \sigma v \rangle (p + p \rightarrow {}^2\text{H}) &= 4.40 \times 10^{-18} \text{ s}^{-1} \\ \lambda_{pd} \equiv \rho N_A \langle \sigma v \rangle (p + {}^2\text{H} \rightarrow {}^3\text{He}) &= 2.58 \times 10^{-2} \text{ s}^{-1} \\ \lambda_{33} \equiv \rho N_A \langle \sigma v \rangle ({}^3\text{He} + {}^3\text{He} \rightarrow {}^4\text{He} + p + p) &= 3.40 \times 10^{-9} \text{ s}^{-1} \end{aligned}$$

Take the screening factors to be unity.

1. First, let's consider the build-up of deuterium. Start from equation (10.45):

$$\frac{d}{dt} Y_d = \frac{1}{2} Y_p^2 \lambda_{pp} - Y_d Y_p \lambda_{pd}. \quad (12.8)$$

Assume that over the timescale to establish the PP chain, the abundance of hydrogen  $Y_p$  is constant, and that all the  $\lambda$  are constant as well. Under these assumptions, solve for  $Y_d(t)$  and show that it approaches a constant value

$$\left. \frac{Y_d}{Y_p} \right|_{\text{equil.}} = \frac{1}{2} \frac{\lambda_{pp}}{\lambda_{pd}}.$$

What is this abundance? What is the timescale to reach this equilibrium?

2. Now consider the evolution of  ${}^3\text{He}$  via production by  $d + p$  and destruction via  ${}^3\text{He} + {}^3\text{He}$ :

$$\frac{d}{dt} Y_3 = Y_p Y_d \lambda_{pd} - Y_3^2 \lambda_{33}. \quad (12.9)$$

Assume that  $Y_p$  is constant and use the equilibrium value of  $Y_d$  to solve this equation for  $Y_3(t)$ . Show that  $Y_3$  approaches a constant value

$$\left. \frac{Y_3}{Y_p} \right|_{\text{equil.}} = \left( \frac{1}{2} \frac{\lambda_{pp}}{\lambda_{33}} \right)^{1/2}.$$

What is this value? What is the timescale for the abundance of  ${}^3\text{He}$  to reach 99% of this equilibrium value? Is the assumption that deuterium is at its equilibrium abundance a valid one?

3. Using the equilibrium value of  ${}^3\text{He}$ , show that the rate of helium production via the  ${}^3\text{He} + {}^3\text{He}$  is

$$\frac{d}{dt} Y_4 = \frac{1}{4} Y_p^2 \lambda_{pp}.$$



# 13

## *Post-Main Sequence Evolution: Low-mass Stars*

Now we will discuss the sequence of events following core H burning and the Main Sequence phase of stellar evolution. The post-main sequence phase is marked by dramatic changes in the structure and appearance of stars, relative to the somewhat quiescent main sequence phase during which stars change slowly in time. This rapid evolution is brought on by the competition between the inward pull of gravity and the outward resistance to this pull generated by nuclear burning. The nuclear burning in the post-main sequence phase is far more complicated, structurally, than that during the core H-burning phase, marked by episodes of shell burning, core burning, core and shell burning, double shell burning, etc. This complicated dance of varied burning stages has a dramatic impact on stars, leading to various instabilities and rapid movement through the H-R diagram.

First, let's consider what happens immediately following the cessation of H burning in the cores of stars.

### *13.1 Helium core contraction and the Schönberg-Chandrasekhar Limit*

Once the core of a main sequence star exhausts all the H available in region about the threshold temperature for the pp chain and CNO cycles, nuclear energy generation throughout the core will cease. This implies that the emergent luminosity as a function of radius through the core is zero,  $L_r \approx 0$  for  $r < R_c$ , where  $R_c$  is the core radius. According to the heat equation (3.61), if  $L_r \approx 0$  then the temperature gradient  $\partial T / \partial r$  must also be zero throughout the core. In other words, the cores is isothermal.

Without a temperature gradient to help, in order for the core to maintain hydrostatic balance and strong density gradient must develop to balance the pressure gradient,  $dP/dr \sim d\rho/dr$ . Using arguments similar to that used to arrive at the virial theorem, we can then

show that there is a maximum mass for isothermal cores to maintain hydrostatic equilibrium, i.e., to resist contraction.

Beginning with the equation of hydrostatic equilibrium,

$$\frac{dP}{dr} = -\frac{Gm}{r^2}\rho,$$

let's integrate this equation over the core, first multiplying by  $4\pi r^3$ ,

$$\int_0^{R_c} 4\pi r^3 dP = - \int_0^{R_c} \frac{Gm}{r} \rho 4\pi r^2 dr = E_{\text{grav},c} = -\alpha \frac{GM_c^2}{R_c}.$$

In the final equality, we are making an analogy to the polytropic relations and  $\alpha > 1$ . We can break up the first integral as

$$\begin{aligned} \int_0^{R_c} 4\pi r^3 dP &= 4\pi r^3 P \Big|_0^{R_c} - 3 \int_0^{R_c} P 4\pi r^2 dr \\ &= 4\pi R_c^3 P(R_c) - \frac{3k_B T_c}{\mu m_u} \int_0^{R_c} \rho 4\pi r^2 dr \\ &= 4\pi R_c^3 P(R_c) - \frac{3k_B T_c}{\mu_c m_u} M_c, \end{aligned}$$

where  $M_c$  is the total mass of the core and  $\mu_c$  is its mean molecular weight. Further, we have used the condition of isothermality to pull the constant core temperature  $T_c$  out of the integral. Equating this with the core gravitational potential energy from above,

$$4\pi R_c^3 P(R_c) - \frac{3k_B T_c}{\mu_c m_u} M_c = -\alpha \frac{GM_c^2}{R_c}. \quad (13.1)$$

Solving for the pressure at the edge of the isothermal core,

$$P(R_c) = \frac{3}{4\pi} \frac{k_B T_c}{\mu_c m_u} \frac{M_c}{R_c^3} - \frac{\alpha}{4\pi} \frac{GM_c^2}{R_c}. \quad (13.2)$$

In terms of the inverse core radius, this equation has a maximum at

$$P_{c,\max}(R_c) = C \frac{T_c^4}{\mu_c^4 M_c^2}, \quad (13.3)$$

where  $C$  is a constant. What this implies, then, is that the pressure due to the weight of the envelope resting atop the core must be less than this for the core to remain isothermal and hydrostatic.

If the core makes up only a small fraction of the total mass of the star, we can use our virial scalings to estimate the pressure of the envelope at  $R_c \sim 0$ ,

$$P_{\text{env}}(R_c) \approx \frac{GM^2}{R^4}$$

where  $M$  and  $R$  are now the total mass and radius of the star. At the point where the envelope joins the core,  $R_c$ , the temperature must be

Q: Give a simple physical explanation for why there is a maximum for  $P(R_c)$ .

the same as that of the isothermal core, and we can again estimate this from the virial scalings,

$$T_c \sim \frac{\mu_{\text{env}} m_u}{k_B} \frac{GM}{R}.$$

Using this to eliminate the total radius from expression for the envelope pressure, we have

$$P_{\text{env}}(R_c) \sim \frac{T_c^4}{\mu_{\text{env}}^4 M^2}. \quad (13.4)$$

Thus, in order to maintain  $P_{\text{env}} < P_{c,\text{max}}$  we find

$$\frac{M_c}{M} \lesssim C \left( \frac{\mu_{\text{env}}}{\mu_c} \right)^2, \quad C = 0.37. \quad (13.5)$$

This is the *Schönberg-Chandrasekhar Limit* for isothermal cores. Numerically, using  $\mu_c = 4/3$  and  $\mu_{\text{env}} \approx 0.6$ , this give a maximum core mass to total mass ratio of about 0.1.

If the He core exceeds this mass limit, it cannot remain hydrostatic and isotermal and must contract on roughly KH time:

$$\tau_{\text{KH}} \approx \frac{GM_c^2}{R_c(L - L_{\text{shell}})}$$

where  $L_{\text{shell}}$  is the luminosity due to nuclear reactions in the shell surrounding the He core.

High mass stars have smaller He core at TAMS because the  $\mu$  gradient that develops shrinks the convective core.

Unless He core is extremely small, core will contract. (Chandrasekhar-Schoenberg Limit,  $M = 0.1 \text{ Msun}$ ).

The start of H-shell fusion: Low and intermediate mass stars

Mirror principle - virial theorem Consequence of the star trying to remain in thermal and radiative equilibrium whilst the core contracts. The thermonuclear reactions in the shell still provide a thermostat - but they can't halt the contraction of the core. Except by providing sufficient luminosity to prevent the core from cooling (i.e., keeping the core isothermal).

Post-MS evolution divided at  $2 M_{\odot}$ .

As we showed in Section ??, stars more massive than  $\sim 1.2 M_{\odot}$  burn via the CNO cycle and have convective cores. For these stars, the region of burning is very centrally concentrated, by virtue of the strong temperature dependence of the CNO reactions ( $\sim T^{18}$ ). The convection, however, continually mixes down fresh fuel to burn, efficiently depleting a much more extended region in the core of an upper-main sequence star. Thus, these stars have helium cores at TAMS roughly twice as massive as lower-main sequence stars (0.2

$M_{\odot}$  vs.  $0.1 M_{\odot}$ ). This distinction has a crucial impact on the post-main sequence evolution of these stars. The edge of the H-depleted He cores, and the unburnt H sitting atop it, is at lower temperatures relative to the threshold temperature for H-burning ( $\sim 10^7$  K) than in lower-mass stars. Since there is still a temperature gradient throughout the star, there will be an outward flux of radiative energy necessitating that the core contract and, somewhat paradoxically, heat up. All the while, the star is still subject to the constraints of the virial theorem,

$$E_{\text{tot}} = E_{\text{kin}} + E_{\text{grav}} = \frac{1}{2}E_{\text{grav}}.$$

Thus, if the core heats up, increasing  $E_{\text{kin}}$ , the gravitational potential energy much *decrease* to compensate. Since the mass is fixed, the star as a whole contracts, increasing  $T_{\text{eff}}$  resulting in a blueward loop in the H-R diagram. The contraction of the He core ends when the H shell atop it reaches temperatures sufficient for fusion.

#### Movie

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**EXERCISE 13.1**— On the Red Giant Branch, the Sun will reach enormous luminosities, peaking at about 1000 times the current solar luminosity. Assuming an average RGB phase luminosity of  $10^2 L_{\odot}$ , how long would the RGB phase last if the luminosity were powered entirely by the gravitational contraction of the inner He core (i.e., no H shell burning)? You may take the He core mass to be  $0.1 M_{\odot}$  and its radius to be  $10^{10}$  cm. The actual time the Sun will spend on the RGB is about 0.5 Gyr. How does your answer compare?

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## 13.2 The Triple-Alpha Reaction

The consumption of  ${}^4\text{He}$  is hindered by the lack of stable nuclei with mass numbers  $A = 5$  and  $A = 8$ . The nucleus  ${}^8\text{Be}$  is, however, long-lived by nuclear standards: its decay width is  $\Gamma = 68$  eV, implying a decay timescale  $\hbar/\Gamma = 9.7 \times 10^{-17}$  s. (For comparison, the lifetime of  ${}^5\text{Li}$  is  $\sim 10^{-22}$  s.) Thus if the reaction  ${}^4\text{He} + {}^4\text{He} \rightarrow {}^8\text{Be}$  can proceed quickly enough, a small amount of  ${}^8\text{Be}$  can accumulate allowing the reaction  ${}^8\text{Be} + {}^4\text{He} \rightarrow {}^{12}\text{C}$  to proceed.

The reaction  $2 {}^4\text{He} \rightarrow {}^8\text{Be}$  is endothermic, with  $Q = -92$  keV. As a result, the peak energy for Coulomb barrier transmission (see the discussion following eq. [10.18]) must reach

$$E_{\text{pk}} = \frac{E_{\text{G}}^{1/3}(kT)^{2/3}}{4^{1/3}} = -Q.$$

Substituting  $E_{\text{G}} = 979 \text{ keV} A(Z_1 Z_2)^2$  and solving for  $T$  gives  $T = 1.2 \times 10^8 \text{ K}$  as the temperature required to build up any substantial

amount of  ${}^8\text{Be}$ . To reach such a temperature requires a helium core mass of  $\gtrsim 0.45 M_{\odot}$ .

Once a sufficient temperature is reached, the reaction  $2 {}^4\text{He} \rightarrow {}^8\text{Be}$  comes into equilibrium with the decay,  ${}^8\text{Be} \rightarrow 2 {}^4\text{He}$ . We can use a Saha-like equation (cf. eqns. [6.35] and [6.36]) to get the abundance of  ${}^8\text{Be}$ . Writing

$$2\mu_4 = \mu_8 - Q$$

and substituting the expression for  $\mu$ , eq. (6.7) gives

$$\frac{n_8}{n_4} = \left( \frac{2\pi\hbar^2}{m_u kT} \right)^{3/2} \left( \frac{8}{4^2} \right)^{3/2} \frac{X_4 \rho}{4m_u} \exp \left( -\frac{92 \text{ keV}}{kT} \right). \quad (13.6)$$

Here we use  $n_8$  and  $n_4$  to mean the number densities of, respectively,  ${}^8\text{Be}$  and  ${}^4\text{He}$ ; also  $X_4$  denotes the mass fraction of helium. Scaling eq. (13.6) to  $\rho = \rho_5 10^5 \text{ g cm}^{-3}$  and  $T = T_8 10^8 \text{ K}$ , we obtain

$$\frac{n_8}{n_4} = 2.8 \times 10^{-5} T_8^{-3/2} X_4 \rho_5 \exp \left( -\frac{10.68}{T_8} \right). \quad (13.7)$$

At  $X_4 \rho_5 = 1$  and  $T_8 = 1$ ,  $n_8/n_4 = 6.5 \times 10^{-10}$ .

If the reaction  ${}^8\text{Be} + {}^4\text{He} \rightarrow {}^{12}\text{C}$  were non-resonant, the reaction rate for this  $n_8$  would be far too slow to account for the amount of  ${}^{12}\text{C}$  synthesized in stars. Hoyle proposed, therefore, that there should be an excited state of the  ${}^{12}\text{C}$  nucleus into which the reaction would proceed. Both  ${}^4\text{He}$  and  ${}^8\text{Be}$  have spin and parity  $J^\pi = 0^+$ ; hence for  $s$ -wave capture (angular momentum  $\ell = 0$ ), the state in  ${}^{12}\text{C}$  should also have  $J^\pi = 0^+$ . What energy should the level have? The  $Q$ -value for  ${}^8\text{Be} + {}^4\text{He} \rightarrow {}^{12}\text{C}$  is 7.367 MeV; the Gamow energy for this reaction is  $E_G = 1.67 \times 10^5 \text{ keV}$  and hence the peak energy is  $E_{\text{pk}} = 146 \text{ keV}$ , with a width  $\Delta = 4(E_{\text{pk}} kT/3)^{1/2} = 82 \text{ keV}$ . The proposed level should therefore have an energy within  $2\Delta$  of 7.513 MeV.

