ENERGY TRANSPORT WITHIN A STAR

8.1 Introduction

Up to now, we have considered how energy is generated within the interior of stars by the processes of gravitational contraction and nuclear fusion. However, for a star to shine, the heat and photons generated deep in the interior of the star have to be transferred to the surface, where they are radiated away as electromagnetic radiation. This is the subject of this lecture.

Three different energy transport mechanisms can operate within the interior of a star. **Radiation** can transfer energy from the core to the surface, as photons are continuously absorbed and re-emitted as they interact with the plasma. **Conduction**, whereby heat is transferred at the microscopic level by collisions between particles. **Convection** is the bulk motion of cells, with hot, buoyant mass elements rising towards the surface, while cooler ones sink.

The efficiency of the first two processes depends on the mean free path of, respectively, photons and particles (i.e. electrons); since the photon mean free path is normally greater than that of electrons, radiative diffusion normally dominates over conduction, although the latter becomes dominant in special physical conditions, such as those prevailing in white dwarfs.

8.2 Radiative Transport

The diffusion of energy from the star's core to the surface via radiation is described by the Eddington equation for radiative equilibrium:

$$\boxed{\frac{dT}{dr} = -\frac{3}{4} \cdot \frac{1}{ac} \cdot \frac{\kappa\rho}{T^3} \cdot \frac{L_r}{4\pi r^2}} \tag{8.1}$$

where a is the radiation constant, $a = 4\sigma/c$, σ is Stefan-Botltzmann constant, κ is the opacity and the other symbols have their usual meanings.

This equation can be derived in different ways; here we are going to use a simple physical argument.

Consider a small cell within the interior of a star, say with volume 1 cm^3 (see Figure 8.1). In LTE, the flux from below is $F_2 \sim \sigma T_2^4$ (blackbody emission). The flux from above is $F_1 \sim \sigma T_1^4$. Hence, the net flux through the element is $F \sim \sigma(T_2^4 - T_1^4)$. Generalising, we have:

$$F \sim -\frac{d}{dr}\sigma T^4 . \tag{8.2}$$

The flux through the volume element has to be multiplied by 'the transparency' of the layer, which is approximately (for small volume elements) $1/\kappa\rho$, that is the photon mean free path (as we saw in Lecture 5). Thus the total flux through the volume element is:

$$F \sim -\frac{1}{\kappa\rho} \cdot \frac{d}{dr}\sigma T^4 . \tag{8.3}$$

The third step in the derivation just equates the flux through a spherical surface of radius r to the total luminosity:

$$F = \frac{L_r}{4\pi r^2} \tag{8.4}$$

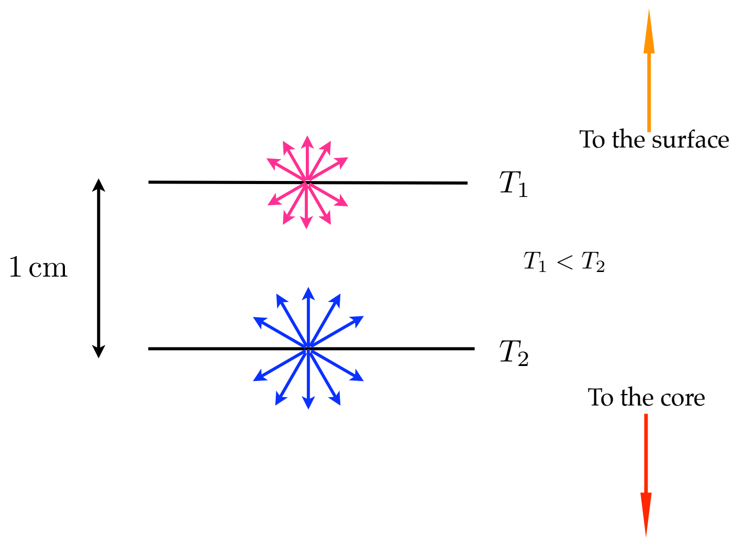

Figure 8.1: Blackbody emission at different distances r from the core of a star.

Combining these three equations we have:

$$L_r \sim -\frac{1}{\kappa\rho} \cdot 4\pi r^2 \cdot \frac{d}{dr}\sigma T^4 \;\; = \;\; -\frac{1}{\kappa\rho} \cdot 4\pi r^2 \cdot \frac{ac}{4} \cdot 4\,T^3 \cdot \frac{dT}{dr} \tag{8.5}$$

which can be rearranged in the form:

$$\frac{dT}{dr} \sim -\frac{1}{ac} \cdot \frac{\kappa\rho}{T^3} \cdot \frac{L_r}{4\pi r^2}\,. \tag{8.6}$$

Eqs. 8.1 and 8.6 differ only in the factor of 3/4 which comes from a proper integration over all angles.

In its 8.5 form (with a 4/3 factor on the right-hand side), the Eddington equation for radiative equilibrium gives the luminosity of star in terms of its radius, temperature, and temperature gradient when energy transport is primarily by radiative diffusion.

In its 8.1 form, the equation gives the temperature gradient required to carry the entire star's luminosity by radiation. A star, or a region within a star, in which this holds is said to be in radiative equilibrium, or simply *radiative*.

