SOLAR CONVECTION

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Abstract. The hydrodynamics of solar convection is reviewed. In particular, a discussion is given of convection on the scale of granulation; i.e., the energy carrying convection patterns in the solar surface layers, and its penetration into the stable layers of the solar photosphere. Convection on global and intermediate scales, and interaction with rotation and magnetic fields is discussed briefly.

1. Introduction

The dynamics of convection in the solar plasma is governed by the equations of (magneto-)hydrodynamics, and also, in the solar surface layer, by the equations of radiative transfer. Because of the extremely weak molecular dissipative processes (viscosity, heat conductivity, electrical resistivity), the resulting behavior of the plasma is quite complicated. The non-linear transport terms (representing advection of plasma properties) dominate the behavior 'almost everywhere', forcing dissipative processes to become significant only in the fractionally small volume where extreme gradients develop. This kind of behavior is next to impossible to describe by analytical methods, but is marginally tractable on present-day computers. Thus, recent advances in the understanding of solar convection, and its interaction with rotation and magnetic fields, have come partly through the development of computer codes capable of numerically simulating the behavior of these systems.

Traditionally, one refers to convection on three size scales on the Sun; global convection, on the scale of hundreds of Mm, extending throughout the entire convection zone; supergranular convection, on the scale of tens of Mm, and granular convection, on the scale of Mm. Recently (November et al., 1981, 1982), mesogranulation has been introduced as a notation for observed velocity and temperature fluctuations on a scale intermediate between those of granulation and supergranulation.

There is a corresponding set of size scales for magnetic activity on the Sun; with activity complexes (Bumba, 1975) on the scale of global convection, active regions on the scale of supergranulation, and the solar filigree (Dunn and Zirker, 1973; Mehlretter, 1974; Muller, 1983) on the scale of granulation. Most likely, these magnetic phenomena are dynamically related to the corresponding scales of convection, but details of these relations are not yet well understood. This review concentrates mostly on convection on a granular scale, but larger scale convection, and interaction with magnetic fields are also touched upon.

The equations of hydrodynamics and their basic physical interpretation are discussed in Section 2. The application to convection on a granular scale, and its influence on the solar photosphere is reviewed in Section 3. Section 4 briefly reviews some recent
numerical work on convection on a global scale and also discusses the less well studied intermediate scales ('super-' and 'meso-granulation'). The interaction of convection on various scales with magnetic fields is discussed in Section 5.

2. Hydrodynamics of Solar Convection

2.1. Basic Equations

For the purpose of the following discussion, we write the equations of (non-relativistic) fluid dynamics with radiative energy transport in a form suitable to describe convection on the scale of granulation, mesogranulation, and supergranulation on the Sun (rotation, radiation pressure, and $v/c$ terms unimportant):

$$
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{u}) = 0 ,
$$

$$
\frac{\partial \mathbf{u}}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{u} = -\nabla P + \rho \mathbf{g} + \rho \mathbf{f}_{\text{visc}} ,
$$

$$
\frac{\partial H}{\partial t} + \mathbf{u} \cdot \nabla H = Q \frac{\partial P}{\partial t} + \mathbf{u} \cdot \nabla P ,
$$

$$
\rho = \rho(P, T) ,
$$

$$
H(P, T) = E(P, T) + P/\rho ,
$$

$$
Q = Q_{\text{rad}} + Q_{\text{visc}} ,
$$

$$
Q_{\text{rad}} \equiv -\nabla \cdot F_{\text{rad}} = \int_0^\infty \int_0^{4\pi} \rho \kappa_v (I_{v\Omega} - S_v) \, d\Omega \, dv ,
$$

$$
\frac{dI_{v\Omega}}{ds} = \rho \kappa_v (I_{v\Omega} - S_v) ,
$$


Although this set of coupled differential equations may appear complicated (and may require substantial effort to solve numerically), the individual equations are conceptually simple, and the qualitative properties of solar convection may be fruitfully discussed with reference to the individual equations.

2.2. The Continuity Equation

It turns out to be convenient, both for the purpose of numerical work (Nordlund, 1982) and for a qualitative discussion, to apply a horizontal Fourier transform to the continuity equation, and to the equation of motion. For a discussion of the advantages of such spectral methods in (non-shocking) hydrodynamics calculations, see Orzag (1971) and
Herring et al. (1974). Symbolically, we may write, for example,

$$\rho u(x, y, z) = \sum_{lm} \psi_{lm}(z) e^{i(kx + my)} \quad (9)$$

(although, in practice, explicit sine and cosine transforms are numerically more efficient). With this representation, it is easy to show (Nordlund, 1978) that, for slowly changing flows, the ratio of the horizontal to vertical velocities of a certain Fourier component is inversely proportional to the horizontal wavenumber (which we may denote $k$):

$$\frac{u_{\text{hor}}}{u_{\text{vert}}} \approx \frac{1}{k H_\psi}, \quad (10)$$

where $H_\psi$ is the scale height of the vertical mass flux, $(\psi_z)_k = (\rho u_z)_k$. In other words, the larger the convective cell, the larger the horizontal velocities must be, for given vertical velocities. The reason is that the rapidly decreasing vertical mass flux in the ascending parts of the cell must be balanced by a horizontal outflow, and for larger cells a given scale height $H_\psi$ becomes a smaller fraction of the size of the cell. In the case of granulation, the ratio of horizontal to vertical velocity in the photosphere is just slightly larger than unity, but for supergranular flows the ratio must be an order of magnitude larger. This is the reason why supergranulation is seen primarily as a horizontal velocity field.