Such a level was indeed detected by Fowler, with  $J^\pi = 0^+$  and  $E = 7.654 \text{ MeV}$ . Radiative decay from this level is hampered: the ground state also has  $J^\pi = 0^+$ , so the decay is forbidden; and the decay to the  $J^\pi = 2^+$  state at 4.44 MeV has a decay width of only  $\Gamma_{\text{rad}} = 3.67 \text{ meV}$ . This level's primary decay is indeed back to  ${}^8\text{Be} + {}^4\text{He}$ . If the forward rate is fast enough, however, than a population of  ${}^{12}\text{C}$  in this excited state can accumulate. The total rate to the ground state would then be  $n_{12^*} \Gamma_{\text{rad}} / \hbar$ , where  $n_{12^*}$  is the number density of  ${}^{12}\text{C}$  nuclei in the excited state.

To compute  $n_{12^*}$ , we again can use the equation for chemical equilibrium:  $\mu_8 + \mu_4 = \mu_{12^*} - Q_*$ . Here  $Q_* = -287 \text{ keV}$  is the difference in energy between the  ${}^4\text{He}$  and  ${}^8\text{Be}$  nuclei and the energy of  ${}^{12}\text{C}$  in the excited state. Again using eq. (6.7) to expand  $\mu$ , we obtain

$$n_{12^*} = \frac{n_{Q,12^*}}{n_{Q,4} n_{Q,8}} n_4 n_8 \exp \left( -\frac{287 \text{ keV}}{kT} \right). \quad (13.8)$$

Substituting for  $n_8$  using equation (13.6), this becomes

$$n_{12*} = \left( \frac{2\pi\hbar^2}{m_u k T} \right)^3 \left( \frac{12}{4^3} \right)^{3/2} \left( \frac{X_4 \rho}{4m_u} \right)^3 \exp \left( -\frac{397 \text{ keV}}{kT} \right). \quad (13.9)$$

Multiplying eq. (13.9) by  $\Gamma_{\text{rad}}/\hbar$  and scaling to  $\rho_5, T_8$ , we obtain the net rate, per unit volume, at which  ${}^3\text{He} \rightarrow {}^{12}\text{C}$ ,

$$\text{rate} = 4.37 \times 10^{31} \text{ cm}^{-3} \text{ s}^{-1} \frac{(X_4 \rho_5)^3}{T_8^3} \exp \left( -\frac{44.0}{T_8} \right). \quad (13.10)$$

Multiplying the rate by  $Q = 7.275 \text{ MeV}$ , the net  $Q$ -value, gives the volumetric heating rate  $\rho \varepsilon_{3\alpha}$ , or

$$\varepsilon_{3\alpha} \approx 5.1 \times 10^{21} \text{ erg g}^{-1} \text{ s}^{-1} \frac{X_4^3 \rho_5^2}{T_8^3} \exp \left( -\frac{44.0}{T_8} \right). \quad (13.11)$$

The temperature exponent is  $d \ln \varepsilon_{3\alpha} / d \ln T = 44.0/T_8 - 3$ , that is,  $\varepsilon_{3\alpha} \sim T^{41}$  at  $T = 10^8 \text{ K}$ .

### *The helium flash and the horizontal branch*

The extreme temperature sensitivity of the triple-alpha reaction motivates a return to analyzing the stability of reactions in a stellar environment. We saw in an earlier problem that the “gravithermal” specific heat

$$C_* \equiv T \left. \frac{\partial S}{\partial T} \right|_M < 0$$

for an ideal gas. The physical cause is that an increase in entropy leads to an increase in radius, and the resulting  $P dV$  work results in a reduced central temperature.

For conditions in low-mass stars at the time of helium ignition (nearly pure He at  $T \approx 10^8 \text{ K}, \rho \approx 10^5 \text{ g cm}^{-3}$ ); at that density the temperature at  $k_B T = \varepsilon_F$  is (eq. [6.27])  $4.1 \times 10^8 \text{ K}$ . Thus, He ignition takes place under semi-degenerate conditions. To understand how the star responds, let’s assume a homologous expansion—that is, one in which the ratios  $r(m)/R(M)$  remains constant. In other words, we assume that the structure of the star retains its functional form and we are merely “rescaling” our radial length. To compute a gravithermal specific heat, we use Jacobians (see Appendix A.1) to transform

$S(T, P)$  to  $S(T, M)$ :

$$\begin{aligned} T \left( \frac{\partial S}{\partial T} \right)_M &= T \frac{\partial(S, M)}{\partial(T, M)} \\ &= T \frac{\partial(S, M)}{\partial(T, P)} \frac{\partial(T, P)}{\partial(T, M)} \\ &= T \left( \frac{\partial S}{\partial T} \right)_P - T \left( \frac{\partial S}{\partial P} \right)_T \left( \frac{\partial M}{\partial T} \right)_P \left( \frac{\partial P}{\partial M} \right)_T \\ &= C_P \left[ 1 - \left( \frac{\partial T}{\partial P} \right)_S \left( \frac{\partial P}{\partial T} \right)_M \right], \end{aligned} \quad (13.12)$$

where in the last identity we have used

$$\left( \frac{\partial T}{\partial P} \right)_S \left( \frac{\partial S}{\partial T} \right)_P \left( \frac{\partial P}{\partial S} \right)_T = -1$$

and a similar expression relating  $P$ ,  $T$ , and  $M$ . We could continue to use this technique of Jacobians to further transform  $(\partial P / \partial T)_M$ ; but an easier method is to expand the equation of state as

$$\ln P = \chi_\rho \ln \rho + \chi_T \ln T, \quad (13.13)$$

where  $\chi_\rho \equiv (\partial \ln P / \partial \ln \rho)_T$ , and similarly for  $\chi_T$ . In an homologous expansion or contraction, the polytropic index stays constant, so that from equations (4.16) and (4.17) we have

$$\left( \frac{\partial \ln P}{\partial \ln R} \right)_M = -4, \quad \left( \frac{\partial \ln \rho}{\partial \ln R} \right)_M = -3. \quad (13.14)$$

Using these relations, we have

$$\begin{aligned} \left( \frac{\partial \ln T}{\partial \ln R} \right)_M &= \chi_T^{-1} \left[ \left( \frac{\partial \ln P}{\partial \ln R} \right)_M - \chi_\rho \left( \frac{\partial \ln \rho}{\partial \ln R} \right)_M \right] \\ &= \frac{-4 + 3\chi_\rho}{\chi_T}. \end{aligned}$$

We can use this along with equations (13.13) and (13.14) to obtain

$$\begin{aligned} \left( \frac{\partial \ln P}{\partial \ln T} \right)_M &= \left( \frac{\partial \ln P}{\partial \ln R} \right)_M \left( \frac{\partial \ln T}{\partial \ln R} \right)_M^{-1} \\ &= \frac{4\chi_T}{4 - 3\chi_\rho}. \end{aligned} \quad (13.15)$$

Inserting eq. (13.15) into equation (13.12), we finally get the expression for the gravothermal specific heat under a homologous expansion or contraction,

$$C_* = C_P \left[ 1 - \nabla_{\text{ad}} \frac{4\chi_T}{4 - 3\chi_\rho} \right]. \quad (13.16)$$

For an ideal gas,  $\chi_T = \chi_\rho = 1$  and  $\nabla_{\text{ad}} = 2/5$ , so that  $C_* < 0$ , as required for stability. As the core becomes degenerate, however,

$\chi_\rho \rightarrow 5/3$  and  $\chi_T \rightarrow 0$ ; as a result, the addition of heat to the core causes the temperature to rise, causing the rate of heating from nuclear reactions to increase even further. The ignition of helium in low-mass stars is therefore somewhat unstable and proceeds via “flashes.”

Eventually, the burning of helium heats the core enough that degeneracy is lifted, and the star settles onto its “helium main-sequence.” On an HR diagram, low-mass stars with core helium burning lie on the “horizontal branch.”

### The Asymptotic Giant Branch

After exhaustion of helium burning in the core, the star has a semi-degenerate C/O core, surrounded by a He burning shell with a superincumbent H-burning shell. Burning by these shells adds to the mass of the C/O core. As on the giant branch, the luminosity increases due to the high pressure at the edge of the core, and this high-luminosity produces a deep convective envelope, so that the star again lies along the Hayashi line in the HR diagram.

Because of the large gravitational acceleration at the edge of the core, the burning shells become thin in radial extent. This tends to make the burning unstable and leads to *thermal pulses*. To understand why the burning is unstable, let’s revisit equation (13.16). This equation still holds for expansion or contraction of a thin layer, with one critical difference: the volume of the shell is  $4\pi r_c^2 D$ , where  $D$  is the thickness of the shell and  $r_c$  is the core radius, which is fixed. Hence if we expand or contract the shell, keeping the mass in the shell fixed, the logarithmic change in density is

$$d \ln \rho = -\frac{dD}{D} = -\frac{dR}{D} = -\frac{R}{D} d \ln R.$$

Replacing the coefficient 3 of  $\chi_\rho$  with  $R/D$  in equation (13.16) gives us the specific heat during a homologous expansion or contraction of a shell of thickness  $D$ ,

$$C_{*,\text{shell}} = C_P \left[ 1 - \nabla_{\text{ad}} \frac{4\chi_T}{4 - (R/D)\chi_\rho} \right]. \quad (13.17)$$

Even for an ideal gas, if  $D < R/4$ , say, then  $C_{*,\text{shell}} > 0$  and the burning in the shell is thermally unstable.

When the burning shells are thin, the temperature, and hence rate of burning, become dependent on the local gravitational acceleration. Because the underlying core is degenerate, this means the luminosity depends almost entirely on the core mass: an empirical fit is<sup>1</sup>

$$\frac{L}{L_\odot} = 5.9 \times 10^4 \left( \frac{M_c}{M_\odot} - 0.52 \right). \quad (13.18)$$

<sup>1</sup> B. Paczyński. Evolution of Single Stars. I. Stellar Evolution from Main Sequence to White Dwarf or Carbon Ignition. *Acta Astronomica*, 20:47, 1970

The rate at which mass is added to the core is set by the rate at which  $^1\text{H}$  is processed into  $^4\text{He}$ ,

$$\frac{dM_c}{dt} = \frac{L}{q}. \quad (13.19)$$

Here  $q$  is the mass-specific energy release from the fusion of  $^1\text{H}$  into  $^4\text{He}$ :  $q = 26.72 \text{ MeV}/(4m_u) = 6.68 \times 10^{18} \text{ erg g}^{-1}$ . Thus as the core mass grows, the luminosity increases and the rate at which mass is added to the core increases. The high luminosity drives a strong wind from the stellar envelope. An empirical fit to the mass loss rate is<sup>2</sup>

$$\frac{dM}{dt} = -8.0 \times 10^{-13} M_\odot \text{ yr}^{-1} \left( \frac{L}{L_\odot} \frac{g_\odot}{g} \frac{R_\odot}{R} \right). \quad (13.20)$$

Here  $g = GM/R^2$  is the gravitational acceleration at the stellar surface. The coefficient has been increased beyond that in the original formula to fit the higher mass-loss rates observed from supergiants<sup>3</sup>. The envelope is thus consumed at the base by hydrogen and helium burning shells and expelled at the top by a radiative wind. This process ends when the envelope is consumed, leaving behind a degenerate C/O white dwarf that gradually cools.

<sup>2</sup> D. Reimers. On the absolute scale of mass-loss in red giants. I - Circumstellar absorption lines in the spectrum of the visual companion of Alpha-1 HER. *A&A*, 61:217–224, October 1977

<sup>3</sup> K.-P. Schröder and E. Sedlmayr. The galactic mass injection from cool stellar winds of the 1 to  $2.5 M_\odot$  stars in the solar neighbourhood. *A&A*, 366:913–922, February 2001

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EXERCISE 13.2— This exercise gives an illustration of the physics that sets the white dwarf initial mass function. Note that the numbers listed here are rather crude.

1. Combine equations (13.18) and (13.19) into a differential equation for the core mass as a function of time. Solve the equation. Since we don't know the initial core mass  $M_{c0}$  at the end of core He burning (other than that we are assuming it is greater than  $0.52 M_\odot$ ), let's leave that as a free parameter in the problem. What is the characteristic timescale for the core to increase in mass?
  2. Solve equation (13.20) for the *total* stellar mass as a function of time, assuming the initial total mass at the start of the AGB phase is  $M_0$ . To make the problem concrete, assume that the surface effective temperature on the AGB is fixed at 4000 K, and use the result of problem 1 to get the luminosity as a function of initial core mass,  $M_{c0}$ .
  3. For a star that starts its AGB phase with  $M_0 = 1.0 M_\odot$  and  $M_{c0} = 0.55 M_\odot$ , what is the final white dwarf mass?
-

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*Rotation and Magnetic Fields*



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*The Standard Solar Model*



# 16

## *Evolution of High-mass Stars*

Humphreys-Davidson limit

Atmospherical Eddington limit

Decoupling of core and envelope evolution

Nuclear burning stages

Helium is “burned” by  $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$

Deeper into ignition masses

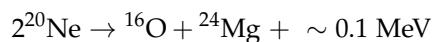
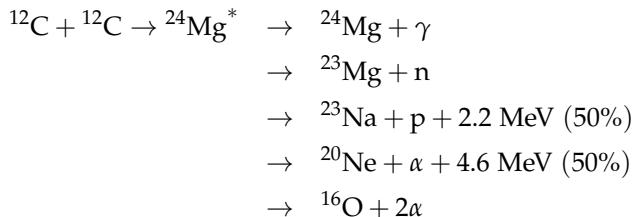
Neutrino-cooled stars - prevents the nuclear runaway in degenerate stars

Photodisintegration

Electron capture.

Nuclear Statistical equilibrium: problem from HKT.

Heavy ion-reactions: Carbon burning



*Neutrino cooling*

$$\epsilon_\nu \approx 1.1 \times 10^7 T_8^8 \text{ erg g}^{-1} \text{ s}^{-1} \quad (16.1)$$

*Electron capture*

Beta-decay:

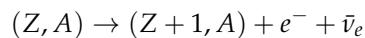


Table 16.1: Nuclear burning stages.