The Eddington equation is a valid description as long as the condition of LTE holds. This is clearly *not* the case at the stellar surface, or the photosphere, since this is the where the photons escape (and therefore the photon mean free path is no longer small compared to the distance over which dT/dr is small—recall our discussion at 5.2). Thus, near the stellar surface the diffusion approximation is no longer justified and one needs to solve the much more complicated equations of radiative transfer. Fortunately, the LTE and diffusion approximations are valid over almost the entire stellar interior.

Sometimes eq. 8.1 is written in terms of the mass. Recalling (Figure 7.1) that $dm = 4\pi r^2 \rho\, dr$, we have:

$$\frac{dT}{dm} = -\frac{3}{4} \cdot \frac{1}{ac} \cdot \frac{\kappa}{T^3} \cdot \frac{L_r}{(4\pi r^2)^2} \tag{8.7}$$

8.2.1 The Luminosity of the Sun

We can use the Eddington equation in the eq 8.5 form:

$$L_r = -\frac{4}{3}\frac{1}{\kappa\rho} \cdot 4\pi r^2 \cdot ac \cdot T^3 \cdot \frac{dT}{dr}$$

to estimate the luminosity of the Sun. Integrating the above equation and using the mean temperature of the Sun (as derived in lecture 7.1) as an approximation, we have:

$$L_\odot \simeq \frac{1}{3}\,\mu \cdot 4\pi r_\odot \cdot ac \cdot \langle T_\odot \rangle^4 \,.$$

$$L_\odot \simeq \frac{4}{3}{\cdot}0.1{\cdot}3\times 7{\times}10^{10}{\cdot}3{\times}10^{10}{\cdot}7.6{\times}10^{-15}\cdot\left(5\times10^6\right)^4 \quad \mathrm{cm^3\,s^{-1}\,erg\ cm^{-3}\ K^{-4}\,K^4}$$

where $\mu \sim 0.1\,\mathrm{cm}$.

$$L_\odot \simeq 6.4\times 10 \cdot 1\times 10^5 \cdot 625\times 10^{24} \qquad \mathrm{erg\ s^{-1}}$$

$$L_\odot \simeq 4\times 10^{33} \qquad \mathrm{erg\ s^{-1}}$$

For comparison, $L_\odot = 3.8\times 10^{33}$ erg s^{-1} – not bad!

8.3 Convection

Convection is a familiar phenomenon in our everyday lives: for example, our daily weather is caused by convection in the Earth's atmosphere. The surface of the Sun (Figure 8.2) is not smooth; instead we see bright granules separated by darker intergranular lanes. We know from Doppler velocity measurements that the motion of the bright regions is mostly outwards, while in the dark intergranular regions the gas is moving downwards. The motions and temperature inhomogeneities seen in the granulation pattern are attributed to the hydrogen convection zone just below the solar photosphere. The bulk motions of the gas and associated magnetic fields are thought to be the source of the mechanical energy flux that heats the solar chromosphere and corona.

Referring back to eq. 8.1, it can be readily appreciated that an increase in opacity κ in a stellar atmosphere will lead to a larger temperature gradient

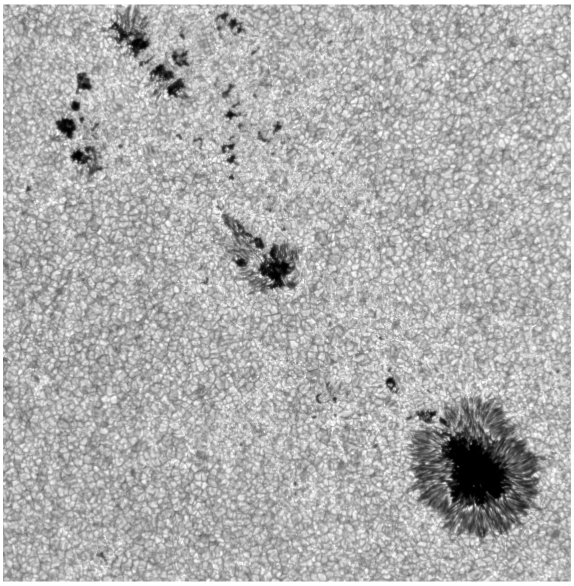

Figure 8.2: High resolution image of the solar photosphere, showing granulation and sunspots.

dT/dr (if the same luminosity is to be maintained). We also saw (Figure 5.5) that as the temperature decreases from $T \sim 10^7$ K, the opacity rises steeply with a Kramer's law $\langle \kappa \rangle \propto T^{-3.5}$. Thus, as we move from the core to the outer regions within a stellar interior, the temperature gradient is expected to become increasingly steep. A very steep temperature gradient is unstable, whether in a star or the Earth's atmosphere.

This can be appreciated by considering the consequences of displacing a volume element of gas, at equilibrium radius r inside a star—where T, P and ρ are the temperature, pressure and density, to a radius $r + dr$, where the ambient parameters are $T + dT$, $P + dP$ and $\rho + d\rho$ (Figure 8.3).