2.3. The anelastic approximation

The qualitative difference between pressure waves, gravity waves, and slowly changing flows is apparent through a consideration of the continuity equation, Equation (1). In pressure waves, the converging flow ($\nabla \cdot (\rho u) < 0$, $\partial \rho / \partial t > 0$) causes an overpressure that halts and reverses the compression phase of the wave. In slowly changing flows, and in slow gravity waves, the continuity equation acts only as a constraint on the flow, allowing only flows with (approximately) divergence-free mass fluxes:

$$\nabla \cdot (\rho u) \approx 0. \quad (11)$$

Adopting this anelastic approximation (Batchelor, 1953; Charney and Ogura, 1960) one finds that the pressure is (apart from boundary conditions) uniquely determined by the divergence of the (volume) forces (including the 'inertial' force-term $-\rho \mathbf{u} \cdot \nabla \mathbf{u}$):

$$\nabla^2 P = \nabla \cdot \{\text{volume forces}\} = \nabla \cdot \{\rho (\mathbf{g} + \mathbf{f}_{\text{visc}} - \mathbf{u} \cdot \nabla \mathbf{u})\} \quad (12)$$

(cf. Nordlund, 1982). This particular pressure has a gradient that conspires with the force field to keep

$$\frac{\partial}{\partial t} (\nabla \cdot (\rho \mathbf{u})) = 0, \quad (13)$$

i.e., the mass flux divergence remains zero. By Equation (12), the pressure responds
instantaneously to changes in the force field; i.e., the system behaves as if the sound speed was infinite. From the foregoing discussion is clear that the anelastic approximation will serve to exclude the possibility of pressure waves from the system, while still allowing (divergence-free) flows.

The anelastic approximation has been used in many forms (Batchelor, 1953; Charney and Ogura, 1960; Ogura and Phillips, 1962; Gough, 1969; Latour et al., 1976; Nordlund, 1978; Gilman and Glatzmaier, 1981), often with an expansion of the equations in terms of some (small) measure of the fluctuations.

Note, however, that according to the discussion above, the validity of the anelastic approximation has to do with how slowly or rapidly the density changes in an Eulerian frame, which is not necessarily directly related to the magnitude of the velocity. Indeed, a stationary hydrodynamical system with arbitrary velocities satisfies Equation (11) exactly. Criteria having to do with velocity or perturbation amplitudes come (directly or indirectly) from a translation of velocity amplitude to time scale, using some characteristic scale length. In the case of the solar granulation, the quasi-stationary radiative cooling of the ascending gas in the granules make the patterns evolve on a much longer time scale (tens of minutes) than, e.g., the travel time over a pressure scale height (of the order of a minute).

2.4. THE EQUATION OF MOTION

The horizontal flows required by the continuity equation (from ascending to descending parts of the convection cells) must be supported by a corresponding pressure pattern since, by the equation of motion, only pressure gradients are available to accelerate the horizontal flow away from updrafts and to decelerate horizontal flows coming towards downdrafts. Thus, we may expect pressure excess over both updrafts and downdrafts, with pressure deficiency in between. (This discussion applies to layers near the upper boundary of an unstable layer, where the strongest horizontal velocities occur, because of the rapid decrease of the vertical mass flux.) The pressure excess over ascending and descending material decreases the buoyancy of ascending material ('buoyancy braking', Massaguer and Zahn, 1980; Latour et al., 1983) but also increases the weight of the descending material, an effect that is not brought out by single-mode calculations such as those of Massaguer and Zahn and Latour et al. This asymmetric buoyancy modification is one effect that contributes to the asymmetry between updrafts and downdrafts in numerical simulations of convection in stratified media (Hurlburt, 1983; Hurlburt et al., 1984, 1985; Toomre et al., 1984; Hurlburt and Toomre, 1985). Another effect (Nordlund, 1978) is produced by the advection of fluid properties by the non-linear term in the equation of motion; with horizontal flows from updrafts to downdrafts in a stratified medium, ascending material will tend to spread with time, while descending material will tend to narrow. (This is again a consequence of the advection of properties by a low-viscosity fluid flow; in this case the flow topology itself is the advected quantity.)

The non-linear terms in the equation of motion and the energy equation produce a
steepening of the gradients of advected quantities in regions where the fluid velocity decreases downstream. Because the viscosity of the solar plasma is very small, viscous effects that halt the steepening set in on scales that cannot (by many orders of magnitude) be resolved in numerical models. To prevent the buildup of steeper gradients than can be numerically resolved, the numerical scheme must contain suitable dissipative (diffusive) terms. Sufficiently large ordinary diffusion terms ('turbulent viscosity') will do the job, but will also influence the resolved motions appreciably. Other types of diffusive terms, which influence resolved motions less, while still affecting the smallest resolved scales sufficiently to prevent excessive buildup of gradients, are to be preferred. Such terms are often referred to as 'sub-grid-scale eddy diffusivities' (e.g., Smagorinsky, 1963). A particular form of such diffusive terms suitable for the semi-spectral representation Equation (9) was derived by Nordlund (1982, Section 3.3).

2.5. WAVES AND WAVE GENERATION

The interaction of waves (pressure waves, gravity waves) and flows in the solar envelope is, for several reasons, a difficult problem:

(1) Wave generation by the turbulent convective flow is a complicated process, where so far only exploratory work based on order of magnitude estimates has been made (e.g., Goldreich and Keeley, 1977b).

(2) The large fluctuations in temperature, pressure, and thermodynamic quantities (e.g., heat capacity and sound speed) associated with the convection. Sufficiently large fluctuations must influence the propagation and dispersion properties of waves; e.g., because the propagation speed becomes a strong function of horizontal position. Some remaining systematic discrepancies between the observed and calculated frequencies of solar p-modes in the 5-min period range (Christensen-Dalsgaard, 1984, 1985) may possibly be related to the question of the wave propagation speed in a very inhomogeneous medium. The upper approximately 500 km of the solar convection zone is very inhomogeneous, and presumably large scale waves propagate with a different speed than in a plane-parallel medium with the same average temperature stratification.