Fuel	Process	$T_{\text{ign}}$ [ $10^6$ K]	Ash	$Q_{\text{net}}$ [MeV/nucleon]	$n(T)$	$\tau_{\text{nucl}}$ yr
H	p-p chain	4	He	6.55	4	–
H	CNO cycle	15	He	6.25	18	$1.1 \times 10^7$
He	triple- $\alpha$	120	C,O	0.61	41	$2.0 \times 10^6$
C	C-fusion	600	Ne,Mg,Na,O	0.54	–	$2.0 \times 10^3$
Ne	Ne dissoc.	900	O,Mg,Si	0.002	–	0.7
O	O-fusion	1000	S,Si,P,Mg	0.30	–	2.6
Si	QSE	3000	Fe,Ni,Cr,Ti	<0.18	–	0.05

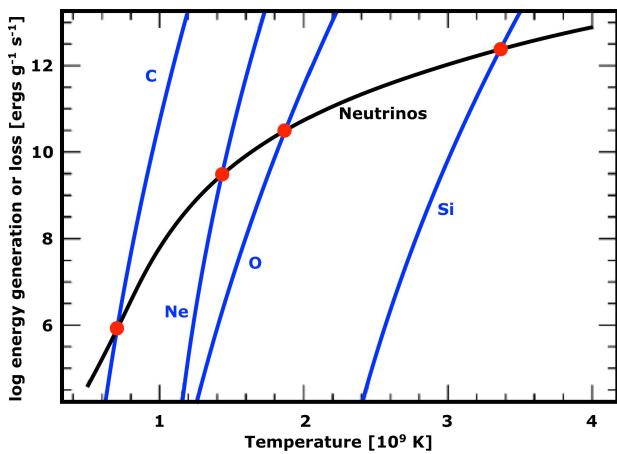


Figure 16.1: Nuclear energy generation and neutrino cooling.

Electron capture:

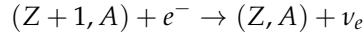


Figure 16.2: Electron capture threshold densities

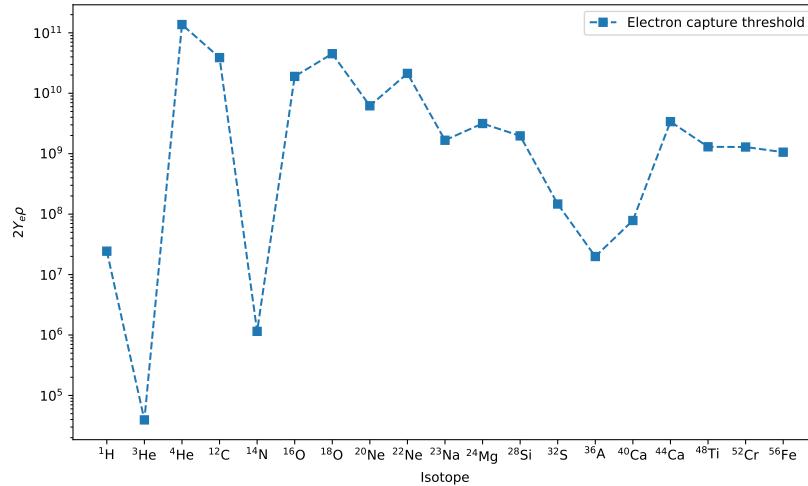
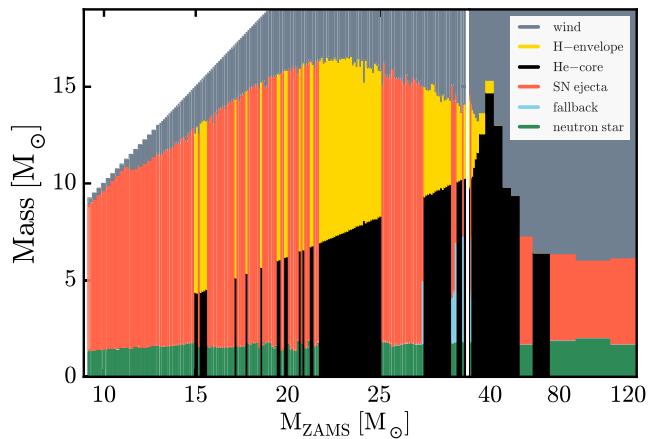


Figure 16.3: Final stellar mass vs. ZAMS mass



### Nuclear statistical equilibrium

The Saha equation is

$$\frac{n_+ n_-}{n_0} = \frac{g_+ g_-}{g_0} \left( \frac{m_- k_B T}{2\pi\hbar^2} \right)^{3/2} \exp \left( -\frac{Q}{k_B T} \right). \quad (16.2)$$



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*Core-collapse Supernovae*



# 18

## *Binaries*

### 18.1 Accretion

Suppose we have two objects orbiting a common center of mass. By convention, the more massive object is known as the *primary* and will be denoted by a subscript “1”; the less massive object is the *secondary* and is denoted by a subscript “2”. From Kepler’s law, the orbital separation is

$$a^3 = G(M_1 + M_2) \left( \frac{P}{2\pi} \right)^2, \quad (18.1)$$

with a numerical value  $a = 0.51 R_\odot (m_1 + m_2)^{1/3} (P/\text{hr})^{2/3}$ , where we’ve scaled our masses to solar values,  $m = M/M_\odot$ .

In a co-rotating frame, there is an equipotential surface with a saddle point, the *inner Lagrange point*, between the two stars. This surface forms two lobes, the *Roche lobes*, that meet at this point. See the crude sketch in Figure 18.1. Although the Roche lobes are not spherical, we can define the radius of an equivalent spherical volume; for the secondary, this is

$$\frac{R_2}{a} \approx 0.462 \left( \frac{M_2}{M_1 + M_2} \right)^{1/3}. \quad (18.2)$$

As an aside, equations (18.2) and (18.1) imply that the average density of the secondary, if it fills its Roche lobe, is

$$\bar{\rho} = \frac{3M_2}{4\pi R_2^3} \approx 111 \text{ g cm}^{-3} \left( \frac{\text{hr}}{P} \right)^2.$$

and does not depend explicitly on the masses of the two stars.

If matter is transferred from  $M_2$  to  $M_1$  how does the system respond? Let’s first write down the angular momentum of the system,

$$J = (M_1 a_1^2 + M_2 a_2^2) \omega = M_1 M_2 \left( \frac{G a}{M_1 + M_2} \right)^{1/2}. \quad (18.3)$$

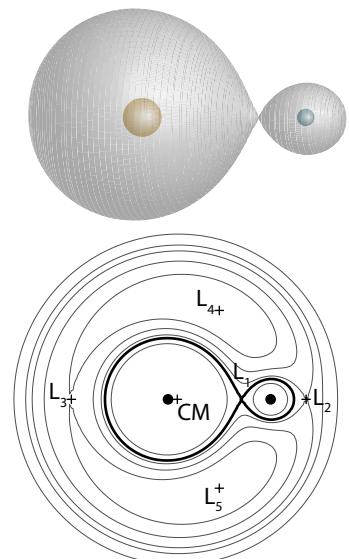


Figure 18.1: Sketch of the Roche lobes and potential for  $M_2 = 0.1M_1$ .

Here we've used equation (18.1) and the relations  $a_1 = aM_2/(M_1 + M_2)$ ,  $a_2 = aM_1/(M_1 + M_2)$ . Let's make the assumption that no mass is lost from the system and that the mass transfer  $\dot{M}$  is from  $M_2$  to  $M_1$ :

$$\begin{aligned}\dot{M}_1 + \dot{M}_2 &= 0; \\ \dot{M}_2 &< 0.\end{aligned}$$

Taking the logarithm of equation (18.3) and differentiating, we then obtain

$$\frac{\dot{a}}{a} = 2\frac{\dot{J}}{J} + 2\left(\frac{\dot{M}}{M_2}\right)\left(1 - \frac{M_2}{M_1}\right). \quad (18.4)$$

We see explicitly that for  $M_2 < M_1$ , the response of mass transfer is to increase the orbital separation  $a$  if there is no external torque on the system.

What about the size of the lobe that the secondary inhabits? There are two countervailing tendencies:  $a$  increases, which acts to increase  $R_2$ , but  $M_2$  decreases, which acts to decrease  $R_2$ . Taking the logarithm of eq. (18.2),

$$\frac{\dot{R}_2}{R_2} = \frac{\dot{a}}{a} + \frac{1}{3}\frac{\dot{M}_2}{M_2} = 2\frac{\dot{J}}{J} + 2\frac{\dot{M}}{M_2}\left(\frac{5}{6} - \frac{M_2}{M_1}\right). \quad (18.5)$$

Hence for  $M_2 < (5/6)M_1$ , the volume of the secondary's Roche lobe increases; if the secondary doesn't expand in response to mass loss and there are no external torques, then the secondary will lose contact with the inner Lagrange point and mass transfer will cease. Alternatively, if  $M_2 > (5/6)M_1$ , then the Roche lobe will clamp down on the secondary; this tends to drive the mass transfer at an even greater rate, and the process is unstable.

In general, there are three physical mechanisms for driving mass transfer.

*Gravitational radiation* For  $P \lesssim 2$  hr, gravitational radiation from the orbit produces a negative torque on the system:

$$\frac{\dot{J}}{J} = -\frac{32}{5}\frac{G^3}{c^5}\frac{M_1M_2(M_1 + M_2)}{a^4}. \quad (18.6)$$

Note that for this short of an orbital period, the binary consists of two degenerate stars (e.g., WD-WD, NS-WD).

*Magnetic braking* At somewhat longer periods  $P \lesssim 1$  d, the companion can be a main sequence star that has a tidally locked rotation period. Main-sequence stars have winds, and these winds carry angular momentum. Because of the tidal locking, this also introduces a negative torque on the system.

*Evolution of the secondary* For wider binary orbits, the secondary star can make contact with the Roche lobe as it becomes a giant star.

Finally, if the secondary is sufficiently evolved, it may have a strong enough wind that accretion can still occur, even if the orbit is so wide that the secondary doesn't fill its Roche lobe.

---

EXERCISE 18.1— Suppose we have a  $1.6 M_{\odot}$  neutron star and a  $0.8 M_{\odot}$  companion. The companion is evolved and has an effective temperature  $T_{\text{eff}} = 3000 \text{ K}$ . We'll assume that  $T_{\text{eff}}$  is fixed. On the giant branch, the luminosity is powered by hydrogen shell burning, the ashes of which are added to the core (i.e., the core mass  $M_c$  increases due to hydrogen burning). The energy released, per gram of hydrogen consumed, is  $Q = 6 \times 10^{18} \text{ erg g}^{-1}$ . For such an evolved giant the luminosity depends mainly on the core mass  $M_c$  and may be approximated as [Ritter, 1999]

$$\frac{L}{L_{\odot}} = 10^{6.3} \left( \frac{M_c}{M_{\odot}} \right)^8.$$

For this system, find the orbital period  $P$  and the mass transfer rate  $\dot{M}$  (in units of solar masses per year) if the secondary has a core mass  $M_c = 0.2 M_{\odot}$ .

---

Matter that crosses the inner Lagrange point will find itself in orbit about the primary. The material still has enough angular momentum that it won't fall directly onto the primary, and so an *accretion disk* will form. In order for the matter to accrete, there must be enough friction in the disk so that the gravitational energy can be radiated away and angular momentum transported outward.

---

EXERCISE 18.2— Suppose our binary consists of two white dwarfs with a short orbital period, less than 2 hours, so that gravitational radiation produces a torque on the system according to eq. (18.6). Because the secondary is also degenerate, its radius depends on mass as

$$R_2 = K \left( \frac{M_2}{M_{\odot}} \right)^{-1/3},$$

where  $K \approx 2 \times 10^9 \text{ cm}$ .

1. Show that if the secondary just fills its Roche lobe, then  $M_2 \propto P^{-1}$  and find  $M_2$  if  $P = 1 \text{ hr}$ .
  2. Show that
- $$\frac{\dot{M}}{M_2} = -\frac{\dot{J}/J}{2/3 - M_2/M_1}$$
3. Using the above relations and eq. (18.6), find  $\dot{M}$  if  $M_1 = 1 M_{\odot}$ . Scale the orbital period to units of 1 hour.
-

## 18.2 The Eddington Limit

There is a characteristic luminosity at which the pressure from radiation balances the gravitational force. This tends to act as a limit to accretion. To derive this limit, known as the *Eddington luminosity*, consider a spherically symmetric shell of matter. Radiation enters the shell and scatters isotropically (Thomson scatters) from electrons. The momentum flux (momentum per unit time per unit area) entering the shell is just

$$\mathbf{P} = \frac{1}{c} \frac{L}{4\pi r^2} \mathbf{e}_r.$$

Here  $L$  is the luminosity and  $r$  is the radius of the shell. Since the scattering is presumed isotropic, the rate at which the fluid element's momentum changes (the impulse imparted to it by the radiation) is just

$$\frac{dp}{dt} = P\sigma_{\text{Th}}n_e,$$

where  $\sigma_{\text{Th}}$  is the cross-section for Thomson scattering. Since  $dp/dt$  is just a force per unit volume, we can balance it with the gravitational force per unit volume,

$$\mathbf{f}_g = -\frac{GM\langle A \rangle m_u n_{\text{ion}}}{r^2} \mathbf{e}_r,$$

and solve for  $L$  to obtain

$$L_{\text{Edd}} = \frac{4\pi GM c}{(\sigma_{\text{Th}}/m_u) Y_e}. \quad (18.7)$$

Here I have used  $n_e = \langle Z \rangle n_{\text{ion}}$  and  $Y_e = \langle Z \rangle / \langle A \rangle$ . Note that  $L_{\text{Edd}}$  is independent of distance from the star. Numerically,

$$L_{\text{Edd}} = 1.3 \times 10^{38} \text{ erg s}^{-1} \left( \frac{M}{M_\odot} \right) Y_e^{-1} = 3.2 \times 10^4 L_\odot \left( \frac{M}{M_\odot} \right) Y_e^{-1}.$$

We can now look at what the luminosity supplied by accretion is. First there is the just the gravitational energy release. The gravitational release, per nucleon, is

$$\frac{GMm_u}{R} = \begin{cases} 0.069 \text{ MeV} & \text{WD with } R = 2 \times 10^9 \text{ cm} \\ 138 \text{ MeV} & \text{NS with } R = 10^6 \text{ cm} \end{cases}$$

In both cases we used  $M = 1 M_\odot$ . For hydrogen-rich accretion onto a white dwarf, steady nuclear burning (assuming it exists!) will dwarf the luminosity from accretion. The situation is reversed for a neutron star, for which the energy from nuclear burning is a perturbation to the large release of gravitational binding energy.

We can compute the accretion rate at which the luminosity supplied by either nuclear burning or release of gravitational binding

energy would equal the Eddington value:

$$\dot{M}_{\text{Edd}} = \begin{cases} \frac{4\pi GMc}{Q(\sigma_{\text{Th}}/m_u)Y_e} & \text{nuclear, } L = \dot{M}Q \\ \frac{4\pi R c}{(\sigma_{\text{Th}}/m_u)Y_e} & \text{gravitational, } L = \frac{GM\dot{M}}{R} \end{cases}.$$

For nuclear burning, taking  $M = 1 M_\odot$ ,  $Q = 6 \times 10^{18} \text{ erg g}^{-1}$  (hydrogen burning), and  $Y_e = 1$  gives  $\dot{M}_{\text{Edd}} = 3.3 \times 10^{-7} M_\odot \text{ yr}^{-1}$ .