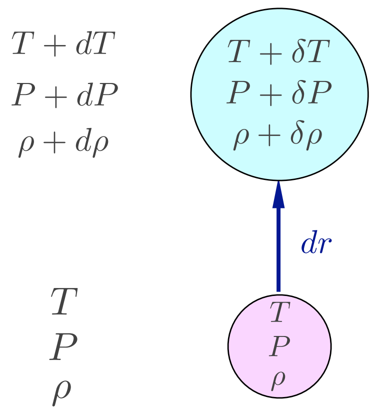

Figure 8.3: Illustration of the onset of convection. If $\rho + \delta\rho > \rho + d\rho$ the element will sink back to its former equilibrium position at radius r. But if $\rho + \delta\rho < \rho + d\rho$, the element will be buoyant and convection will ensue.

In the following treatment, we make two assumptions:

1. Pressure equilibrium: the element maintains the same pressure as its surroundings, and
2. The process is adiabatic, that is there no heat exchange between the volume element and its surroundings.

This is the same as saying that the timescale for removing pressure imbalance is short compared to the timescale for the establishment of thermal equilibrium. The ideal gas law:

$$P = \frac{\rho k T}{\mu m_H} \tag{8.8}$$

where μ is the mean molecular weight, m_H is the mass of the hydrogen atom, and the product μm_H is the mean mass of the gas particles, can be written as:

$$P = K\rho^{\gamma} \tag{8.9}$$

where K is a constant and γ is the ratio of the specific heats:

$$\gamma = \frac{C_P}{C_V} \qquad C_P \equiv \left.\frac{\partial Q}{\partial T}\right|_P, \quad C_V \equiv \left.\frac{\partial Q}{\partial T}\right|_V \tag{8.10}$$

which measure the amount of heat required to raise the temperature of a unit mass of material by a unit temperature interval at constant pressure and at constant volume respectively. Note that $C_P > C_V$ (and therefore $\gamma > 1$), because at constant pressure some of the energy input goes into increasing the volume of the gas and hence more energy is required to raise the temperature by 1 K.

What happens next depends on the difference $\Delta\rho \equiv (\rho + \delta\rho) - (\rho + d\rho)$. In the gravitational field g of the star, our rising volume element will experience a force $f = -g\Delta\rho$. Thus, if $\Delta\rho$ is +ve, the element will sink back to its equilibrium position at radius r but, if $\Delta\rho$ is −ve, the element will be buoyant and will rise further. This is the onset of convection.

Thus, the condition for convective instability is set by the density gradient within the star: if the gradient is less than that experienced by the volume element rising adiabatically, the star will be unstable against convection. Since we have assumed pressure equilibrium with its surrounding, we could also state the same instability criterion in terms of the temperature gradient (see eq. 8.8) as:

$$\left|\frac{dT}{dr}\right|_{\rm rad} > \left|\frac{dT}{dr}\right|_{\rm ad}. \tag{8.11}$$

Eq. 8.11 is known as the Schwarzschild criterion for convective instability (see Figure 8.4). It tells us that if the temperature profile within a star is super-adiabatic, the star is unstable against convection.

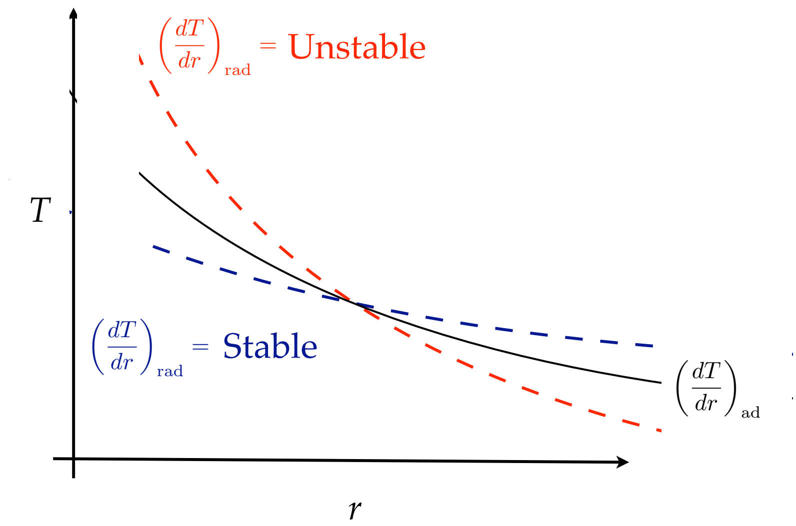

Figure 8.4: The Schwarzschild criterion for convective instability.