(3) The convective velocity field, which may cause scattering and aberration of waves. Smaller scale waves, especially of a size comparable to the size of granulation, may be expected to suffer refraction, reflection, and advection by the velocity field associated with the convection. To my knowledge, this problem has not been studied in the solar context.

(4) The modulation of the convective energy flux by large scale waves, which changes the stability properties of standing wave modes. Attempts to estimate modulation of convection by pressure waves have been made by, e.g., Unno (1976), Gough (1977), and Goldreich and Keeley (1977a). However, these attempts were based on a mixing length description of convection, and thus the results are of uncertain validity.

(5) Radiative transfer effects in the surface layers, which are crucial both to the convection and to wave propagation and reflection. Non-adiabatic effects on pressure wave propagation in the solar photosphere have been studied by Christensen-Dalsgaard and Frandsen (1983), who concluded that effects of departures from radiative
equilibrium in the mean state may change the stability properties of standing waves drastically. Non-adiabatic effects on internal gravity waves have been studied (in the Newtonian cooling approximation) by Mihalas and Toomre (1982).

It would be perfectly possible to study these problems by solving the linear wave equations separately in a background medium obtained, for example, from anelastic simulations of the convection. This would not be a very useful approach, though, because the solution of the wave equations under these conditions (large amplitude fluctuations of the thermodynamic quantities and flows with significant non-linear terms) would be about as time consuming as a solution of the full set of equations. A separate solution of the linear wave equations would not, however, provide any information on the back reaction of the waves on the convection. Thus there is really no point in aiming for less than a solution of the full problem, as the next logical step. The main complication in doing the full problem probably lies with implementing boundary conditions that behave reasonably. The time steps would not have to be reduced very much compared to the anelastic simulations of granular convection, since the convective velocities are a significant fraction of the sound speed in some places.

2.6. Convective penetration into stable layers

The penetration of motions from a convectively unstable zone into neighboring stable layers may have significant consequences for the temperature structure of these stable layers. A qualitative understanding of the extent to which such flow penetration occurs may be gained from a consideration of the Poisson equation for the pressure, Equation (2) (cf. Nordlund, 1982, Section 3.6).

In terms of the horizontal Fourier components \( P_{lm} \) of the pressure, the relative pressure fluctuations,

\[
P_{lm}(z) = P_{lm}(z)/P_{00}(z),
\]

obey the ordinary differential equation

\[
\frac{d^2 P_{lm}}{dz^2} + \alpha \frac{d P_{lm}}{dz} - k^2 P_{lm} = 0,
\]

where \( \alpha \) is the inverse of the pressure scale height,

\[
\alpha = \frac{d \ln P_{00}}{dz},
\]

where \( \alpha \) is the inverse of the pressure scale height,

\[
k^2 = (l^2 + m^2).
\]

For the purpose of the discussion, it is convenient to disregard induced temperature perturbations in the stable layers, and consider only the penetration of pressure perturbations induced from the unstable layers. This is not unrealistic for the stable layer above the convection zone, where the radiative transport of energy efficiently reduces
the temperature perturbations, and where relatively large pressure fluctuations in the top layers of the unstable zone have a strong influence in the stable layers. Below the convection zone, conditions are different, and here the induced temperature perturbations may be important in reducing the penetration of motions into the stable layer (undershoot). However, as we shall see, even disregarding temperature perturbations, the penetration is much weaker below a strongly stratified convection zone than above. The general solutions to the differential equation (15) are of the type

\[ p_{lm} \approx e^{\alpha \pm \varepsilon}, \]  

where

\[ \alpha = -\alpha/2 \pm (\alpha^2/4 + k^2)^{1/2}. \]  

In the limit of large \( k \) (disturbances small compared with the scale height; \( k \gg \alpha \)), we have

\[ \alpha \rightarrow \pm k, \]  
i.e., small scale pressure disturbances at the convection zone boundary decay away over a distance comparable with their size.

In the limit of large scale disturbances, there is a qualitative difference in the situation below and above an unstable layer: In a stable layer above an unstable layer, large scale fluctuations have a long range influence (since \( \alpha_+ \rightarrow k^2/\alpha \)), whereas in a stable layer below an unstable layer large scale disturbances lose importance over about one pressure scale height (\( \alpha_- \rightarrow -\alpha \)). Qualitatively, this is because in a stratified atmosphere the mean state pressure increases rapidly below the unstable layer, which rapidly diminishes the relative important of a certain (absolute) pressure disturbance, while the opposite is true above an unstable layer.

This result has important consequences for the character of the motions generated in stable layers above and below convection zones. A forced motion in a stable layer, if sufficiently slow, sets up propagating internal gravity waves. Internal gravity waves are characterized by a back-and-forth motion, where the restoring force is due to the excess weight of the displaced fluid. However, for sufficiently strong perturbations (as for example in the case of the solar photosphere), the induced motion will have the character of a flow, where the fluid goes from A (an upflow) to B (a downflow), never to return.

The essence of the results cited above is that it is easier to drive (mixing) penetration into stable layers above than into stable layer below a convection zone. The character of the penetration, if sufficiently strong, is one of large scale flows (Nordlund, 1984b) rather than the excitation of internal gravity waves (Mihalas and Toomre 1981, 1982). The qualitative difference between penetration above and below an unstable layer is nicely illustrated in the work by Hurlburt et al. (1985).