For gravitational power, taking  $R = 10^6 \text{ cm}$  and  $Y_e = 1$  gives  $\dot{M}_{\text{Edd}} = 1.5 \times 10^{-8} M_\odot \text{ yr}^{-1}$ . Observed systems do in fact have inferred mass transfer rates that are typically less than these values.

**EXERCISE 18.3**— Consider a star surrounded by an accretion disk of matter. In the disk, which consists of ionized gas, a fluid element gradually spirals inward, so that at each point in the disk, it is approximately in a circular orbit. Show that under these conditions, the energy radiated, per unit mass, by the fluid element over its lifetime in the disk is  $GM/(2R)$ .



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*Type Ia Supernovae*



## 20

# Stellar Pulsations

In this chapter, we'll linearize the perturbed continuity and momentum equations and solve for the frequencies of the normal modes for a star.

### 20.1 Adiabatic, radial pulsations

Imagine that we perturb our fluid in some way. As described in section 9.2, we can describe the *Eulerian* perturbation in some fluid property  $f$ :

$$\Delta f \equiv f(\mathbf{r}, t) - f_0(\mathbf{r}, t), \quad (20.1)$$

where the subscript “o” denotes the unperturbed quantity. Said another way,  $\Delta f$  describes the change, under our perturbation, in some property of the fluid at a fixed location.

We may also describe our perturbation as a *Lagrangian* one, where we compare the same fluid element in both the perturbed and unperturbed systems:

$$\delta f \equiv f(\mathbf{r}, t) - f_0(\mathbf{r}_0, t). \quad (20.2)$$

Under a Lagrangian perturbation the fluid element in the perturbed system in general has a different position  $\mathbf{r}$  than in the unperturbed system,  $\mathbf{r}_0$ .

The two perturbations are related to one another via

$$\delta f = \Delta f + (\delta \mathbf{r} \cdot \nabla) f_0. \quad (20.3)$$

There are a few useful commutation relations that are easily proved:

$$\partial_t \Delta f = \Delta (\partial_t f), \quad (20.4)$$

$$\nabla \Delta f = \Delta \nabla f, \quad (20.5)$$

$$\frac{D}{Dt} \delta f = \delta \frac{Df}{Dt}. \quad (20.6)$$

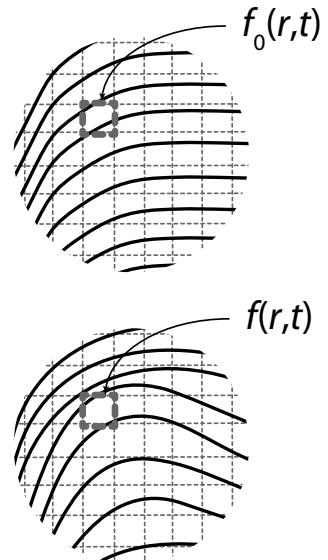


Figure 20.1: An Eulerian perturbation: we compare quantities at corresponding locations.

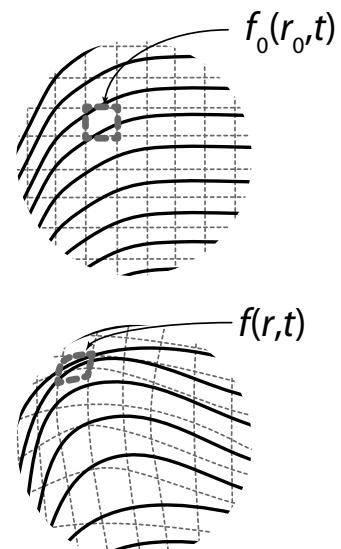


Figure 20.2: A Lagrangian perturbation: we compare quantities for corresponding fluid elements.

And there are operations that do not commute:

$$\partial_t \delta f \neq \delta (\partial_t f), \quad (20.7)$$

$$\nabla \delta f \neq \delta \nabla f, \quad (20.8)$$

$$\frac{D}{Dt} \Delta f \neq \Delta \frac{Df}{Dt}. \quad (20.9)$$

One can further show that  $\delta \mathbf{u} = (D/Dt)\delta \mathbf{r}$ . Also, if the fluid has unperturbed velocity  $\mathbf{u} = 0$ , then  $\Delta \mathbf{u} = \delta \mathbf{u}$ . Finally, for purely radial motion, we can introduce the Lagrangian mass coordinate  $m$ , in which case  $\partial_m \delta f = \delta(\partial_m f)$  and  $\partial_m \Delta f \neq \Delta(\partial_m f)$ .

WE ARE NOW READY TO USE THESE COMMUTATION RELATIONS TO DERIVE a *linear adiabatic wave equation*. By linear, we mean that we shall only keep terms to first order in  $\delta$ . By adiabatic, we mean that we shall only consider the equations of continuity and momentum, and we shall relate the density and pressure perturbations via

$$\frac{\delta P}{P} = \Gamma_1 \frac{\delta \rho}{\rho}. \quad (20.10)$$

Here  $\Gamma_1 \equiv (\partial \ln P / \partial \ln \rho)_s$ .

For simplicity, we'll start with purely radial oscillations. First, let's perturb the equation of continuity, expressed in Lagrangian form (eq. [3.53]),

$$\frac{\partial \ln r}{\partial m} = \frac{1}{4\pi r^3 \rho}.$$

We apply a Lagrangian perturbation to both sides of this equation and expand the right-hand side to first order in  $\delta r$  and  $\delta \rho$ . Since  $\delta$  and  $\partial/\partial_m$  commute, we can interchange them:

$$\begin{aligned} \frac{\partial}{\partial m} \left( \frac{\delta r}{r} \right) &= \delta \left( \frac{\partial \ln r}{\partial m} \right) \\ &= \delta \left( 4\pi r^3 \rho \right)^{-1} \\ &= \left( 4\pi r^3 \rho \right)^{-1} \left( -3 \frac{\delta r}{r} - \frac{\delta \rho}{\rho} \right). \end{aligned}$$

Moving  $(4\pi r^3 \rho)$  to the left-hand side of the equation, and recognizing that

$$4\pi r^3 \rho \frac{\partial}{\partial m} = r \frac{\partial m}{\partial r} \frac{\partial}{\partial m} = r \frac{\partial}{\partial r},$$

we have our first equation,

$$r \frac{\partial}{\partial r} \left( \frac{\delta r}{r} \right) = -3 \frac{\delta r}{r} - \frac{\delta \rho}{\rho}. \quad (20.11)$$

Next, we can perturb the force equation (eq. [3.54])

$$\frac{D^2 r}{Dt^2} = -\frac{Gm}{r^2} - 4\pi r^2 \frac{\partial P}{\partial m}.$$

If the unperturbed state is taken to have  $D\mathbf{r}_0/Dt = D^2\mathbf{r}_0/Dt^2 = 0$ , then a similar linearization yields

$$\rho r \frac{D^2}{Dt^2} \left( \frac{\delta r}{r} \right) = - \frac{\partial P}{\partial r} \left( 4 \frac{\delta r}{r} + \frac{\delta P}{P} \right) - P \frac{\partial}{\partial r} \left( \frac{\delta P}{P} \right), \quad (20.12)$$

which is our second equation.

To proceed further, we write

$$\frac{\delta r}{r} = \zeta(r) \exp(i\sigma t),$$

so that the left-hand side of equation (20.12) becomes  $-\rho r \sigma^2 \zeta(r) e^{i\sigma t}$ , and we can make the substitution  $\partial_r(\delta r/r) \rightarrow e^{i\sigma t} (\partial_r \zeta)$ . We additionally eliminate  $\delta\rho/\rho$  from equation (20.11) using the adiabatic condition, eq. (20.10) and make use of the zeroth-order momentum equation  $dP/dr = -\rho Gm/r^2$  to obtain

$$\frac{d}{dr} \zeta = - \frac{1}{r} \left( 3\zeta + \frac{1}{\Gamma_1} \frac{\delta P}{P} \right) \quad (20.13)$$

$$\frac{d}{dr} \left( \frac{\delta P}{P} \right) = \frac{1}{\lambda_P} \left[ \left( 4 + \sigma^2 \frac{r^3}{Gm} \right) \zeta + \frac{\delta P}{P} \right]. \quad (20.14)$$

Here we introduce the pressure scale height (in the unperturbed system)  $\lambda_P \equiv -(d \ln P / dr)^{-1} = Pr^2 / (\rho Gm)$ . Multiply equation (20.13) by  $\Gamma_1 Pr^4$  and then differentiate with respect to  $r$ , using equation (20.14) to eliminate the spatial derivative of  $\delta P/P$  and equation (20.13) to eliminate  $\delta P/P$  to obtain

$$\frac{d}{dr} \left[ \Gamma_1 Pr^4 \frac{d}{dr} \zeta \right] + \left\{ r^3 \frac{d}{dr} [(3\Gamma_1 - 4) P] \right\} \zeta + \sigma^2 (r^4 \rho) \zeta = 0. \quad (20.15)$$

Notice here that we have *not* assumed that  $\Gamma_1$  is a constant.

Equation (20.15) has the form

$$\mathcal{L}\zeta(r) + \sigma^2 w(r)\zeta(r)$$

where

$$\mathcal{L} \equiv \frac{d}{dr} \left[ u(r) \frac{d\zeta}{dr} \right] + q(r)\zeta(r),$$

with  $u(r) = \Gamma_1 Pr^4$ ,

$$q(r) = r^3 \frac{d}{dr} [(3\Gamma_1 - 4) P],$$

and  $w(r) = r^4 \rho$ . For the imposed boundary conditions, there will in general be solutions for only certain eigenvalues  $\sigma^2$ . Note that  $u(r) > 0$  on the interval  $0 < r < R$ . Furthermore, we require that  $\zeta$  and  $d\zeta/dr$  be finite at  $r = 0$  and  $r = R$ , which means that if  $\zeta_i$  and  $\zeta_j$  are solutions of eq. (20.15), then

$$u(r)\zeta_i^* \frac{d\zeta_j}{dr} \Big|_{r=0} = u(r)\zeta_i^* \frac{d\zeta_j}{dr} \Big|_{r=R} = 0. \quad (20.16)$$

Using these boundary conditions and the form of the operator  $\mathcal{L}$ , we find that

$$\begin{aligned}\int_0^R dr \zeta_i^* \mathcal{L} \zeta_j &= u \zeta_i^* \frac{d\zeta_j}{dr} \Big|_{r=0}^{r=R} - \int_0^R dr \frac{d\zeta_i^*}{dr} u \frac{d\zeta_j}{dr} + \zeta_j q(r) \zeta_i^* \\ &= -u(r) \zeta_j \frac{d\zeta_i^*}{dr} \Big|_{r=0}^{r=R} + \int_0^R dr \zeta_j \frac{d}{dr} \left[ u \frac{d\zeta_i^*}{dr} \right] + \zeta_j q(r) \zeta_i^* \\ &= \int_0^R dr \zeta_j \mathcal{L} \zeta_i^*.\end{aligned}$$

The operator  $\mathcal{L}$  is thus Hermitian. As a result, the eigenvalues  $\sigma^2$  are real and denumerable. There is a minimum eigenvalue  $\sigma_0^2$ . The eigenfunctions corresponding to these eigenvalues are orthogonal in the following sense: if  $\sigma_i^2$  and  $\sigma_j^2$  are eigenvalues of equation (20.15) and  $\zeta_1, \zeta_2$  their corresponding eigenfunctions, then

$$\int_0^R dr w(r) \zeta_i \zeta_j = \int_0^R dr r^4 \rho \zeta_i \zeta_j = 0 \quad \text{if } \sigma_i^2 \neq \sigma_j^2. \quad (20.17)$$

Solutions with larger eigenvalues have more nodes.

## 20.2 Adiabatic non-radial pulsations

Now that we've warmed up with the purely radial pulsations, we'll do the more general case. Rather than going through the separation of variables in spherical coordinates, we'll keep things simple and cartesian. This amounts to looking at a small box in the star. We will also make an *ansatz* that the fluid perturbations don't change the gravitational potential<sup>1</sup>. In this case, the (constant) gravitational acceleration  $\mathbf{g} = -g \mathbf{e}_r$  defines the local vertical, so we will separate our equations into a radial direction, labeled by "r", and a transverse direction, labeled by "t".

Let us first perturb the equation of continuity,

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{u}) = \frac{D\rho}{Dt} + \rho \nabla \cdot \mathbf{u} = 0. \quad (20.18)$$

We take our unperturbed state to be independent of time with  $\mathbf{u} = \mathbf{0}$ .

If we take the Lagrangian perturbation of eq. (20.18), we have

$$\frac{D}{Dt} (\delta\rho + \rho \nabla \cdot \boldsymbol{\xi}) = 0,$$

with  $\mathbf{u} = D\boldsymbol{\xi}/Dt$ . Setting the constant of integration to 0 reduces this to

$$\frac{\delta\rho}{\rho} = -\nabla \cdot \boldsymbol{\xi}.$$

The linearization of the momentum equation,

$$\frac{D^2 \boldsymbol{\xi}}{Dt^2} = -\frac{1}{\rho} \nabla \Delta P + \frac{\Delta\rho}{\rho} \mathbf{g}, \quad (20.19)$$

<sup>1</sup> This is known as the **Cowling approximation**.

looks very similar to what we did in deriving a condition for convection, § 9.2. This time, however, we won't impose the condition that  $\Delta P = 0$ . When we were looking at convective instabilities, we were interested in low-frequency perturbations, in which the pressure has time to equilibrate. Keeping the terms with  $\Delta P$ , the perturbed momentum equation becomes

$$\frac{D^2\xi}{Dt^2} = -\frac{1}{\rho} \nabla \Delta P + g \underbrace{\frac{1}{\Gamma_1} \frac{\Delta P}{P}}_{\text{buoyancy from density perturbation}} + g(\xi \cdot \nabla) \overbrace{\left[ \frac{1}{\Gamma_1} \ln P - \ln \rho \right]}^{-\mathcal{A}}. \quad (20.20)$$

If  $\Delta P \rightarrow 0$ , this reduces to equation (9.20) with  $\mathcal{A}$  being the **Schwarzschild discriminant**.

To decompose our perturbed system into normal modes, with the perturbed quantities varying in time as  $\exp(i\omega t)$ , we first note that in the unperturbed system  $\ln P$  and  $\ln \rho$  depend only on  $r$ ; and nothing depends on the transverse directions. As a result, we can impose periodic boundary conditions in the transverse direction and write

$$\begin{aligned} \Delta P(x, t) &= \Delta P(r) \exp(ik_t \cdot x_t) \exp(i\omega t), \\ \Delta \xi(x, t) &= \xi(r) \exp(ik_t \cdot x_t) \exp(i\omega t). \end{aligned}$$

Here  $k_t$  is a transverse wavenumber and  $x_t$  are the transverse coordinates. The transverse component of equation (20.20) is thus

$$\rho \omega^2 \xi_t = ik_t \Delta P, \quad (20.21)$$

and the radial equation of motion is

$$\rho \omega^2 \xi_r = \partial_r \Delta P + \frac{g\rho}{\Gamma_1 P} \Delta P + \rho N^2 \xi_r. \quad (20.22)$$

We've re-introduced the Brunt-Väisälä frequency<sup>2</sup>

<sup>2</sup> We take  $\Gamma_1$  to be constant.

$$N^2 = -g \frac{d}{dr} \mathcal{A} = g \left( \frac{1}{\Gamma_1} \frac{d}{dr} \ln P - \frac{d}{dr} \ln \rho \right),$$

just as was done in analyzing convection.