It is also instructive to express the convective instability criterion in terms of the parameter γ. We can do this as follows. From eq. 8.11 we have:

$$\left|\frac{d\ln T}{dr}\right|_{\rm rad} > \left|\frac{d\ln T}{dr}\right|_{\rm ad}, \tag{8.12}$$

diving through by $d\ln P/dr$,

$$\left|\frac{d\ln T}{d\ln P}\right|_{\rm rad} > \left|\frac{d\ln T}{d\ln P}\right|_{\rm ad} \tag{8.13}$$

or

$$\left|\frac{d\ln P}{d\ln T}\right|_{\rm rad} < \left|\frac{d\ln P}{d\ln T}\right|_{\rm ad}. \tag{8.14}$$

Eqs. 8.8 and 8.9 can be combined to give the adiabatic relation between pressure and temperature:

$$P\,T^{\frac{\gamma}{1-\gamma}} = K', \tag{8.15}$$

so that

$$\frac{dP}{P} = \frac{\gamma}{\gamma - 1} \frac{dT}{T} \tag{8.16}$$

or

$$\left|\frac{d\ln P}{d\ln T}\right|_{\rm ad} = \frac{\gamma}{\gamma - 1}\,, \tag{8.17}$$

The Schwarzschild instability criterion then becomes:

$$\left|\frac{d\ln P}{d\ln T}\right|_{\rm star} < \frac{\gamma}{\gamma - 1}\,. \tag{8.18}$$

Sometimes, this condition is also given as:

$$\left|\frac{dT}{dr}\right| > \frac{\gamma - 1}{\gamma} \frac{T}{P} \left|\frac{dP}{dr}\right| \tag{8.19}$$

obtained from 8.16 by dividing both sides by dr and rearranging.

The above treatment emphasises the importance of the ratio of the specific heats in determining the stability of a star against convection. Recall that $C_P > C_V$. Their difference is given by:

$$C_P - C_V = \frac{k}{\mu\, m_H}$$

so that:

$$\frac{\gamma}{\gamma - 1} = C_P \frac{\mu\, m_H}{k}\,.$$

Thus, the higher the specific heats, the closer their ratio is to 1; the closer γ is to 1, the higher is the value of the adiabatic $|d\ln P/d\ln T|$ gradient (eq. 8.17), leading to instability (eq. 8.18). Under what conditions are C_P and C_V high? One example is partially ionised gas, where some of the heat supplied to the system may go into further ionizing atoms instead of increasing their kinetic energy. Similarly, the presence of molecules in the mix, would increase the specific heats since some of energy supplied would go into breaking the molecular bonds.

8.3.1 Which Stars Are Convective?

From the discussion above we can develop some physical understanding of when a star will develop a convective layer in its core. From eq. 8.1 we saw that the temperature gradient is proportional to the opacity κ; thus, we expect that in layers where the opacity is very high, the temperature gradient required for radiative energy transport becomes unachievably steep. We also saw from Figure 5.5 that stellar opacities increase dramatically as T decreases from 10^6 to 10^5 K; at $T \sim 10^5$ K the gas is only partly ionised (at typical stellar densities); the rise in κ is produced by the availability of many bound-bound and bound-free transitions.

Hand in hand with this is the increase in the specific heats; consequently their ratio γ tends to 1 as the gas becomes partly, as opposed to fully, ionised. As we have discussed, this will increase the adiabatic $|d \ln P/d \ln T|$ gradient, leading to convection. For both reasons, convection will occur in the outer layers of cool stars. In a G0 V star the convective layer is thin, while main sequence M stars are almost fully convective. Red giants and supergiants are also convective over most of their interiors. Figure 8.5 shows the onset of convection in the Sun at a radius $r \simeq 0.7R_\odot$.

Convection is also important in stellar layers where the ratio $L/4\pi r^2$ is high (cfr. eqs. 8.1 and 8.11), that is where large luminosities are generated over small volumes. This is the situation in the cores of massive stars, given

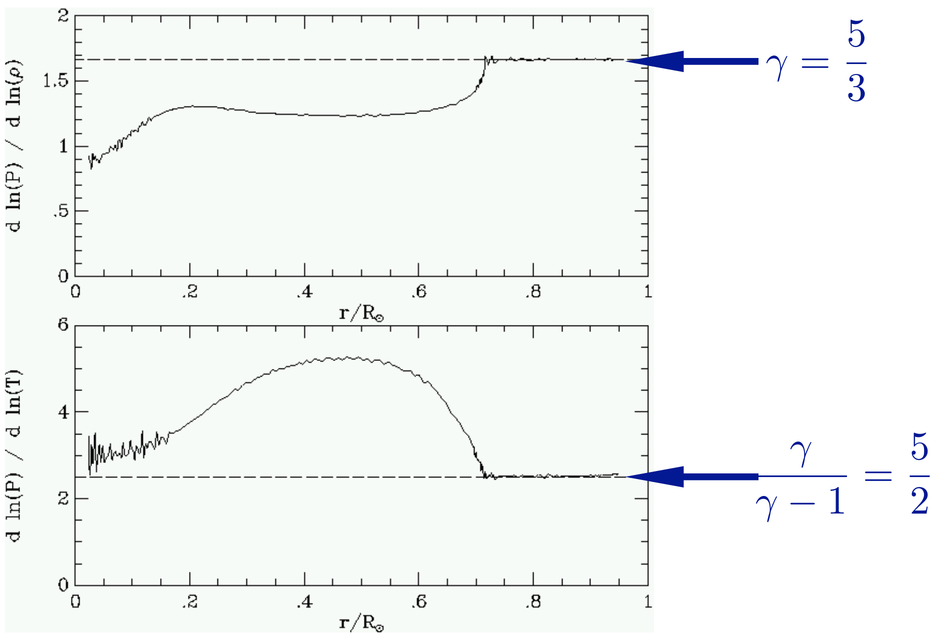

Figure 8.5: Plots of $d \ln P/d \ln \rho$ (top) and $d \ln P/d \ln T$ for the Standard Solar Model. For a monoatomic gas $\gamma = 5/3$.

the steep temperature dependencies of the CNO cycle and the triple-alpha process (Lecture 7.4.2 and 7.4.3).