Superficially, the results of Massaguer et al. (1984) would seem to be in conflict with the conclusions on penetration stated above. Massaguer et al. conclude, on the basis of numerical solutions of anelastic one- or two-mode equations for a strongly stratified
medium, that the penetration below an unstable layer is more pronounced than that above. For three reasons, this conclusion cannot be directly confronted with the results quoted above: (1) because they compare the penetration in terms of physical depth, rather than in terms of the local scale height; (2) because their strong downward penetration occurs only for a planform with downward directed flow at cell center, a topology that is not found in three-dimensional multi-mode calculations; (3) because they do not regard a flow with strong horizontal velocities but weak vertical velocities as penetrating. However, in terms of local density or pressure scale heights the penetration is actually much stronger in their stable layer above than in their stable layer below, especially for the planform with upward directed flow at cell center, and for the case with two coexisting oppositely oriented planforms (cf. their Figures 1, 3, and 5).

2.7. THE ENERGY EQUATION AND RADIATIVE ENERGY TRANSFER

The energy equation, Equation (3), describes the evolution of temperature (as measured by the enthalpy) as a result of pressure changes \( (\rho \frac{\partial H}{\partial t} + \rho u \cdot \nabla H - \frac{\partial P}{\partial t} - u \cdot \nabla P = 0) \) describing adiabatic advection; \( DS/\Delta t \equiv 0 \), and non-adiabatic heating/cooling, due to contributions from, e.g., radiation and viscous dissipation, Equation (6).

The most important non-adiabatic effect in the solar photosphere is due to the radiative term \( Q_{\text{rad}} \). Indeed, the granulation pattern results from the release (through the \( Q_{\text{rad}} \) term) of excess heat carried to the surface by ascending material in the convection flow. The \( Q_{\text{rad}} \) term is also responsible for a substantial heating of the optically thin parts of the photosphere (cf. below, Section 2.8).

From Equations (3) and the leftmost part of (7), one may estimate the ascent velocity necessary for the advection of enthalpy to balance the solar luminosity:

\[
\rho \Delta H/\Delta z = F_\odot/\Delta z, \tag{21}
\]

\[
u_z \approx \frac{F_\odot}{\rho \Delta H}, \tag{22}
\]

which comes out to be of the order of 2 km s\(^{-1}\) for solar conditions (Nordlund, 1982). Thus, material ascending slower than about 2 km s\(^{-1}\) will suffer net cooling (in an Eulerian frame) at the surface. This is one of the limiting factors in the growth of granules with time; the growth of individual granules (caused by the advection by the horizontal flows from updrafts to downdrafts) is accompanied by a steady increase in pressure excess in the granule with time. The consequent loss of buoyancy leads to a decrease in the ascent velocity. When the ascent velocity falls below the critical value, Equation (22), the temperature excess can no longer be upheld in the center of the granule. This process is most obvious in exceptionally undisturbed granules that grow symmetrically, where the retardation/cooling at the center results in a bright granule with a dark center (‘exploding granules’, Namba and van Rijsbergen, 1977; Bray et al., 1984, Section 2.3.7).
In optically thin layers the radiative flux changes relatively little, and it is more enlightening to discuss the radiative energy exchange in terms of the rightmost part of Equation (7); i.e., the radiative heating/cooling is expressed as the difference between absorption \( \rho \kappa_v I_{v \Omega} \) and emission \( \rho \kappa_v S_v \), integrated over frequency and angle.

2.8. THE ENERGY BALANCE OF THE UPPER PHOTOSPHERE

In one-dimensional hydrostatic models of the solar photosphere, the temperature structure of most of the photosphere is determined by \textit{radiative equilibrium}, which may be written as an equilibrium between absorption \( (\kappa_v I_{v \Omega}) \) and emission \( (\kappa_v S_v) \) of radiation:

\[
Q_{\text{rad}} = \int \int \rho \kappa_v (I_{v \Omega} - S_v) \, d\nu \, d\Omega = 0,
\]

where \( \kappa_v \) is the monochromatic absorption coefficient (m\(^2\) kg\(^{-1}\)), and \( I_{v \Omega} \) and \( S_v \) are the monochromatic radiation intensity and source functions, respectively (W Hz\(^{-1}\) sterd\(^{-1}\)).

Qualitatively, the positive (heating) contributions to this integral come from frequencies and angles where the radiation temperature is higher than the local kinetic temperature, and \textit{vice versa}. In an atmosphere with frequency independent (grey) absorption coefficient, this is basically a balance between heating in out-going rays, and cooling in in-coming rays. The exact solution for the temperature structure is well known (e.g., Mihalas, 1978, p. 72), and would, in the solar case, correspond to an essentially constant temperature (approximately equal to 4600 K) in most of the photosphere.

The presence of hundreds of thousands of spectral lines in the solar spectrum changes the energy balance qualitatively, and complicates the calculation of the temperature structure considerably (cf. the review by Carbon, 1979).

Qualitatively, for models in radiative equilibrium, the energy balance in the layers that are optically thin in the continua is one between heating in the continua and cooling in spectral lines: at a certain height in the atmosphere, the integrand in Equation (23) is positive for all frequencies where the monochromatic optical depth is much smaller than unity (i.e. in the continua, in weak lines, and in the wings of strong lines). However, in frequencies where the optical depth is of the order of unity (at the same certain height in the atmosphere), the out-going intensity is close to the local source function, while the in-coming intensity is significantly lower. Thus, for these frequencies, the integrand in Equation (23) is negative (cooling). As a result, the overall temperature of the upper layers of an atmosphere in radiative equilibrium is reduced (‘surface cooling by spectral lines’), such that the cooling in the spectral lines is compensated for by a corresponding heating in the continuum.

In frequencies where the optical depth is large, the radiation intensity approaches the source function (thermalizes) exponentially with optical depth, and consequently such frequencies contribute little to the integral in Equation (23). On the other hand, since the continuum radiation is ‘blocked’ in these frequencies, the temperature of the continuum forming layers has to be increased (relative to a grey model), to conserve the
total emitted radiation flux (effective temperature). This is usually referred to as ‘backwarming’.