We still have more unknowns than equations. Next, we'll use the perturbed equation of continuity,

$$\frac{\delta \rho}{\rho} = -\nabla \cdot \xi = -\partial_r \xi_r - k_t \cdot \xi_t. \quad (20.23)$$

For adiabatic perturbations,

$$\begin{aligned} \frac{\delta \rho}{\rho} &= \frac{1}{\Gamma_1} \frac{\delta P}{P} \\ &= \frac{1}{\Gamma_1} \left( \frac{\Delta P}{P} + \xi \cdot \nabla \ln P \right) \\ &= \frac{1}{\Gamma_1} \frac{\Delta P}{P} - \xi_r \frac{1}{\Gamma_1 \lambda_P}. \end{aligned}$$

Here we used  $\xi \cdot \nabla \ln P = \xi_r (d/dr) \ln P = -\xi_r \rho g / P = -\xi_r / \lambda_P$ , where  $\lambda_P$  is the pressure scale height. It makes physical sense that this scale will enter: modes with radial wavelengths much less than  $H$  shouldn't be affected by the background stratification.

Inserting the expression for  $\delta\rho/\rho$  in equation (20.23) and using equation (20.21) to eliminate  $\xi_t$ , we obtain

$$\partial_r \xi_r = \frac{1}{\omega^2 \Gamma_1 P} \left( \frac{k_t^2 \Gamma_1 P}{\rho} - \omega^2 \right) \Delta P + \frac{\xi_r}{\Gamma_1 \lambda_P}. \quad (20.24)$$

The quantity  $S^2 = k_t^2 \Gamma_1 P / \rho = k_t^2 c_s^2$  is called the **Lamb frequency**.

Collecting equations (20.24) and (20.22), we now have two coupled first order differential equations for  $\xi_r(r)$  and  $\Delta P(r)$ ,

$$\partial_r \xi_r = \frac{1}{\Gamma_1 P} \frac{S^2 - \omega^2}{\omega^2} \Delta P + \frac{\xi_r}{\Gamma_1 \lambda_P} \quad (20.25)$$

$$\partial_r \Delta P = \rho(\omega^2 - N^2) \xi_r - \frac{\Delta P}{\Gamma_1 \lambda_P}. \quad (20.26)$$

From the form of these equations, let's try a solution in the form

$$\Delta P(r) = \Delta P e^{Kr}; \quad \xi_r(r) = \xi_r e^{Kr}. \quad (20.27)$$

If  $K$  is imaginary, then we have oscillatory solutions. If  $K$  is real, then we must choose  $K$  so that the radial functions decay exponentially, or **evanesce**. Substituting equation (20.27) into eq. (20.26) gives

$$\Delta P = \left[ \frac{\rho(\omega^2 - N^2)}{1 + 1/K\lambda_P} \right] \frac{\xi_r}{K}.$$

Substituting this expression for  $\Delta P$  into eq. (20.25) and eliminating  $\xi_r$  gives us the dispersion equation,

$$K^2 = \frac{k_t^2}{S^2} \frac{(S^2 - \omega^2)(\omega^2 - N^2)}{(1 + 1/K\lambda_P)\omega^2} + \frac{1}{K\lambda_P}, \quad (20.28)$$

which relates the mode frequency  $\omega$  to the transverse wavenumber  $k_t$  and the radial wavenumber  $K$ .

In the limit where the radial wavelength is much less than a pressure scale height,  $|K|\lambda_P \gg 1$ , equation (20.28) simplifies to

$$K^2 = \frac{k_t^2 (S^2 - \omega^2)(\omega^2 - N^2)}{S^2 \omega^2}. \quad (20.29)$$

If  $N^2 < \omega^2 < S^2$ ,  $K^2 > 0$  and the radial perturbations evanesce. To have a mode, we require either that  $\omega^2 < N^2, S^2$  or  $\omega^2 > N^2, S^2$ .

The analysis is similar in spherical coordinates, except that instead of plane waves,  $e^{k_t \cdot x_t}$ , we'll have spherical harmonics  $Y_{\ell m}$ . In the above dispersion relation, make the substitution  $k_t^2 \rightarrow \ell(\ell+1)/r^2$  and  $S^2 \rightarrow S_\ell^2$ .

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EXERCISE 20.1— Use dimensional analysis and virial estimates to sketch  $S$  (for  $\ell = 1$ ) and  $N$  on a plot of frequency against radius for the sun. Indicate where modes with  $\omega^2 < N^2, S^2$  can propagate, and likewise for  $\omega^2 > N^2, S^2$ .

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## 21

# *Stellar Atmospheres*

In the atmosphere of the star, the optical depth approaches unity, and we can no longer treat the radiation field as being isotropic. Let's consider the time-independent problem ( $\partial_t \rightarrow 0$ ) of a plane-parallel atmosphere. The *optical depth* for an outward-directed ray is

$$\tau_\nu = \int_z^\infty \rho \kappa_\nu dz'. \quad (21.1)$$

Now the optical depth just the distance divided by the mean free path. Clearly, when  $\tau_\nu < 1$ , a photon has a good chance of reaching a distant observer without any further interactions with the stellar matter. As a result, the intensity takes its final form around  $\tau_\nu \approx 1$ , and this defines the stellar *photosphere*. To get some of the basic properties of the photosphere, rewrite eq. (21.1) in differential form,

$$\frac{d\tau}{dz} = -\rho\kappa. \quad (21.2)$$

This is for a crude estimate, so we neglect the frequency dependence for now. We can use equation (21.2) along with hydrostatic balance to get an estimate of the photospheric pressure,

$$\frac{dP}{d\tau} = - \left( \frac{d\tau}{dz} \right)^{-1} \rho g = \frac{g}{\kappa}. \quad (21.3)$$

Thus, at  $\tau \approx 1$ , the pressure is  $P_{ph} \approx g/\kappa$ . Since the flux at the photosphere is  $\sigma_{SB} T_{eff}^4$ , we would expect that the local temperature is  $T \approx T_{eff}$ .

### 21.1 The Eddington Approximation

To get an analytical approximation for the atmosphere, we'll first redefine our transfer equation in terms of optical depth (eq. [7.18]).

Here however, we will take the optical depth to be along the  $z$ -direction, so we define  $\mu = \hat{k} \cdot \hat{n}$ , where  $\hat{n}$  is the direction along

the ray. The equation of transfer then becomes

$$\mu \frac{\partial I_\nu}{\partial \tau_\nu} = I_\nu - S_\nu, \quad (21.4)$$

where

$$S_\nu \equiv \frac{1}{\kappa_\nu} \left( \frac{\varepsilon_\nu}{4\pi} + \kappa_\nu^{\text{sca}} J_\nu \right) \quad (21.5)$$

is the *source function*. In local thermodynamical equilibrium (LTE), we can write  $S_\nu = (1 - A_\nu)B_\nu + A_\nu J_\nu$ , where  $A_\nu \equiv \kappa_\nu^{\text{sca}} / \kappa_\nu$  is the *albedo*. Recall that  $J_\nu = (4\pi)^{-1} \int d\Omega I_\nu$  is the angle-average of  $I_\nu$ .

We noted that in thermal equilibrium,  $P_\nu = c^{-1} \int_{-1}^1 d\mu \mu^2 I_\nu = u_\nu / 3$ . This relation holds even when the radiation is not thermal, so long as it is isotropic to terms linear in  $\mu$ . To make this concrete, suppose we write

$$I_\nu(\mu) = I_\nu^{(0)} + \mu I_\nu^{(1)} + \mu^2 I_\nu^{(2)} + \dots$$

Here we are assuming that terms marked (0) are much larger than terms marked (1), etc. To lowest order, the energy density, flux, and momentum flux are then

$$\begin{aligned} u_\nu &= \frac{2\pi}{c} \int_{-1}^1 d\mu I_\nu(\mu) = \frac{4\pi}{c} I_\nu^{(0)}, \\ F_\nu &= 2\pi \int_{-1}^1 d\mu \mu I_\nu(\mu) = \frac{4\pi}{3} I_\nu^{(1)}, \\ P_\nu &= \frac{2\pi}{c} \int_{-1}^1 d\mu \mu^2 I_\nu(\mu) = \frac{4\pi}{3c} I_\nu^{(0)} = \frac{u_\nu}{3}. \end{aligned}$$

The *Eddington approximation* then consists of treating the radiation field as if its anisotropy is linear in  $\mu$  *everywhere*, so that the above relations hold; in particular, it means assuming that  $P_\nu = u_\nu / 3$  everywhere.

## 21.2 A Grey Atmosphere

Finally, to get an analytical approximation to the structure of the solar atmosphere, let's consider a grey atmosphere in LTE, i. e., one for which  $\kappa_\nu^{\text{abs}} = \kappa^{\text{abs}}$  and  $\kappa_\nu^{\text{sca}} = \kappa^{\text{sca}}$  are independent of frequency. Equation (21.4) can then be integrated over all frequencies to become

$$\mu \frac{\partial I}{\partial \tau} = I - S. \quad (21.6)$$

Integrating over all angles (note that we can pull the derivative wrt  $\tau$  out of the integral) gives

$$\frac{1}{4\pi} \frac{\partial F}{\partial \tau} = J - S = 0. \quad (21.7)$$

Why does the right-hand side vanish? Note that  $S - J = (1 - A)(B - J)$ . Clearly  $S = J$  if  $A = 1$  (a pure scattering atmosphere). If  $A \neq 1$ , so that there is some absorption, then the condition of detailed balance, equation (7.15), implies that  $\varepsilon_\nu = 4\pi\kappa^{\text{abs}}B_\nu(T)$ ; inserting this into equation (7.14), factoring out the constant  $\kappa^{\text{abs}}$ , and integrating over  $\nu$  implies that  $B - J = 0$ , and hence  $S - J = 0$ . Note that  $J = B$  does *not* necessarily imply that  $I_\nu = B_\nu$ !

Now multiply equation (21.6) by  $\mu$  and integrate over  $2\pi d\mu$  to obtain

$$c \frac{\partial P}{\partial \tau} = F, \quad (21.8)$$

the integral over  $\mu S$  vanishing because it is odd in  $\mu$ . Equation (21.7) implies that  $F$  is constant; hence we can integrate equation (21.8) at once to obtain

$$cP = F(\tau + \tau_0), \quad (21.9)$$

where  $\tau_0$  is a constant of integration. Of course, this does help us yet; all we have done is introduce a new variable  $P$ , the radiation pressure. This is where the Eddington approximation comes in. We set  $P = u/3 = 4\pi J/(3c)$  in equation (21.9) to obtain  $4\pi J = 3F(\tau + \tau_0)$ . Since  $J = S$ , we can then write equation (21.6) as

$$\mu \frac{\partial I}{\partial \tau} = I - \frac{3}{4\pi} F(\tau + \tau_0). \quad (21.10)$$

Since  $F$  is constant, this first-order differential equation is now solvable,

$$\begin{aligned} I(\mu, \tau = 0) &= \frac{1}{\mu} \int_0^\infty \frac{3}{4\pi} F(\tau + \tau_0) e^{-\tau/\mu} d\tau, \\ &= \frac{3}{4\pi} F(\mu + \tau_0). \end{aligned} \quad (21.11)$$

Now at  $\tau = 0$ , all of the flux must be outward-directed ( $\mu > 0$ ), so  $I(\mu < 0, \tau = 0) = 0$  if the star is not irradiated by another source. Note that the Eddington approximation is clearly violated here. Still, we will see later that this approximation is not too terrible.

To determine  $\tau_0$ , multiply  $I(\mu, \tau = 0)$  by  $\mu$  and integrate equation (21.11) over all angles to find

$$F = 2\pi \int_0^1 \mu I(\mu, 0) d\mu = \frac{1}{2} \int_0^1 3F(\mu + \tau_0) \mu d\mu = F \left( \frac{1}{2} + \frac{3}{4}\tau_0 \right). \quad (21.12)$$

We therefore find  $\tau_0 = 2/3$ . Now, since we are in LTE,  $P = aT^4/3$ . Further, let us define an effective temperature by the relation  $F = \sigma_{\text{SB}} T_{\text{eff}}^4$ . Substituting these definitions and the value of  $\tau_0$  into equation (21.9) gives us the atmospheric temperature structure,

$$T^4(\tau) = \frac{3}{4} T_{\text{eff}}^4 \left( \tau + \frac{2}{3} \right). \quad (21.13)$$

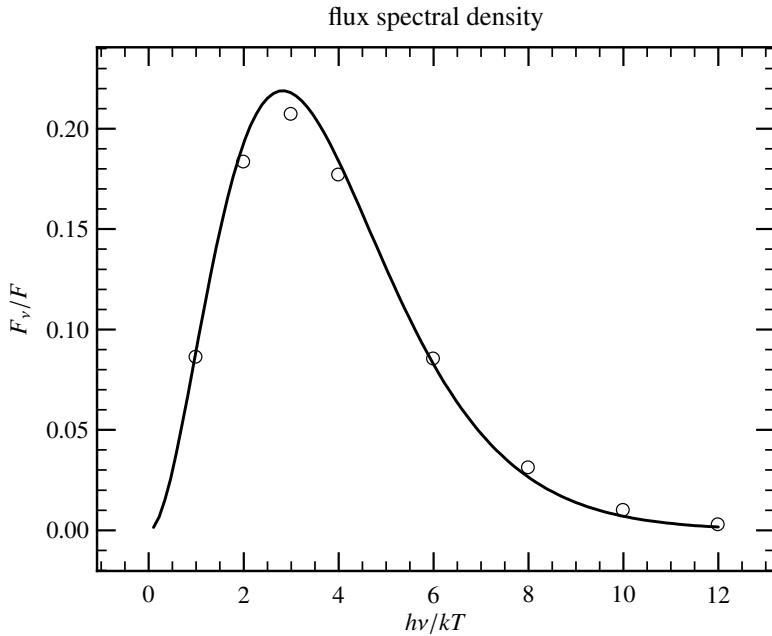


Figure 21.1: Spectral distribution from a grey atmosphere. The open circles are from Chandrasekhar, *Radiative Transfer*; the solid line is the Planck distribution.

Thus  $T(\tau = 0) = 2^{-1/4} T_{\text{eff}}$  and  $T(\tau = 2/3) = T_{\text{eff}}$ .