Figure 8.6 summarises pictorially the points made in this section.

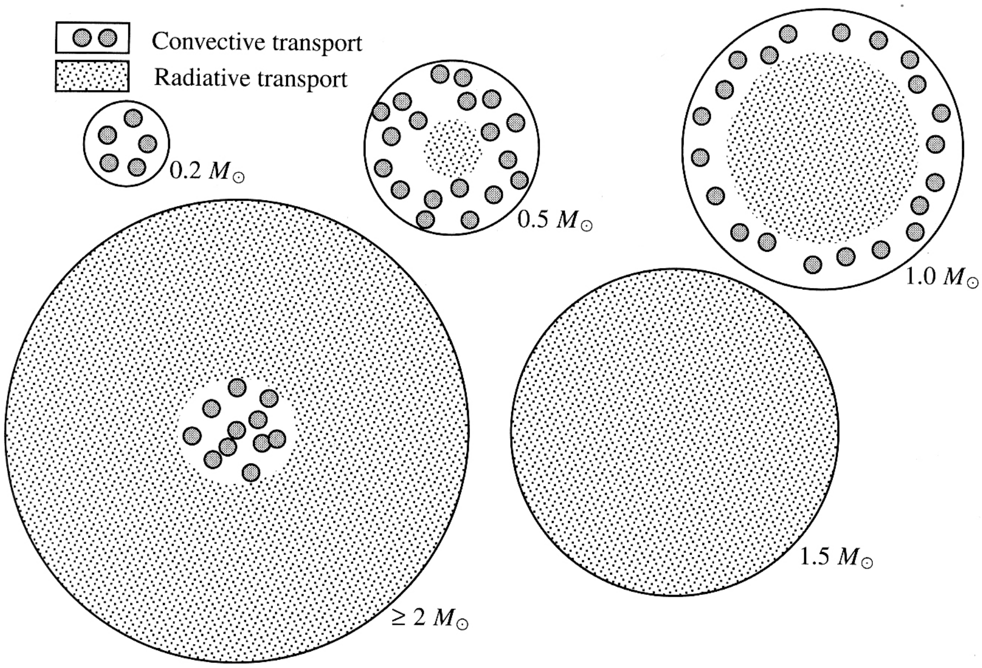

Figure 8.6: Zones of convection and radiation in main-sequence stars of various masses. The lowest mass stars are completely convective. A radiative core develops at $M \simeq 0.4 M_\odot$, and a star is fully radiative at $M \simeq 1.5 M_\odot$. The core region is again convective for masses $M \gtrsim 2 M_\odot$. The relative sizes of the stars shown here are approximately correct, while on the main sequence. (Figure reproduced from Bradt, H., Astrophysics Processes, CUP).

8.4 Hydrostatic Equilibrium

Having established the conditions under which convection will occur, it may be of interest to estimate the convective flux F_c, that is the energy transported by convective cells per unit time through a unit surface area of the star. In order to do so, we need to first consider the concepts of hydrostatic equilibrium and pressure scale height.

8.4.1 Hydrostatic Equilibrium

What stops a star from collapsing under its own gravity? The inward force of gravity must be balanced by an equal force of opposite sign; this force is provided by pressure. Considering a small cylindrical element of mass dm and height dr located at a distance r from the centre of a spherically symmetric star (Figure 8.7), we have:

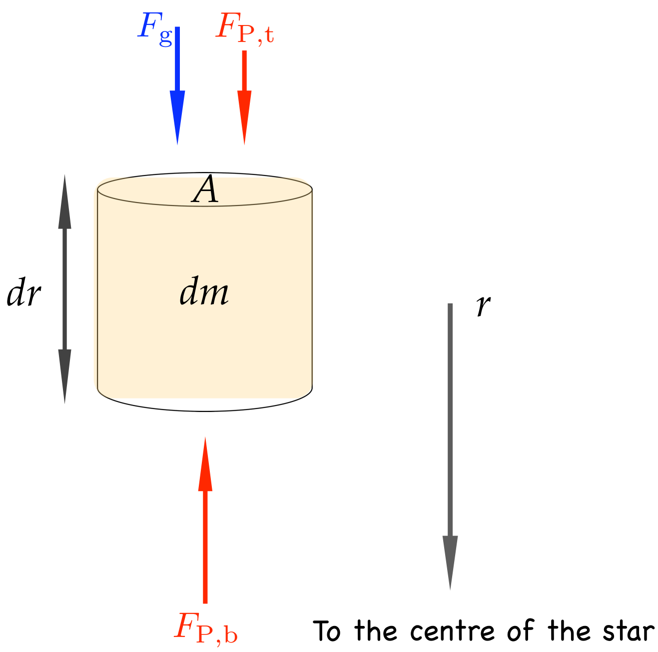

Figure 8.7: In hydrostatic equilibrium $|F_g + F_{P,t}| = |F_{P,b}|$.