The surface cooling by spectral lines (and the backwarming) determines the temperature structure of the upper photosphere in one-dimensional, static models of stellar atmospheres in radiative equilibrium (cf. Athay, 1970c; Gustafsson et al., 1975; Kurucz, 1970). In particular, such (‘line blanketed’) models of the solar photosphere have a temperature structure much closer to empirical models than models with only continuum absorption included.

The energy balance of the upper photosphere is qualitatively different in 3-D hydrodynamical models of stellar atmosphere (Nordlund, 1984b; Nordlund and Dravins, 1986). Because of the presence of a velocity field in the photosphere, the temperature structure is no longer determined by radiative effects alone. Rather, the energy balance is one between convective cooling and radiative heating. Ideal gas ascending adiabatically from the continuum forming layers at the base of the photosphere (where the temperature is of the order of 6000 K and the pressure is of the order of 10 kPa) up to the temperature minimum (where, empirically, the temperature is of the order of 4100–4300 K and the pressure is of the order of 0.1 kPa) would have its temperature reduced by a factor of about $10^{2(2/5)} = 10^{0.8} = 6!$ Thus, clearly, significant radiative heating is required to heat gas ascending from the continuum layers up to the temperature minimum.

In the calculation of one-dimensional, plane-parallel model stellar atmospheres, it is possible to take into account line blanketing from hundreds of thousands of spectral lines, either by statistical methods (‘opacity sampling’), or by a reordering of the absorption into monotonic functions of frequency (‘opacity distribution functions’), cf. the review by Carbon (1979). In the 3-D hydrodynamical models, the radiative transfer has to be treated along a large number of rays, for a large number of time steps and both of the above mentioned methods become prohibitively expensive. A further simplification is possible if one assumes that the depth-dependence of the monochromatic opacity is approximately the same for all frequencies. Then, one may classify the opacities into bins (as in the ODF method), but then average over the source function in each bin before solving the transfer equation, Equation (8). Thus, the radiative transfer equation needs to be solved only for a small number of bins (e.g., four) along each ray at any one time in the simulation, which makes the computational cost of the radiative transfer problem affordable (Nordlund, 1982, 1984b).

Though the assumption of frequency independent depth-dependence of the monochromatic opacity is a reasonable first approximation, there are important opacity contributions for which this is not a good approximation. Weak lines of neutral iron are one such case, since neutral iron is abundant in the (cool) upper photosphere, whereas iron is essentially ionized in the lower photosphere. This makes the dept-dependence of the weak iron line absorption radically different from that of, e.g., the $H$-continuum absorption. The absorption of continuum radiation by optically thin iron lines may, in fact, contribute significantly to the heating of the upper photosphere (cf. Nordlund, 1984b, 1985b).
3. Numerical Simulations of Granular Convection

Numerical simulations of granular convection have been performed by myself (Nordlund, 1982, 1984b), by Cloutman (1979), and by Gigas and Steffen (1984). Cloutman used a two-dimensional model, which unfortunately had a horizontal extension too small for a realistic representation of solar granulation. Steffen and Grigas used a model with cylindrical symmetry, with radiative transfer included, and attempted to find stationary solutions. They used a method of characteristics, but did not include any viscous effects (viscosity or shocks). Probably due to the lack of viscous effects, and/or because of the symmetry constraints, they were able to find stationary solutions only for cases where the size of the cylinder was sufficiently small.

My own simulations were performed using a three-dimensional model with $32 \times 32 \times 32$ degrees of freedom per variable, using a semi-spectral representation with a horizontal period of 3 Mm, and a vertical extent of 1.6 Mm. The anelastic approximation was used to filter out pressure waves, and the radiative transport of energy was approximated by the binning of frequencies according to opacity as described above (Section 2.8).

3.1. The Granulation Pattern

The granulation pattern, visible on the Sun on scales of the order of a Mm (cf. Bray et al., 1984), consists of bright granules formed by ascending hot gas, and dark intergranular lanes of descending cool gas, with relatively strong horizontal flows from the granules to the intergranular lanes. The intensity contrast between the hot granules and the cool intergranular lanes is closely linked to the $\Delta H$ necessary to carry the solar flux $F_\odot$ (cf. Equation (21)). This $\Delta H$ corresponds to a $\Delta T$ (peak to peak) of some 5000 K immediately below the visible surface! Most of this temperature contrast is masked by the strong temperature sensitivity of the continuum opacity, which raises (lowers) the surface of continuum optical depth unity in granules (inter-granular lanes).

The amplitude of the visible intensity fluctuation depends on details of, e.g., the temperature dependence of the opacity. Also, in the numerical simulations, the resulting intensity fluctuations may depend on the numerical resolution, and on uncertainties in the absorption coefficients. The tests carried out thus far, however, have not indicated that such uncertainties could influence $\Delta I_{\text{rms}}$ significantly.

The actual intensity fluctuations obtained in the numerical simulations are of the order of 25–30% ($\Delta I_{\text{rms}}$) at 500 nm, and 20–25% at 600 nm (Nordlund, 1984a). This would seem to be in conflict with empirical determinations of $\Delta I_{\text{rms}}$, which generally fall in the range 5–15% at 500 nm (based on actually observed values of typically 2–8%) (cf. the review by Wittman, 1979). These measurements are typically corrected for the influence of ‘seeing’ (atmospheric + instrumental image degradation) by assuming gaussian point spread functions, with a subjectively estimated width. Empirically determined spread functions (Levy, 1971; Deubner and Mattig, 1975; Ricort and Aime 1979; Ricort et al., 1981) have shapes that are far from gaussian. Instead, empirical spread functions – usually one-dimensional (‘line’) spread functions are displayed – have extended wings.
which contain a substantial fraction of the light. The smallest discernable detail ('resolution') is determined by the width of the 'core' part of the spread function, while the degradation of the contrast is determined by the amount of light in the wings of the spread function ('scattering'). It follows that it is not possible to estimate and correct for the degradation of $\Delta I_{\text{rms}}$ from an estimate of the 'resolution' alone. The determination of a spread function must be done independently, for example from the intensity transition across the lunar limb during a partial eclipse (Levy, 1971; Deubner and Mattig, 1975), or else we must await high resolution observations from space (cf. Spruit and van Ballegooijen, 1985).