To get the spectral distribution, go back to equation (21.4) and (assuming the atmosphere has some absorption so that the matter and radiation can come into equilibrium) insert  $S_\nu = B_\nu(T)$ ; solving for  $I_\nu$  at  $\tau = 0$  then gives

$$I_\nu(\mu, \tau = 0) = \frac{1}{\mu} \int_0^\infty B_\nu[T(\tau)] e^{-\tau/\mu} d\tau. \quad (21.14)$$

A plot of the spectral distribution for the emergent flux is shown (*open circles*) in Fig. 21.1. For comparison, a plot of the Planck distribution (*solid line*) is also shown. Both fluxes are normalized to the total flux. Note that  $I_\nu(\mu, \tau = 0)$  depends on angle; rays propagating at a slant will have a lower intensity. As a result, when we observe the sun, the edge of the visible disk appears darker than the center, a phenomenon known as *limb darkening*.

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EXERCISE 21.1— Compute the reduction in intensity  $I$  as a function of viewing angle.

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### 21.3 Some examples

#### Requirement for convection in the atmosphere

Let's construct a simple atmosphere model using our temperature structure, eq. (21.13). Although we assume a grey opacity, we will

let it vary with temperature and density,  $\kappa(\rho, T) = \kappa_0 \rho^r T^s$ . For an ideal gas, we can rewrite this in terms of pressure and temperature,  $\kappa(P, T) = \kappa_0 (\mu m_u / k)^r P^r T^{s-r}$ . Substituting this into equation (21.3), and using equation (21.13), we obtain,

$$\frac{dP}{d\tau} = \frac{g}{\kappa_0} \left( \frac{k}{\mu m_u} \right)^r P^{-r} \left[ \frac{3}{4} T_{\text{eff}}^4 \left( \tau + \frac{2}{3} \right) \right]^{(r-s)/4}.$$

This is easily integrated: we'll take  $P(\tau = 0) = 0$  and obtain

$$P(\tau) = (\text{const}) \tau^{(r-s+4)/(1+r)}$$

so that

$$\frac{d \ln P}{d\tau} = \frac{r-s+4}{4(1+r)\tau}. \quad (21.15)$$

For the temperature,

$$\frac{d \ln T}{d\tau} = \frac{1}{4} \frac{d \ln T^4}{d\tau} = \frac{1}{4(\tau + 2/3)}. \quad (21.16)$$

We now combine eqn. (21.15) and (21.16) to obtain

$$\frac{d \ln T}{d \ln P} = \left( \frac{1+r}{r-s+4} \right) \left( \frac{\tau}{\tau + 2/3} \right). \quad (21.17)$$

For convection to happen,  $d \ln T / d \ln P > (\partial \ln T / \partial \ln P)_s = 1/(1+n)$ , where  $n = 3/2$  for an ideal gas. That is,

$$n > \frac{r-s+4}{1+r} - 1 = \frac{3-s}{1+r}, \quad (21.18)$$

is required for convection to happen somewhere. Table 21.1 illustrates the behavior of  $d \ln T / d \ln P$  for various opacity sources. The fact that the  $H^-$  opacity increases with temperature forces the temperature gradient to steepen with increasing pressure and ensures that low-mass stars have outer convective zones.

source	$r$	$s$	$\frac{3-s}{1+r}$
Thomson	0	0	3
free-free	1	-7/2	13/4
$H^-$	1/2	9	-4

Table 21.1: Right-hand side of eq. (21.18) for various opacities

### An irradiated atmosphere

Many extra-solar planets are in rather tight orbits and as a result are strongly irradiated. The following example is a simplified treatment<sup>1</sup>. At a distance  $D$  from the star, the luminous flux is  $\sigma_{\text{SB}} T_{\star}^4 (R_{\star}/D)^2$ . The incident intensity is then  $(\sigma_{\text{SB}}/\pi) W T_{\star}^4$ , where  $W = (R_{\star}/D)^2$ ,

<sup>1</sup> D. G. Hummer. The effect of reflected and external radiation on stellar flux distributions. *ApJ*, 257:724–732, June 1982; and I. Hubeny, A. Burrows, and D. Sudarsky. A Possible Bifurcation in Atmospheres of Strongly Irradiated Stars and Planets. *ApJ*, 594:1011–1018, September 2003

since this will give the flux when integrated over all forward directions.

An classic approximation in stellar atmospheres is to write the intensity as a sum of two streams,

$$I_\nu(\mu) = I_\nu^+ \delta\left(\mu - \frac{1}{\sqrt{3}}\right) + I_\nu^- \delta\left(\mu + \frac{1}{\sqrt{3}}\right). \quad (21.19)$$

The reason for the choice of  $\mu$  becomes apparent when we compute the mean intensity, the flux, and the pressure:

$$\begin{aligned} J_\nu &= \frac{1}{4\pi} \int d\phi d\mu I_\nu = \frac{1}{2} (I_\nu^+ + I_\nu^-) \\ F_\nu &= \int d\phi d\mu \mu I_\nu = \frac{2\pi}{\sqrt{3}} (I_\nu^+ - I_\nu^-) \\ P_\nu &= \frac{1}{c} \int d\phi d\mu \mu^2 I_\nu = \frac{2\pi}{3c} (I_\nu^+ + I_\nu^-). \end{aligned}$$

You will recognize by comparing  $P_\nu$  with  $J_\nu$  that this formalism automatically satisfies the Eddington approximation, since  $J_\nu = (c/4\pi)u_\nu$  (cf. § 7.1 and eq. [A.17]).

Following the standard method, we take successive moments of our equation of transfer (for a grey atmosphere),

$$\mu \frac{dI}{d\tau} = I - S,$$

to obtain

$$\frac{dF}{d\tau} = 4\pi(J - S) \quad (21.20)$$

$$c \frac{dP}{d\tau} = F. \quad (21.21)$$

In LTE,  $J - S = 0$  and therefore  $F = \text{const}$ . We therefore integrate eq. (21.21) and use the Eddington approximation,  $cP = (4\pi/3)J$ , to obtain

$$J(\tau) = \frac{3}{4\pi} F\tau + J_0. \quad (21.22)$$

To determine  $J_0$ , we use our two stream approximation to write  $J_0 = (\sqrt{3}/4\pi)F + I^-$ . Since  $F$  is constant, we set it to its value at great depth in the star. Let us characterize  $F$  by temperature  $T_{\text{int}}$  via  $F \equiv \sigma_{\text{SB}} T_{\text{int}}^4$ . Finally, we set  $I^-$  to the incident intensity,  $I^- = (\sigma_{\text{SB}}/\pi)WT_\star^4$  and note that in radiative equilibrium,  $J = B = (\sigma_{\text{SB}}/\pi)T^4(\tau)$ .

Collecting terms, we have the equation for the temperature structure,

$$T^4(\tau) = \frac{3}{4} T_{\text{int}}^4 \left( \tau + \frac{1}{\sqrt{3}} \right) + WT_\star^4. \quad (21.23)$$

If  $WT_\star \gg T_{\text{int}}$ , then the temperature is nearly isothermal to a depth  $\tau_h \approx W(T_\star/T_{\text{int}})^4$ . The assumption of a grey atmosphere is, however,

quite poor: the incident photons are peaked in the optical, whereas the local temperatures are in the infrared. Even taking a mean opacity is not sufficient.

### 21.4 Line formation and the curve of growth

Spectral lines are the diagnostics of a stellar atmosphere's temperature, pressure, and composition. We'll briefly treat here how the ambient conditions set the line shape.

#### The classical oscillator

Suppose we have a classical charged harmonic oscillator. The instantaneous power emitted by the oscillator is

$$P(t) = \frac{2e^2}{3c^3} |\dot{\mathbf{u}}|^2, \quad (21.24)$$

and when averaged over a cycle is

$$\langle P(t) \rangle = \frac{e^2}{3c^3} x_0^2 \omega^4, \quad (21.25)$$

since  $\dot{\mathbf{u}} = -\omega^2 \mathbf{x}_0 \cos \omega t$ . Since the oscillator is radiating, it is losing energy and is damped. Let us write the damping as  $\mathbf{F}_{\text{rad}} \cdot \mathbf{u}$  and integrate over a cycle,

$$-\int_{t_1}^{t_2} dt \frac{2e^2}{3c^3} \dot{\mathbf{u}} \cdot \dot{\mathbf{u}} = -\frac{2e^2}{3c^3} \dot{\mathbf{u}} \cdot \mathbf{u} \Big|_{t_1}^{t_2} + \frac{2e^2}{3c^3} \int_{t_1}^{t_2} dt \dot{\mathbf{u}} \cdot \mathbf{u}.$$

The first term vanishes and we can therefore identify

$$\mathbf{F}_{\text{rad}} = \frac{2e^2}{3c^3} \dot{\mathbf{u}} = -m \left( \frac{2e^2 \omega^2}{3c^3 m} \right) \mathbf{u}$$

as the radiation damping term with the term in parenthesis being the damping constant  $\gamma$ . If there is an driving electric field on our oscillator, then its equation of motion becomes

$$m\ddot{\mathbf{x}} = -m\omega_0^2 \mathbf{x} + eEe^{i\omega t} - m\gamma\dot{\mathbf{x}}. \quad (21.26)$$

Using a trial function  $\mathbf{x} \propto e^{i\omega t}$  gives

$$\mathbf{x} = \frac{e}{m} \frac{Ee^{i\omega t}}{(\omega_0^2 - \omega^2) + i\omega\gamma}.$$

Taking the second derivative w.r.t. time of  $\mathbf{x}$ , substituting into eq. (21.24), and averaging over a cycle gives the power radiated by the oscillator,

$$\langle P(t) \rangle = \frac{e^4 \omega^4 E^2}{3c^2 m^2} \frac{1}{(\omega_0^2 - \omega^2)^2 + \gamma^2 \omega^2}.$$

Dividing  $\langle P(t) \rangle$  by the incident power per unit area,  $cE^2/(8\pi)$ , gives the cross-section,

$$\sigma = \frac{8\pi}{3} \frac{e^4}{m^2 c^3} \frac{\omega^4}{(\omega_0^2 - \omega^2)^2 + \gamma^2 \omega^2}. \quad (21.27)$$

Now, for  $\omega \approx \omega_0$ , we can expand  $(\omega_0^2 - \omega^2)^2 \approx 4\omega_0^2(\omega_0 - \omega)^2$ ; furthermore, we identify  $2e^2\omega_0^2/(3c^3m) = \gamma$  and equation (21.27) becomes

$$\sigma = \pi \left( \frac{e^2}{mc} \right) \frac{\gamma}{(\omega_0 - \omega)^2 + (\gamma/2)^2}. \quad (21.28)$$

The line profile is Lorentzian, with a width  $\gamma$ . In terms of wavelength, the width is

$$\Delta\lambda = \left| \frac{d\lambda}{d\omega} \right| \gamma = \frac{2\pi c}{\omega^2} \gamma = 1.2 \times 10^{-4} \text{ Å}.$$

This width is independent of the transition frequency (it is just the classical electron radius), and it is very, very small. In a stellar atmosphere, the width is set by interactions and doppler broadening.

Suppose we model the oscillator as being started and stopped by impacts; in between impacts it just goes as  $e^{i\omega_0 t}$ . To get the spectrum, we take the Fourier transform,

$$F(\omega, t) = \int_0^t dt' \exp[i(\omega_0 - \omega)t'],$$

where  $t$  is some time between impacts. Now if the impacts are distributed randomly and are uncorrelated, then the distribution of wait times follows a Poisson distribution,

$$W(t) dt = e^{-t/\tau} dt/\tau,$$

where  $\tau$  is the average time between collisions. Using this to compute the energy spectrum, we obtain

$$E(\omega) = \frac{1}{2\pi\tau} \int_0^\infty dt F(\omega, t) F^*(\omega, t) W(t) = \frac{1}{\pi\tau} \frac{1}{(\omega_0 - \omega)^2 + (1/\tau)^2};$$

the line profile is again Lorentzian, with a FWHM  $2/\tau$ .

We might be inclined to treat the atoms as hard spheres, but this gives a large  $\tau$ , or equivalently a narrow line width. We are therefore led to consider longer-range interactions for setting the intrinsic line width. Table 21.2 lists such interactions. The picture is similar to our considerations of collisions in §8.1. For a given impact parameter, the interaction perturbs the energy levels; by integrating over a distribution of impact parameters one gets the intrinsic damping. Of course, we should really use a quantum mechanical calculation. We can scale

our cross-section to the classical result (eq. [21.28]), however, by writing

$$\sigma_\nu = \left( \frac{\pi e^2}{m_e c} \right) f \phi_\nu, \quad (21.29)$$

where  $\phi_\nu$  is the line profile (dimension  $\sim \text{Hz}^{-1}$ ) and  $f$  is a dimensionless cross-section called the **oscillator strength**.

perturbation	form	source	affects	Table 21.2: Interactions in stellar atmospheres
linear Stark	$C_2 r^{-2}$	$e^-$ , $p$ , ions	$H$ ( $H\alpha$ , $H\beta$ , ...)	
quadratic Stark	$C_4 r^{-4}$	$e^-$	non-hydrogenic ions	
van der Waals	$C_6 r^{-6}$	atoms, $H$	most atomic lines, esp. in cool stars	

### The Curve of Growth

A classical technique in the analysis of stellar spectra is to construct the *curve of growth*, which relates the equivalent width of a line  $W_\nu$  to the opacity in the line. This discussion follows Mihalas, *Stellar Atmospheres*.

Let's first get the opacity in the line. Write the cross-section for the transition  $i \rightarrow j$  as

$$\sigma_\nu = \left( \frac{\pi e^2}{m_e c} \right) f_{ij} \phi_\nu,$$

where the first term is the classical oscillator cross-section,  $f_{ij}$  is the oscillator strength and contains the quantum mechanical details of the interaction, and  $\phi_\nu$  is the line profile. Now recall that the opacity is given by  $\kappa_\nu = n_i \sigma_\nu / \rho$ , where  $n_i$  denotes the number density of available atoms in state  $i$  available to absorb a photon. Furthermore, we need to allow for *stimulated emission* from state  $j$  to state  $i$ . With this added, the opacity is (I'm writing it as  $\chi_\nu$  to distinguish it from the *continuum opacity*)

$$\rho \chi_\nu = \left( \frac{\pi e^2}{m_e c} \right) f_{ij} \phi_\nu n_i \left[ 1 - \frac{g_i n_j}{g_j n_i} \right]. \quad (21.30)$$

If we are in LTE, then the relative population of  $n_i$  and  $n_j$  follow a Boltzmann distribution,

$$1 - \frac{g_i n_j}{g_j n_i} = 1 - \exp \left( - \frac{h\nu}{kT} \right).$$

This ensures we have a positive opacity. If our population were inverted, i. e., more atoms in the upper state  $j$ , then the opacity would be negative and we would have a *laser*.