$$A\,dP = -G\frac{M_r\,dm}{r^2} \tag{8.20}$$

where: (i) M_r is the mass enclosed within radius r; (ii) A is the area of the base of the cylinder; and (iii) dP is the difference in pressure (defined as force per unit area, i.e. $P \equiv F_p/A$) between the top and bottom faces of the cylinder. Since we are considering an infinitesimally small cylinder, we can assume that the density is constant within the cylinder and express the mass in terms of the density: $dm = \rho\,A\,dr$:

$$A\,dP = -G\frac{M_r\,\rho\,A\,dr}{r^2}\,. \tag{8.21}$$

Dividing both sides by the volume of the cylinder $dV = A\,dr$, we obtain the equation of hydrostatic equilibrium:

$$\boxed{\frac{dP}{dr} = -G\,\frac{M_r\,\rho}{r^2}} \tag{8.22}$$

Since $G\,M_r/r^2 \equiv g$, the local acceleration of gravity at radius r, we can also write:

$$\frac{dP}{dr} = -\rho\,g\,. \tag{8.23}$$

In order for a star to be in hydrostatic equilibrium, a *negative pressure gradient* must exist within the star, with the pressure being larger in the interior than near the surface.

8.4.2 Dynamical Timescale

What would happen if pressure suddenly vanished and the only force acting on a star were gravity? The star would collapse on a free-fall timescale given by:

$$t_{\rm ff} = \sqrt{\frac{2R}{g}} \tag{8.24}$$

where R is the star radius (recall $r = \frac{1}{2}at^2$ for a body experiencing acceleration a). With the above definition of g, and expressing M_R in terms of the mean density $\langle\rho\rangle$, $M_R = \frac{4}{3}\pi R^3\,\langle\rho\rangle$, we have:

$$t_{\rm ff} = \sqrt{\frac{2R\,R^2}{GM_R}} = \sqrt{\frac{2R\,R^2\,3}{G\,4\pi R^3\,\langle\rho\rangle}} = \sqrt{\frac{3}{2\pi}}\sqrt{\frac{1}{G\,\langle\rho\rangle}} \tag{8.25}$$

The last term in the equation is often referred to as the *dynamical timescale*:

$$\boxed{t_{\rm dyn} \sim \sqrt{\frac{1}{G\rho}}} \tag{8.26}$$

Although we have derived it by using the unrealistic example of the free-fall time in a star, the dynamical timescale is an important concept often used in dimensional treatment of astrophysical situations. It describes the time taken for changes in one part of a body to be communicated to the rest of that body (and thus $t_{\mathrm{dyn}} \approx t_{\mathrm{sc}}$ where t_{sc} is the sound crossing time). Another way to think about the dynamical timescale is as the time required for a system to move from one equilibrium state to another after a sudden change. Thus, for example, an interstellar cloud cannot collapse to form stars over a timescale shorter than t_{dyn}; similarly, a sudden burst of star formation in a galaxy cannot take place over a timescale shorter than t_{dyn}.

8.4.3 Pressure Scale Height

Returning to the equation of hydrostatic equilibrium, we define the pressure scale height H_p as:

$$\frac{1}{H_p} \equiv -\frac{1}{P}\frac{dP}{dr}\,, \tag{8.27}$$

so that we can express the variation in pressure with radius as:

$$P = P_0\, e^{-r/H_p}\,. \tag{8.28}$$

In other words, H_p is the radial distance over which the pressure drops by a factor of e. Recalling eq. 8.23, it can be seen that:

$$H_P = \frac{P}{\rho g}\,. \tag{8.29}$$

8.5 Mixing Length Theory

Let us now return to the question of the energy transport by convection. A buoyant bubble will rise until the temperature inside gradually adjusts to the (lower) temperature of its surroundings. The distance that a hot cell rises, or a cold one sinks, is referred to as the *mixing length*, and is normally expressed in terms of the pressure scale height:

$$\ell = \alpha H_P$$

where $\alpha \sim 1$ is a free parameter (meaning simply that we do not have a physical theory to relate ℓ to H_P). In other words, the mixing length is of the order of the pressure scale height.

As the cell travels towards the stellar surface by one mixing length, the excess heat flow per unit volume from the bubble into its surroundings can be written as:

$$\delta q = (C_P \, \delta T)\rho \,, \tag{8.30}$$

where C_P is the specific heat at constant pressure and

$$\delta T = \delta \left(\frac{dT}{dr} \right) \, dr \tag{8.31}$$

is the temperature difference between the surroundings and the rising bubble. The convective flux is simply:

$$F_c = \delta q \, \langle v_c \rangle \tag{8.32}$$

where $\langle v_c \rangle$ is the average speed of the convective bubble. We can obtain an expression for $\langle v_c \rangle$ by equating the kinetic energy of the rising bubble to some fraction β $(0 < \beta < 1)$ of the work per unit volume done by the buoyant force over the distance ℓ:

$$\frac{1}{2}\rho \langle v_c \rangle^2 = \beta \, \langle f_{\rm net} \rangle \, \ell \tag{8.33}$$

so that:

$$\langle v_c \rangle = \left(\frac{2\beta \langle f_{\rm net} \rangle \, \ell}{\rho} \right)^{1/2} \,. \tag{8.34}$$

where β is another free parameter.