Deubner and Mattig (1975), who estimated $\Delta I_{\text{rms}} \approx 12.8\%$ at $\lambda = 607$ nm, fitted their observations of the intensity transition across the lunar limb with the sum of two gaussians, a procedure that may severely underestimate the 'wing' component of the spread function. A re-evaluation of the same data, using the sum of two Lorenzians to fit the spread function resulted in a revised $\Delta I_{\text{rms}}$ of at least 20% at $\lambda = 607$ nm (Nordlund, 1984a), largely in agreement with the results of the numerical simulations.

The blueshift and widths of photospheric spectral lines (cf. discussion below) together constitute another, independent measure of the intensity fluctuations in the granulation: since the velocity amplitude determines the spectral line widths (of heavy elements like iron), and the 'product' of velocity and intensity fluctuation determines the spectral line shift, the shift and the width together constitute an indirect measure of the intensity fluctuation. The accuracies of laboratory wavelengths of weak Fe lines are (marginally) sufficient for a comparison of observed and synthetic spectral line shifts (Nordlund, 1980; Dravins et al., 1981; Nordlund, 1984b). The comparison shows good agreement of the blueshifts and widths (cf. Figure 3 and 4 below, and Figure 7(a)–(d) of Dravins et al., 1981). Detailed analysis of limb-effect curves (cf. Nordlund, 1984b, Figure 6.11) could provide another test, although the behavior at the extreme limb is uncertain (both observationally and theoretically).

The geometrical properties of synthetic granulation images obtained from the numerical simulations (Nordlund, 1984b) and granulation images obtained from high resolution observations were compared by Wöhl and Nordlund (1985) and found to be in basic agreement.

3.2. Temperature Fluctuations in the Photosphere

The typical characteristics of the velocity field and the temperature fluctuations in the photosphere are illustrated in Figure 1, which shows contour plots of the temperature and vector plots of the velocity in selected horizontal and vertical planes in the numerical model. The horizontal temperature fluctuation pattern at the $z = 0$ level is a good proxy of the observable intensity fluctuation.

At larger heights, the temperature fluctuation pattern looks very different. This is the result of two competing effects; adiabatic expansion cooling (or compression heating), and radiative energy exchange, which is typically a heating contribution. The combined result is to produce relative temperature minima above small granules (where the ascent velocities are high), and just outside the edges of larger granules (where the horizontal
Fig. 1. Shaded contour plots overlaid with vector plots, showing a snapshot of the temperatures and horizontal velocities in four horizontal planes, from numerical simulations of the solar granulation (after Nordlund, 1984b). The panels show the temperature at heights 300 (a), 200 (b), 100 (c), and 0 km (d). The zero-point of the depth scale corresponds approximately to \( t_{500 \text{ nm}} = 1 \) in a plane parallel model. The contour levels are chosen to span the interval between maximum and minimum temperature with equidistant levels.

velocities are high and, therefore, the pressure is low, and where also the ascent velocities are large). The vertically inclined temperature inhomogeneities resulting from these competing effects leave characteristic signatures in slit-spectra of photospheric spectral lines. Such signatures were observed and correctly interpreted by Evans (1964).

Thus, while the velocity field associated with the granulation penetrates the entire
photosphere, the temperature fluctuation pattern in most of the photosphere bears little resemblance to the granulation pattern. Due to the cooling of gas penetrating into a stable layer, there is a negative correlation with the continuum intensity, but because of the temperature fluctuations associated with the horizontal motion, this negative correlation is weak.

3.3. The sub-surface layers

Below the visible surface, the character of the flow is qualitatively different from that at and above the surface. This is illustrated in Figure 2, which shows the temperatures

![Temperature at depth = 400 km](image1)

![Pressure at depth = 400 km](image2)

![Vertical vel. at depth = 400 km](image3)

![Horizontal vel. at depth = 400 km](image4)

Fig. 2. Shaded contour plots showing snapshots of the temperature (a), pressure (b), and vertical velocity (c) in a horizontal plane at a depth of 400 km. In panel c, the contours corresponding to $u_x > 0$, $u_z = 0$ and $u_x < 0$ are shown full drawn, dashed-dotted, and dashed, respectively. The last panel (d) shows a vector plot of the horizontal velocities, overlaid with the contour lines from panel b.
and horizontal velocities at depth, at one instant in time in the numerical simulations of granular convection (Nordlund, 1984b). The downflow below the surface is concentrated into narrow 'fingers', with unstructured upflow in between. This is similar to the narrow structures that develop in experiments and simulations of Rayleigh–Taylor instabilities. In addition, the strong density stratification aids in concentrating the downflows and dispersing the upflows.