Now for the line profile. In addition to damping, there is also Doppler broadening from thermal (or convective) motion. Let the line

profile (here we'll switch to  $\nu$ , rather than  $\omega$ ) be Lorentzian,

$$\phi = \frac{\Gamma/(4\pi)}{(\nu - \nu_0)^2 + (\Gamma/[4\pi])^2}.$$

In a Maxwellian distribution, the probability of having a line-of-sight velocity in  $(u, u + du)$  is

$$\mathcal{P}(u) du = \frac{1}{\sqrt{\pi}u_0} \exp\left(-\frac{u^2}{u_0^2}\right),$$

where  $u_0 = (2kT/m)^{1/2} = 12.85 \text{ km s}^{-1}$  ( $T/10^4 \text{ K}$ ) (for H) is the mean thermal velocity. The atom absorbs at a shifted frequency  $\nu(1 - u/c)$ , so the mean cross section is

$$\sigma_\nu = \int_{-\infty}^{\infty} \sigma\left(1 - \frac{u}{c}\right) \mathcal{P}(u) du. \quad (21.31)$$

After some algebraic manipulations, we have the cross-section

$$\begin{aligned} \sigma_\nu &= \left(\frac{\sqrt{\pi}e^2}{m_e c}\right) f_{ij} \frac{1}{\Delta\nu_D} \left\{ \frac{a}{\pi} \int_{-\infty}^{\infty} \frac{\exp(-y^2) dy}{(v - y)^2 + a^2} \right\} \\ &\equiv \frac{1}{\Delta\nu_D} H(a, v) \end{aligned} \quad (21.32)$$

where  $\Delta\nu_D \equiv \nu u_0 / c$  is the doppler width,  $a = \Gamma/(4\pi\Delta\nu_D)$  is the ratio of the damping width  $\Gamma$  to the doppler width, and  $v = \Delta\nu/\Delta\nu_D$  is the difference in frequency from the line center in units of the doppler width. The function  $H(a, v)$  is called the *Voigt* function.

Let's combine the line opacity with the continuum opacity and solve the equation of transfer. For simplicity, we are going to assume pure absorption in both the continuum and the line. Under these conditions, the source function is (see the notes on the Eddington atmosphere)  $S_\nu = B_\nu$ , the Planck function. For a plane-parallel atmosphere, the equation of transfer is then

$$\mu \frac{dI_\nu}{d\tau_\nu} = I_\nu - B_\nu \quad (21.33)$$

where  $\mu$  is the cosine of the angle of the ray with vertical. Solving equation (21.33) for the emergent intensity at  $\tau_\nu = 0$  gives

$$I_\nu(\mu) = \frac{1}{\mu} \int_0^\infty B_\nu[T(\tau_\nu)] \exp(-\tau_\nu/\mu) d\tau_\nu. \quad (21.34)$$

The opacity is given by

$$\kappa_\nu = \kappa_\nu^C + \chi_\nu, \quad (21.35)$$

where  $\kappa_\nu^C$  is the continuum opacity and  $\chi_\nu = \chi_0 \phi_\nu$  is the line opacity, with

$$\chi_0 = \frac{1}{\rho} \left(\frac{\pi e^2}{m_e c}\right) f_{ij} n_i \left(1 - e^{h\nu_\ell/kT}\right)$$

being the line opacity at the line center  $\nu_\ell$ .

As a further simplification, we can usually ignore the variation with  $\nu$  in  $\kappa_\nu^C$  over the width of the line. As a more suspect approximation (although it is not so bad in practice), let's assume that  $\beta_\nu \equiv \chi_\nu/\kappa_C$  is independent of  $\tau_\nu$ . With this assumption we can write  $d\tau_\nu = (1 + \beta_\nu)d\tau$ , where  $\tau = -\rho\kappa^C dz$ . Finally, let's assume that in the line forming region, the temperature does not vary too much, so that we can expand  $B_\nu$  to first order in  $\tau$ ,

$$B_\nu[T(\tau)] \approx B_0 + B_1\tau,$$

where  $B_0$  and  $B_1$  are constants. Inserting these approximations into equation (21.34), multiplying by the direction cosine  $\mu$  and integrating over outward bound rays gives us the flux,

$$\begin{aligned} F_\nu &= 2\pi \int_0^1 \int_0^\infty [B_0 + B_1\tau] \exp\left[-\frac{\tau}{\mu}(1 + \beta_\nu)\right] (1 + \beta_\nu) d\tau d\mu \\ &= \pi \left[ B_0 + \frac{2}{3} \frac{B_1}{1 + \beta_\nu} \right]. \end{aligned} \quad (21.36)$$

Far from the line-center,  $\beta_\nu \rightarrow 0$ , implying that the continuum flux is

$$F_\nu^C = \pi \left[ B_0 + \frac{2B_1}{3} \right].$$

Hence the depth of the line is

$$A_\nu \equiv 1 - \frac{F_\nu}{F_\nu^C} = A_0 \frac{\beta_\nu}{1 + \beta_\nu}, \quad (21.37)$$

where

$$A_0 \equiv \frac{2B_1/3}{B_0 + 2B_1/3} \quad (21.38)$$

is the depth of an infinitely opaque ( $\beta_\nu \rightarrow \infty$ ) line.

**EXERCISE 21.2** — Explain why an infinitely opaque line ( $A_0 = 0$  in eq. [21.38]) is not completely black.

Now that we have the depth of the line  $A_\nu$  we can compute the *equivalent width*,

$$W_\nu \equiv \int_0^\infty A_\nu d\nu = A_0 \int_0^\infty \frac{\beta_\nu}{1 + \beta_\nu} d\nu. \quad (21.39)$$

Let's change variables from  $\nu$  to  $v = \Delta\nu/\Delta\nu_D = (\nu - \nu_\ell)/\Delta\nu_D$ . Since  $H(a, v)$  is symmetrical about the line center, we will just integrate over  $\Delta\nu > 0$ , giving

$$W_\nu = 2A_0\Delta\nu_D \int_0^\infty \frac{\beta_0 H(a, v)}{1 + \beta_0 H(a, v)} dv, \quad (21.40)$$

with  $\beta_0 = \chi_0 / (\kappa^C \Delta\nu_D)$ .

It's useful to understand the behavior of  $W_\nu$  in various limits.

First, at small line optical depth ( $\beta_0 \ll 1$ ) only the core of the line will be visible. In the core of the line,  $H(a, v) \approx \exp(-v^2)$  so we insert this into equation (21.40) and expand the denominator to give

$$\begin{aligned} W_\nu^* \equiv \frac{W_\nu}{2A_0\Delta\nu_D} &= \int_0^\infty \sum_{k=1}^\infty (-1)^{k-1} \beta_0^k e^{-kv^2} dv \\ &= \frac{1}{2} \sqrt{\pi} \beta_0 \left[ 1 - \frac{\beta_0}{\sqrt{2}} + \frac{\beta_0^2}{\sqrt{3}} - \dots \right]. \end{aligned} \quad (21.41)$$

Here  $W_\nu^*$  is the *reduced equivalent width*. Notice that since  $\beta_0 \propto 1/\Delta\nu_D$  (cf. eq. [21.32]), the equivalent width  $W_\nu$  is independent of  $\Delta\nu_D$  in this *linear regime*. Physically, in the limit of small optical depth, each atom in state  $i$  is able to absorb photons, and the flux removed is just proportional to the number of atoms  $n_i$ .

As we increase  $\beta_0$  eventually the core of the line saturates—no more absorption in the core is possible. As a result, the equivalent width should be nearly constant until there are so many absorbers that the damping wings contribute to the removal of flux. In the *saturation regime*, the Voigt function is still given by  $e^{-v^2}$ , but we can no longer assume  $\beta_0 \ll 1$ , so our expansion in equation (21.41) won't work. Let's go back to our integral, eq. (21.40), change variables to  $z = v^2$ , and define  $\alpha = \ln \beta_0$  to find

$$W_\nu^* = \frac{1}{2} \int_0^\infty \frac{z^{-1/2}}{e^{z-\alpha} + 1} dz.$$

This may not look like an improvement, but you might notice that it bears a resemblance to a Fermi-Dirac integral (see the notes on the equation of state). That means that very smart people figured out tricks to handle these integrals and all we have to do is look up what they did. In this case we have Sommerfeld to thank. In this saturation regime,

$$W_\nu^* \approx \sqrt{\ln \beta_0} \left[ 1 - \frac{\pi^2}{24(\ln \beta_0)^2} - \frac{7\pi^4}{384(\ln \beta_0)^4} - \dots \right]. \quad (21.42)$$

Note that the amount of flux removed is basically  $2A_0\Delta\nu_D$ : the line is maximally dark across the gaussian core.

Finally, if we continue to increase the line opacity, there will finally be so many absorbers that there will be significant flux removed from the wings. Now the form of the Voigt profile is  $H(a, v) \approx (a/\sqrt{\pi})v^{-2}$ , so our integral (eq. [21.40]) in this *damping regime* be-

comes

$$\begin{aligned} W_v^* &= \int_0^\infty \left(1 + \frac{\sqrt{\pi}v^2}{\beta_0 a}\right)^{-1} dv \\ &= \frac{1}{2} (\pi a \beta_0)^{1/2}. \end{aligned} \quad (21.43)$$

Note that since  $a\beta_0 \propto \Delta\nu_D^{-2}$ ,  $W_v$  is again independent of the doppler width in this regime.

Now that we have this “curve of growth”,  $W_v^*(\beta_0)$ , why is it useful? Since it only involves the equivalent width, it is possible to construct the curve of growth empirically without a high-resolution spectrum. Next, let’s put some of the factors back into the quantities in the curve of growth. First, for a set of lines, the population of the excited state depends on the Boltzmann factor  $\exp(-E/kT)$ . Second, we can expand out the Doppler width in both  $W_\lambda^*$  and  $\beta_0$ ,

$$\log \left( \frac{W_\lambda}{\Delta\lambda_D} \right) = \log \left( \frac{W_\lambda}{\lambda} \right) - \log \left( \frac{u_0}{c} \right) \quad (21.44)$$

$$\log \beta_0 = \log(g_i f_{ij} \lambda) - \frac{E}{kT} + \log(N/\kappa^C) + \log C \quad (21.45)$$

where  $C$  contains all of the constants and the continuum opacity. The temperature  $T$  is picked as a free parameter, and is picked to minimize scatter about a single curve that is assumed to fit all of the lines. What is measured then is  $\log(W_\lambda/\lambda)$  and  $\log(g_i f_{ij} \lambda)$ ; by comparing them to theoretical curves one gets an estimate of  $\log(u_0/c)$ , the mean velocity of atoms (may be thermal or turbulent). Since the continuum opacity  $\kappa^C$  usually depends on the density of H, one gets from equation (21.45) an estimate of the abundance of the line-producing element to H.

EXERCISE 21.3— There is a subtlety involved when an atmospheric opacity is scattering-dominated, because scattering does not change the photon energy. Suppose we have an atmosphere where the Thomson scattering dominates the opacity, and the absorption of a photon is inverse bremsstrahlung (free-free), for which you can get the cross section expression from equation (7.26). Note that we do not want the Rosseland mean here, we want to know what happens to a photon of a specific frequency. Finally, we are after scalings here, so don't get hung up on the precise value of numerical prefactors.

1. Is the opacity scattering-dominated at all frequencies?
  2. Trace a photon of frequency  $\nu$  back into the atmosphere. How deep (in terms of the scattering optical depth) does it go before being absorbed? Is there a single well-defined photosphere for all frequencies? (*Hint: the photon is taking a random walk into the star.*)
  3. Now, suppose the emergent intensity is still Planckian, but with a temperature that is the local temperature at the depth where the photon was last absorbed. Obtain an expression for  $T$  and  $\rho$  as a function of scattering optical depth, and use this to derive an approximate expression for the spectrum at high frequencies. How does it compare to a blackbody at temperature  $T_{\text{eff}}$ ? *Hint:* you may find the article by Illarionov and Sunyaev (Astrophys. & Space Science 19: 61 [1972]) helpful.
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# A

## Technical Notes

### A.1 Thermodynamical derivatives

A common task in stellar physics is transforming between different derivatives with respect to different thermodynamical quantities. For example, you may have expressions for  $(\partial\kappa/\partial T)_\rho$  and  $(\partial\kappa/\partial\rho)_T$ , but you need  $(\partial\kappa/\partial T)_P$  and  $(\partial\kappa/\partial P)_T$ . There is a straightforward way to handle transforming from  $(\rho, T)$  space to  $(P, T)$  space, and that is using Jacobians. Despite the utility of this technique, it is not commonly discussed in astrophysical texts<sup>1</sup>.

The *Jacobian* is defined as the determinant of a matrix of partial derivatives,

$$\begin{aligned} \frac{\partial(a,b)}{\partial(c,d)} &\equiv \det \begin{bmatrix} \left(\frac{\partial a}{\partial c}\right)_d & \left(\frac{\partial a}{\partial d}\right)_c \\ \left(\frac{\partial b}{\partial c}\right)_d & \left(\frac{\partial b}{\partial d}\right)_c \end{bmatrix} \\ &= \left(\frac{\partial a}{\partial c}\right)_d \left(\frac{\partial b}{\partial d}\right)_c - \left(\frac{\partial a}{\partial d}\right)_c \left(\frac{\partial b}{\partial c}\right)_d. \end{aligned} \quad (\text{A.1})$$

Because interchanging any the rows (or the columns) causes the determinant to change sign,

$$\frac{\partial(b,a)}{\partial(c,d)} = -\frac{\partial(a,b)}{\partial(c,d)} \quad (\text{A.2})$$

and

$$\frac{\partial(a,b)}{\partial(d,c)} = -\frac{\partial(a,b)}{\partial(c,d)}. \quad (\text{A.3})$$

Further,

$$\frac{\partial(a,s)}{\partial(c,s)} = \left(\frac{\partial a}{\partial c}\right)_s \left(\frac{\partial s}{\partial s}\right)_a - \left(\frac{\partial a}{\partial s}\right)_c \left(\frac{\partial s}{\partial c}\right)_s = \left(\frac{\partial a}{\partial c}\right)_s, \quad (\text{A.4})$$

and

$$\frac{\partial(a,b)}{\partial(a,b)} = \left(\frac{\partial a}{\partial a}\right)_b \left(\frac{\partial b}{\partial b}\right)_a - \left(\frac{\partial a}{\partial b}\right)_a \left(\frac{\partial b}{\partial a}\right)_b = 1. \quad (\text{A.5})$$

<sup>1</sup> L. D. Landau and E. M. Lifshitz. *Statistical Physics, part 1*. Pergamon Press, Oxford, 3 edition, 1980

Hence we can write thermodynamical derivative in terms of Jacobians, for example,

$$\left(\frac{\partial T}{\partial P}\right)_S = \frac{\partial(T, S)}{\partial(P, S)}. \quad (\text{A.6})$$

Finally, when multiplying two Jacobians, one can “cancel” identical upper and lower parts,

$$\frac{\partial(a, b)}{\partial(c, d)} \frac{\partial(c, d)}{\partial(s, t)} = \frac{\partial(a, b)}{\partial(s, t)}, \quad (\text{A.7})$$

as can be readily checked by expanding out both the left and right hand sides.