The net force on the bubble is just

$$f_{\rm net} = -g \, \delta\rho \tag{8.35}$$

We can express $\delta\rho$ in terms of δT using the the ideal gas law

$$\frac{\delta P}{P} = \frac{\delta\rho}{\rho} + \frac{\delta T}{T} \tag{8.36}$$

or, since we have pressure equilibrium between the bubble and its surroundings (i.e. $\delta P = 0$),

$$\delta\rho = -\frac{\rho}{T} \, \delta T \,. \tag{8.37}$$

Thus,

$$f_{\rm net} = \frac{\rho g}{T} \, \delta T \tag{8.38}$$

and

$$\langle f_{\rm net} \rangle = \frac{1}{2}\,\frac{\rho g}{T}\,\delta T \tag{8.39}$$

taking the average of $f_{\rm net}$ over the distance ℓ travelled by the convective cell. Substituting 8.39 into 8.34, we have:

$$\langle v_c \rangle^2 = \beta\,\frac{\rho g}{\rho T}\,\ell\,\delta T \tag{8.40}$$

$$\langle v_c \rangle^2 = \beta\,\frac{g}{T}\,\alpha H_p\,\delta T \tag{8.41}$$

$$\langle v_c \rangle^2 = \beta\,\frac{g}{T}\,\alpha\,\frac{P}{\rho g}\,\delta T \tag{8.42}$$

Using the ideal gas law: $P = (\rho k T)/(\mu m_H)$, we have:

$$\langle v_c \rangle^2 = \beta\,\alpha\,\frac{k}{\mu m_H}\,\delta T \tag{8.43}$$

Using 8.31 with $dr = \ell = \alpha H_p = \alpha P/\rho g$ and proceeding as before for P, we have:

$$\langle v_c \rangle^2 = \beta\,\alpha\,\frac{k}{\mu m_H}\,\delta\left(\frac{dt}{dr}\right)\alpha\frac{kT}{\mu m_H}\,\frac{1}{g} \tag{8.44}$$

$$\langle v_c \rangle^2 = \beta\,\alpha^2\left(\frac{k}{\mu m_H}\right)^2\frac{T}{g}\,\delta\left(\frac{dt}{dr}\right) \tag{8.45}$$

Substituting 8.45 into 8.32 and using again 8.30 and 8.31, we finally obtain an expression for the convective flux:

$$F_c = \rho C_P\left(\frac{k}{\mu m_H}\right)^2\left(\frac{T}{g}\right)^{3/2}\beta^{1/2}\left[\delta\left(\frac{dT}{dr}\right)\right]^{3/2}\alpha^2 \tag{8.46}$$

What we have described here is an example of mixing-length theory. The unsatisfactory aspects of it are the ‘fudge-factors’ α and β which are generally adjusted to fit observations. However, it is a significantly simpler treatment than the full 3-D (magneto-)[1]hydrodynamical calculations of convective flows (see Figure 8.8).

In closing, convection is a complicated topic and remains an active area of research, in particular because it can have considerable impact on a wide

[1]Yes, we have not mentioned them yet, but magnetic fields are clearly important when we are dealing with bulk motions of ionised gas!

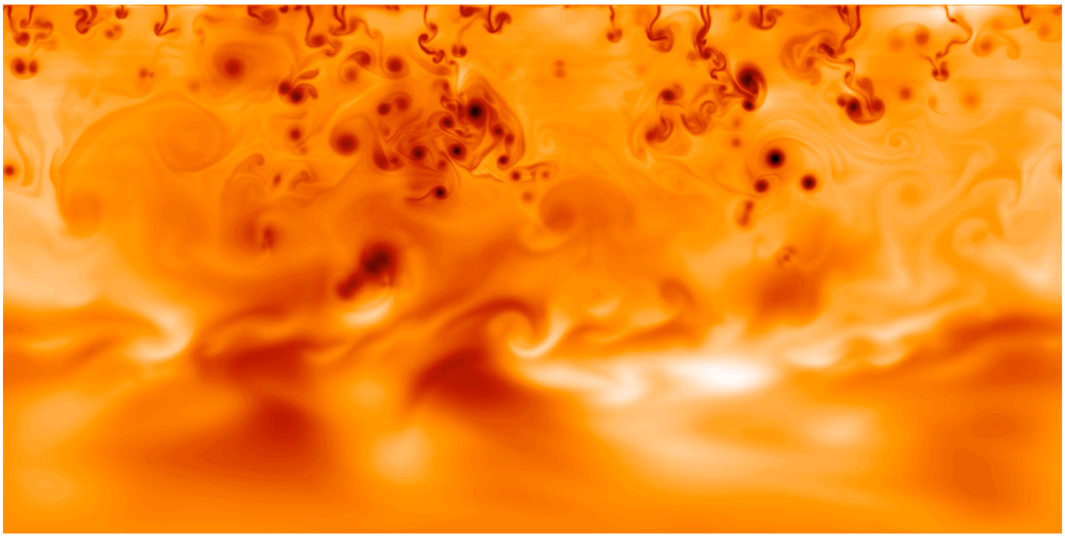

Figure 8.8: Computer simulation of convection.

range of stellar properties and on stellar evolution. The lack of a good theory of convection, and of the amount of energy that can be transferred by convection, is at present an important limitation in our understanding of stellar structure.