The fluid that goes down into these finger-shaped downdrafts typically has a non-zero angular momentum with respect to the center of the downdraft (as a result of the random positions and shapes of the surrounding granules which supply the inflow). This sets up a circular motion around the center of the downdraft, and the circular velocity is amplified as the downdraft narrows ('bath-tub' or 'inverted tornado'). The centripetal force that balances the acceleration of the fluid in its path around the center of the downdraft is supplied by a negative pressure gradient directed towards the center, corresponding to a pressure deficiency at the center of the downdraft. These pressure minima can become large enough to cause difficulties in the numerical simulations. These 'inverted tornados' are interesting from a purely hydrodynamics point of view, but the fact that any vertical magnetic field lines in the surrounding photosphere must be carried towards, and 'sucked into', these downdrafts also makes the phenomenon potentially very important as a source of hydromagnetic disturbances.

3.5. SYNTHETIC SPECTRAL LINE PROFILES

Time and area averaged synthetic spectral line profiles provide an important diagnostic for granular convection, particularly since such profiles do not require high spatial resolution. Indeed, using full disc profiles, comparisons of synthetic full disc spectral line profiles with observed stellar spectral line profiles are possible (Dravins, 1982; Dravins and Lind, 1984; Nordlund and Dravins, 1986).

Figure 3 (from Nordlund, 1984b) shows a comparison of synthetic and observed solar center disc spectral line profiles for four Fe I lines. Accurate oscillator strengths from the Oxford group (Blackwell et al., 1976, 1979a, b, 1982) were used, and the iron abundance was calibrated on the weakest line. The iron ionization equilibrium was calculated from the radiation intensities at the most important ionization edges of Fe, using data from Lites (1972; see also Athay and Lites, 1972), and scaled \( H^- \) absorption coefficients, calibrated on the observed solar disc center intensities in the blue and the near UV.

The comparison shows basic agreement, but with a noticeable discrepancy in the cores of the stronger lines. This discrepancy is probably due to too low temperatures in the uppermost layers of the model caused by approximations in the treatment of spectral line blanketing effects (cf. Nordlund, 1984b, Section 7, and Nordlund, 1985b). On the other hand, a quite cool component of the upper photosphere is necessary to account for the strength of the infrared bands of carbon monoxide (Ayres and Testerman, 1981, Ayres et al., 1985), and Fe I line cores may be substantially brightened by velocity induced non-LTE effects (Nordlund, 1985a).

Figure 4 shows the spectral line bisectors corresponding to the line profiles in
Fig. 3. Synthetic (dash-dotted) and observed (full drawn) Fe I spectral lines (from Nordlund, 1984b). The observed line profiles are from the Jungfraujoch atlas (Delbouille et al., 1973), with line center positions derived from the Kitt Peak Tables (Pierce and Breckinridge, 1973). The synthetic spectral lines are averages over a 110 solar min simulation sequence. Accurate $g_f$-values from Blackwell et al. (1976, 1979a, b, 1982) were used, and the iron abundance was chosen to fit the central depth of the weakest line.

Figure 3. The bisectors of the synthetic spectral lines fit the bisectors of the observed lines remarkably well, especially after the slight broadening due to the 5-min oscillations has been taken into account (this changes the asymmetry of the synthetic spectral lines even if one assumes that the broadening itself is symmetric, because the unbroadened spectral line is asymmetric).
4. Convection on Global and Intermediate Scales

4.1. Global Convection

Numerical simulations of global convection have been performed quite extensively. Non-linear convection in a spherical shell has been simulated in the Boussinesq approximation by Durney (1970), Young (1974), Gilman (1977, 1978a, b), and Marcus (1979, 1980a, b). Similar models with magnetic fields included were calculated by Gilman and Miller (1981) and Gilman (1983b). Linear and second-order non-linear stratified models in the anelastic approximation have been calculated by Glatzmaier and
Gilman (1981a, b), while Glatzmaier (1984, 1985a, b) and Gilman and Miller (1985) have calculated fully non-linear stratified models in the anelastic approximation, with dynamically active magnetic fields included. The earlier work has been reviewed by Gilman (1980, 1983a); here I will just make a few remarks on the most recent work.

The calculations by Glatzmaier and Gilman represent a logical evolution of increasingly realistic models of global solar convection. The most recent of these models have a differential rotation in basic agreement with the observed one, and have meridional circulations and pole-equator temperature differences consistent with observational upper limits. These results do not involve any tuning of free parameters. The few free parameters that remain are related to the numerical methods used, in particular to what happens on the sub-grid scale. For sufficiently high numerical resolution, these parameters do not influence the character of the results greatly, as long as physically reasonable values are chosen.

Substantial disagreement with observations remains only (for the velocity field) in the amplitudes of the horizontal velocities near the upper boundary (which however is several Mm below the visible surface) and more indirectly, with the time dependence of the generated magnetic fields, which propagate away from the equator with time.

The differential rotation in these models tends to be constant on cylinder surfaces (cf. Glatzmaier, 1984, Figure 5; 1985a, Figure 2(a)). This is something that may be expected from an extension of the Taylor-Proudman theorem to a stratified system (Busse, 1977). The model results may or may not be consistent with recent Fourier Tachometer observations of the rotational splitting of solar \(p\)-mode oscillations (Brown, 1985). The preliminary analysis by Brown indicates that the differential rotation decreases with depth (which is qualitatively consistent with the hydrodynamical models), but also that the rotation rate at depth becomes more nearly equal to the equatorial rotation rate (which does not seem to be consistent with the models). Should the empirically determined rotation have a depth dependence that is qualitatively different from that of 'constancy on cylinders', then either some essential ingredient in the physics must be missing in the models, or else the yet not included surface layers play an unexpectedly important role.

Since the most prominent disagreement between the behavior of the numerical models and the Sun is the failure of the numerical models to predict magnetic cycles of the right period, and with the right sense of magnetic fields propagation, further discussion of the numerical models of global convection is deferred to Section 5.3.