### *Common thermodynamic derivatives*

Certain thermodynamical derivatives occur commonly in working with fluids. The first set express how the pressure relates to the pair  $\rho, T$ :

$$\chi_T \equiv \frac{T}{P} \left( \frac{\partial P}{\partial T} \right)_\rho, \quad \chi_\rho \equiv \frac{\rho}{P} \left( \frac{\partial P}{\partial \rho} \right)_T. \quad (\text{A.8})$$

For a fixed composition the equation of state can always be expanded about a point  $P_0(\rho_0, T_0)$  as

$$P = P_0 \left( \frac{\rho}{\rho_0} \right)^{\chi_\rho} \left( \frac{T}{T_0} \right)^{\chi_T}, \quad (\text{A.9})$$

or equivalently

$$\frac{dP}{P} = \chi_\rho \frac{d\rho}{\rho} + \chi_T \frac{dT}{T}.$$

Here we are implicitly keeping the composition fixed.

The next set concern how quantities transform under adiabatic changes. For a general equation of state with fixed composition, define the following:

$$\Gamma_1 \equiv \frac{\rho}{P} \left( \frac{\partial P}{\partial \rho} \right)_S = \left( \frac{\partial \ln P}{\partial \ln \rho} \right)_S; \quad (\text{A.10})$$

$$\frac{\Gamma_2 - 1}{\Gamma_2} \equiv \left( \frac{\partial \ln T}{\partial \ln P} \right)_S \equiv \nabla_{\text{ad}}; \quad (\text{A.11})$$

$$\Gamma_3 - 1 \equiv \left( \frac{\partial \ln T}{\partial \ln \rho} \right)_S. \quad (\text{A.12})$$

The nomenclature is historical. These quantities are not independent: for example, one can show (see exercise 2) that

$$\Gamma_1 = [\chi_\rho + \chi_T (\Gamma_3 - 1)].$$

Furthermore,

$$\begin{aligned} \Gamma_3 - 1 &= \frac{P}{\rho T} \frac{\chi_T}{c_\rho}, \\ \frac{\Gamma_2 - 1}{\Gamma_2} &= \frac{\Gamma_3 - 1}{\Gamma_1}. \end{aligned}$$

Furthermore, one can show that the specific heat at constant pressure is  $c_P = (\Gamma_1/\chi_\rho)c_\rho$ .

### *A worked example: derivatives of the opacity*

Here's a simple worked example of how we can use these identities.

Suppose we have an opacity  $\kappa(\rho, T)$  that is expressed in terms of density and temperature, but we need the quantity  $(\partial\kappa/\partial T)_P$ . We can express  $(\partial\kappa/\partial T)_P$  as

$$\begin{aligned} \left(\frac{\partial\kappa}{\partial T}\right)_P &= \frac{\partial(\kappa, P)}{\partial(T, P)} \\ &= \frac{\partial(\kappa, P)}{\partial(T, \rho)} \frac{\partial(T, \rho)}{\partial(T, P)} \\ &= \left(\frac{\partial\rho}{\partial P}\right)_T \left[ \left(\frac{\partial\kappa}{\partial T}\right)_\rho \left(\frac{\partial P}{\partial\rho}\right)_T - \left(\frac{\partial\kappa}{\partial\rho}\right)_T \left(\frac{\partial P}{\partial T}\right)_\rho \right] \\ &= \left(\frac{\partial\kappa}{\partial T}\right)_\rho - \left(\frac{\partial\kappa}{\partial\rho}\right)_T \frac{\partial(\rho, T)}{\partial(P, T)} \frac{\partial(P, \rho)}{\partial(T, \rho)} \\ &= \left(\frac{\partial\kappa}{\partial T}\right)_\rho - \left(\frac{\partial\kappa}{\partial\rho}\right)_T \frac{\chi_T}{\chi_\rho} \frac{\rho}{T}. \end{aligned} \quad (\text{A.13})$$

We have to add to  $(\partial\kappa/\partial T)_\rho$  a term that allows for the density to change with temperature at constant pressure.

#### EXERCISE A.1—

1. Show that

$$\left(\frac{\partial T}{\partial P}\right)_S \left(\frac{\partial S}{\partial T}\right)_P \left(\frac{\partial P}{\partial S}\right)_T = -1$$

2. Show that

$$\left(\frac{\partial P}{\partial s}\right)_\rho = \frac{P}{c_\rho} \chi_T, \quad \left(\frac{\partial P}{\partial \rho}\right)_s = \frac{P}{\rho} [\chi_\rho + \chi_T (\Gamma_3 - 1)].$$

### *A.2 Moments of the radiant intensity*

In the notes on radiative transport (§ 7.1) we have tended to use quantities derived from the specific intensity with a readily interpretable physical meaning, such as the energy density, flux, and radiation pressure. Often, however, it is useful to make this more formal by defining *moments* of the specific intensity, which are just weighted angular averages. For example, integrating  $I_\nu$  over all angles and dividing by  $4\pi$  gives

$$J_\nu \equiv \frac{1}{4\pi} \int_0^{2\pi} d\phi \int_0^\pi \sin\theta d\theta I_\mu = \frac{1}{2} \int_{-1}^1 d\mu I_\nu. \quad (\text{A.14})$$

Here  $\mu = \cos \theta$ . For the first moment, we can multiply  $I_\nu$  by a unit vector  $\mathbf{k}$ , and then dot that into the unit directional vector and integrate over all directions,

$$H_\nu \equiv \frac{1}{4\pi} \int_0^{2\pi} d\phi \int_0^\pi \sin \theta d\theta I_\mu \mathbf{k} \cdot \mathbf{n} = \frac{1}{2} \int_{-1}^1 d\mu \mu I_\nu. \quad (\text{A.15})$$

Finally, we can multiply  $I_\nu$  by a tensor  $\mathbf{k}\mathbf{k}$ ; contracting this along  $\mathbf{n}$  gives

$$K_\nu \equiv \frac{1}{4\pi} \int_0^{2\pi} d\phi \int_0^\pi \sin \theta d\theta I_\mu (\mathbf{k} \cdot \mathbf{n})^2 = \frac{1}{2} \int_{-1}^1 d\mu \mu^2 I_\nu. \quad (\text{A.16})$$

If we further integrate equations (A.14)–(A.16) over all frequencies, we will obtain expressions for the energy density, flux, and radiation pressure,

$$u = \frac{4\pi}{c} J, \quad F = 4\pi H, \quad P = \frac{4\pi}{c} K. \quad (\text{A.17})$$

### A.3 Shocks

In a shock, the properties of the fluid change over a scale of a few mean free paths. Over this distance, the fluid properties—density, pressure, temperature—are not well characterized by smooth, differentiable functions. To understand how shocks arise, imagine a long tube filled with a gas and fitted with a piston at one end. The piston is pushed into the gas; this creates a compressed region of higher density and pressure, and a compression wave propagates into the cylinder. This behavior can be captured by Burgers's equation,

$$\frac{\partial u}{\partial t} + \frac{\partial}{\partial x} \left( \frac{u^2}{2} \right) = 0,$$

which is simple but has a nonlinear term that captures the formation of a shock.

Information about the disturbance is carried by acoustic waves. In the compressed region, however, the sound speed is higher. As a result the back of the transition region travels at a faster velocity than the front of the disturbance, so that over time the disturbance steepens (Fig. A.1). This steepening continues until the thickness of the transition is small enough that diffusive effects—i.e., viscosity—can balance the steepening. This gives a characteristic width to the front. On macroscopic scales, the transition is essentially a discontinuity and is termed a *shock front*.

To illustrate the properties of a shock, we'll follow Zel'dovich and Kompaneets<sup>2</sup> and imagine a piston being pushed into a tube filled with gas, as illustrated in Fig. A.2. The piston moves to the right with velocity  $u$ . We assume that the piston has been pushing the fluid for

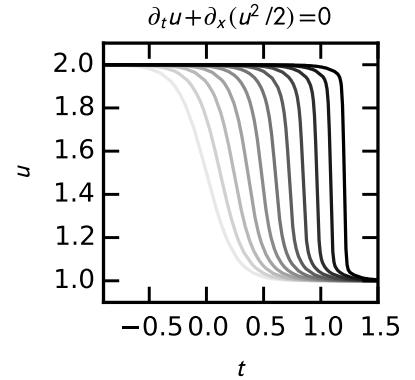


Figure A.1: The plots show the solution of Burgers's equation for a disturbance propagating along the  $x$ -direction. Because the sound speed is greater in the compressed region, the “back” of the disturbance  $B$  moves faster than the front  $F$ : as a result, the disturbance steepens. The disturbance steepens into a shock in a time  $t \approx \Delta / (u_L - u_R)$ , where  $\Delta$  is the initial width of the transition and  $u_L$  and  $u_R$  are the values of  $u$  on the left and right sides of the disturbance.

<sup>2</sup> Ya. B. Zel'dovich and A. S. Kompaneets. *Theory of Detonation*. Academic Press, Inc., 1960

a while at constant velocity, so that any transients have died down, leaving a simple flow structure at  $t = 0$  (Fig. A.2, top panel): far from the piston, the fluid is at rest, with density  $\rho_0$ , pressure  $P_0$ , and internal energy  $\varepsilon_0$ ; a shock propagates into this fluid with velocity  $S$ ; between the shock and the piston the fluid moves with the same velocity,  $u$ , as the piston and has density  $\rho_1$ , pressure  $P_1$ , and  $\varepsilon_1$ .

In a time  $t$ , the piston has moved (Fig. A.2, bottom panel) a distance  $u \cdot t$ , and the shock has moved a distance  $S \cdot t$ ; the mass of the “shocked” fluid has therefore increased by  $\rho_1(S - u)t$ . This must equal the mass swept up by the shock, namely  $\rho_0 St$ . We therefore have our first relation,

$$\rho_1(S - u) = \rho_0 S. \quad (\text{A.18})$$

The fluid swept up by the shock now has velocity  $u$  and thus a momentum  $(\rho_0 St)u$ . This increase in momentum must equal the net impulse imparted to the fluid,  $(P_1 - P_0)t$ . We therefore have our second relation,

$$S_1 - S_0 = S\rho_0 u. \quad (\text{A.19})$$

The fluid swept up by the shock has a change in total energy,  $(\rho_0 St)(\varepsilon_1 + u^2/2 - \varepsilon_0)$ . This must equal the work done on the fluid by the piston, which is the force times displacement  $P_1 ut$ . We therefore have our third relation,

$$\rho_0 S \left[ \varepsilon_1 - \varepsilon_0 + \frac{1}{2} u^2 \right] = P_1 u. \quad (\text{A.20})$$

To abstract the problem, we transform to a frame moving with the shock. In this frame the upstream velocity ahead of the shock has velocity  $u_0 = -S$ ; the downstream velocity is  $u_1 = u - S$ . The difference in velocity is  $|u_0 - u_1| = u$ . In this frame, equations (A.18), (A.19), and (A.20) take on the more familiar forms,

$$\rho_1 u_1 = \rho_0 u_0, \quad (\text{A.21})$$

$$P_1 + \rho_1 u_1^2 = P_0 + \rho_0 u_0^2, \quad (\text{A.22})$$

$$(\rho_1 u_1) \left[ \varepsilon_1 + \frac{1}{2} u_1^2 + \frac{P_1}{\rho_1} \right] = (\rho_0 u_0) \left[ \varepsilon_0 + \frac{1}{2} u_0^2 + \frac{P_0}{\rho_0} \right]. \quad (\text{A.23})$$

Recognize these? Compare with equations (3.1), (3.3), and (3.5), and recognize that equations (A.21), (A.22), (A.23) just express that the downstream mass flux, momentum flux, and energy flux equal their upstream counterparts.

Equations (A.21), (A.22), (A.23), when supplemented with an equation of state, are sufficient to allow solution for  $P_1$  and  $\rho_1$  in terms of  $P_0$ ,  $\rho_0$  and  $u$ . To explore the properties of the shock, we adopt a simple ideal adiabatic equation of state:  $P = (\gamma - 1)\rho\varepsilon = K\rho^\gamma$

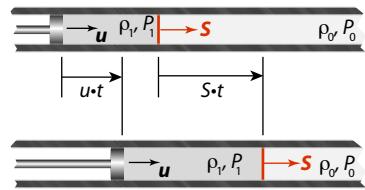


Figure A.2: Schematic of a piston driving a shock. In this schematic, the shock propagates at velocity  $S$ .

with sound speed  $c^2 = \gamma P/\rho$ . The enthalpy is then

$$\varepsilon + \frac{P}{\rho} = \frac{\gamma}{\gamma - 1} \frac{P}{\rho}.$$

Using this equation of state, we can reduce equations (A.21), (A.22), (A.23) into expressions in terms of the pressure ratio  $P = P_1/P_0$ :

$$\text{density ratio, } \frac{\rho_1}{\rho_0} = \frac{(\gamma + 1)P + (\gamma - 1)}{(\gamma - 1)P + (\gamma + 1)}; \quad (\text{A.24})$$

$$\text{entropy increase, } s_1 - s_0 = c_v \ln \left[ \frac{P_1}{P_0} \left( \frac{\rho_0}{\rho_1} \right)^\gamma \right]; \quad (\text{A.25})$$

$$\text{temperature ratio, } \frac{T_1}{T_0} = \frac{P_1 \rho_0}{P_0 \rho_1}; \quad (\text{A.26})$$

$$\text{upstream Mach, } \left( \frac{u_0}{c_0} \right)^2 = \frac{(\gamma + 1)P + (\gamma - 1)}{2\gamma}; \quad (\text{A.27})$$

$$\text{downstream Mach, } \left( \frac{u_1}{c_1} \right)^2 = \frac{(\gamma - 1)P + (\gamma + 1)}{2\gamma P}. \quad (\text{A.28})$$

In the limiting case of a strong shock,  $P \gg 1$ , the density increase is finite:  $\rho_1/\rho_0 \rightarrow (\gamma + 1)/(\gamma - 1)$ ; for  $\gamma = 5/3$ ,  $\rho_1/\rho_0 \rightarrow 4$ . The temperature and entropy jumps across the shock, however, scale with  $P$  and are arbitrarily large. The upstream Mach number is likewise proportional to the pressure ratio,  $\text{Ma}_0^2 \rightarrow (\gamma + 1)P/(2\gamma)$ , but the downstream flow is subsonic,  $\text{Ma}_1^2 \rightarrow (\gamma - 1)/(2\gamma)$ . The shock transforms the ordered (low-entropy) supersonic upstream flow into the disordered (high-entropy) subsonic downstream flow.

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