However, convection remains an important phenomenon in the interiors of stars, and not only as a way of transporting energy to the surface where it is radiated away. Convection is also important as a means of *mixing* between different layers of the star. Mixing happens because of *convective overshooting*: at the top and bottom boundaries of the convection zone, even though there is no net force, a rising or falling cell will arrive with a finite velocity and will overshoot. This was first realised in the 1980s from observations of the nuclear products at the surface of massive stars, indicative of some degree of mixing to layers above the convective core boundary. Stellar rotation also greatly facilitates mixing between different stellar layers (see Figure 8.9). The mixing of freshly synthesised elements into the outer layers of stars has important consequences for their evolution.

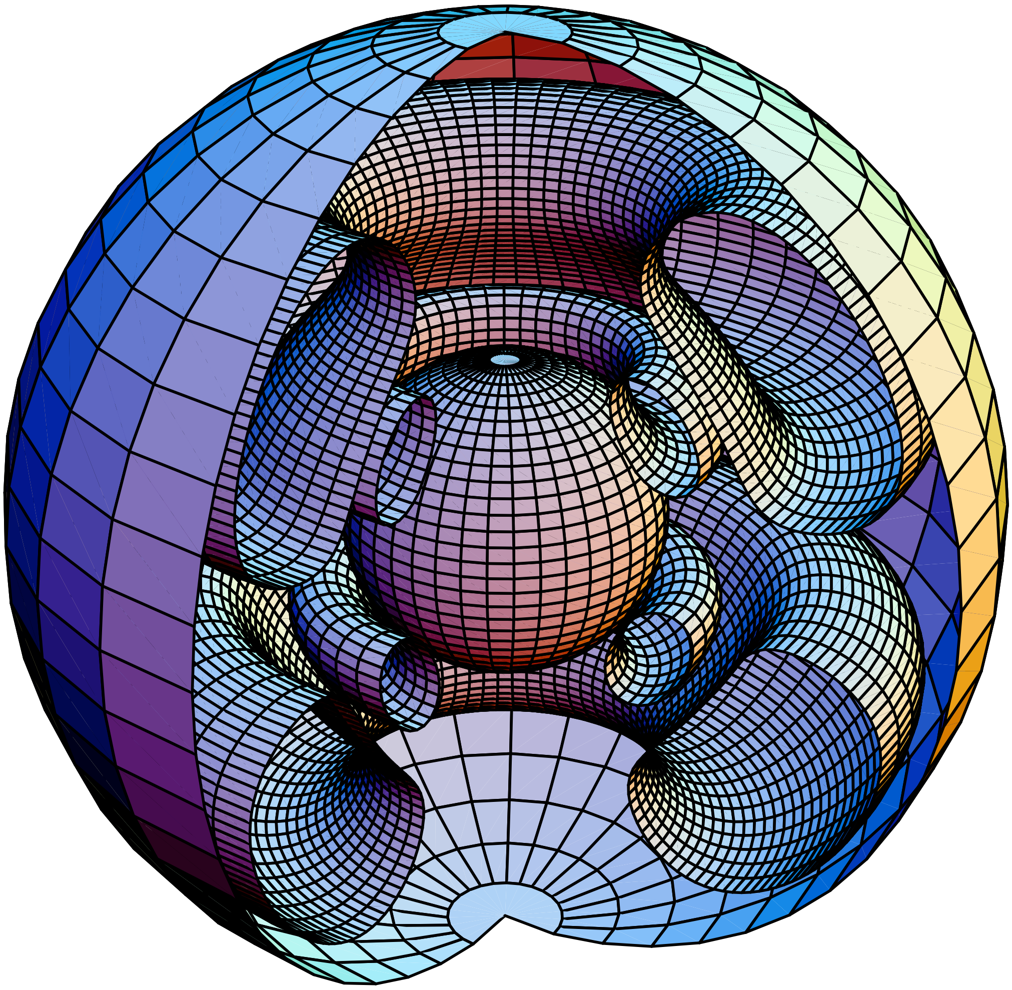

Figure 8.9: Stream lines of meridional circulation in a rotating $20M_\odot$ model with solar metallicity and $v_{\rm rot} = 300$ km s^{-1} at the beginning of the H-burning phase. The stream-lines are in the meridian plane. In the upper hemisphere on the right section, matter is turning counterclockwise along the outer stream line and clockwise along the inner one. The outer sphere is the star surface and has a radius equal to $5.2R_\odot$. The inner sphere is the outer boundary of the convective core. It has a radius of $1.7R_\odot$. (Figure reproduced from Meynet & Maeder 2002, A&A, 390, 561).