4.2. Supergranulation and Mesogranulation

The intermediate scales (3–10 Mm), with supergranulation (Leighton et al., 1962; Simon and Leighton, 1964; Worden and Simon, 1976) and mesogranulation (November et al., 1981, 1982) have received little attention in numerical work, compared with granular convection and global convection. This may be partially due to the difficulties of describing the interaction of convection on a supergranular scale with granular convection and radiation at the solar surface. In view of the very small limits on any horizontal temperature variation associated with supergranulation, and the results of
granulation simulations showing that (on the time-scale of supergranulation) the energy is removed at an approximately constant rate in a layer at, say, 1 Mm below the surface, it would seem possible to use a 'constant heat flux' upper boundary condition for numerical simulations of convection on a horizontal scale of 3–100 Mm, over a depth range covering perhaps 25–50% of the solar convection zone.

The term 'mesogranulation' was introduced by November et al. (1981, 1982) for a velocity and temperature fluctuation field observed on the scale of 5–10 Mm, with a characteristic time-scale of an hour or so. The apparent velocity magnitude is of the order of 100 m s\(^{-1}\). The true character of these fluctuations is somewhat unclear, since apparent velocity shifts and temperature fluctuations of this order of magnitude, and with this characteristic time scale are produced also by a modulation of the 'vigor' of granulation, even though there is no net mass flux involved (cf. Dravins et al., 1981, Figures 11(a)–(b)). The separation of a real velocity/temperature fluctuation field from this apparent one is a difficult observational problem. However, with superb spatial resolution, it might be possible to observationally determine, and correct for, the apparent fluctuations caused by the slow modulation of granular convection.

5. Interaction with Magnetic Fields

The observation of magnetic fields on the Sun and interpretation of their structure is reviewed elsewhere in this volume (Stenflo, Zwaan, Muller). For reviews on magnetoconvection, see Proctor and Weiss (1982) and Nordlund (1984c). Dynamo theory has recently been reviewed by (e.g.) Cowling (1981), Gilman (1983a), and Schüssler (1983). Spruit (1983, 1984a, b) has reviewed the magnetic flux tube picture of solar magnetic fields. However, since the interaction of magnetic fields with convection is crucial in determining the structure of the solar magnetic fields, some mention of magnetic fields is appropriate at this point.

5.1. Interaction with Granular Convection

The 'convective collapse instability' as a mechanism for flux concentration has been discussed extensively (Nordlund, 1976; Parker, 1978; Webb and Roberts, 1978; Spruit and Zweibel, 1979; Spruit, 1979; Unno and Ando, 1979; Hasan, 1983, 1984, 1985; see also Parker, 1979, Section 10 and references therein).

These investigations qualitatively illustrate how a superadiabatic (convectively unstable) mean stratification leads to a 'collapse' of an initially weak field into a strong flux concentration. Proctor (private communication) has pointed out that, rather than being the effect of an instability, this is the quenching of a (somewhat modified) convective instability: when the magnetic field becomes sufficiently concentrated, it stops something that looks very much like a normal convective instability.

The 'convective collapse' carries a weak initial magnetic field into a configuration where the interior of the magnetic flux structure has a temperature that is lower than the mean surroundings over several pressure scale heights, and where, consequently, the gas pressure is much reduced inside the flux concentration, relative to outside.
(Parker, 1955; Spruit, 1976). The gas pressure difference is then only slightly smaller than the external pressure and can support a magnetic field pressure almost corresponding to the surrounding gas pressure. In the Sun, the gas pressure is about 13 kPa at the level where the continuum optical depth is unity, and this gas pressure corresponds to a magnetic flux density of some 0.18 mT (1.8 kG). Thus, if the evacuation is significant down to this level, but not much further, one gets flux densities which are similar to the observed ones.

Most of the investigations quoted above have assumed adiabatic perturbations (no radiative transfer effects). There has also been some confusion about the choice of boundary conditions, and one must conclude that the balance of effects that determines the actual degree of evacuation (and, hence, the observable field strength) in the small scale solar flux concentrations is still not well understood (cf. Nordlund, 1984c, Section 4).

Numerical simulations of the interaction of convection and magnetic fields in strongly stratified models have been performed by myself (Nordlund, 1983, 1986), and by Hurlburt and Toomre (1985).

The calculations of Hurlburt and Toomre are two-dimensional, with a polytropic model covering 2.4 density scale heights, fully compressible, but are not directly relevant to the solar case, since no description of radiative transfer effects were included. Concentration of magnetic flux towards the downdrafts occurs initially because of kinematic effects, but is further enhanced because of the superadiabatic stratification. Since the temperature gradient excess relative to an adiabatic stratification is not very large, the temperature deficiency in the interior of the flux concentrations is not big enough to produce a strong evacuation of the flux structures.

My own calculations are extensions of the numerical simulations of solar granulation, complemented with appropriate equations for the magnetic field. Thus, these calculations are three-dimensional, anelastic, cover about 5.5 density scale heights, and include a realistic description of radiative transfer effects. The behavior of the magnetic flux concentrations obtained in these simulations is in qualitative agreement with the observed properties of small scale flux concentrations of the Sun (Muller, 1983). The magnetic flux, which is unipolar with a conserved average vertical component, is concentrated in the inter-granular lanes, where the presence of the magnetic field helps to reduce the temperature (by screening off the flow, and hence allowing the radiative cooling at the surface to cool the interior further). The resulting gas pressure difference is large enough to concentrate the magnetic field to flux densities of 0.1–0.15 mT (1–1.5 kG) at the observable level. The volume in-between the magnetic flux concentrations are swept free of magnetic field, and here convection proceeds to produce new granules and new intergranular lanes. The magnetic flux concentrations ‘creep’ into the newly formed downdrafts, since the gas pressure in the downdrafts is insufficient to balance the magnetic pressure in the flux concentrations.

Thus, these simulations are capable of illustrating the qualitative properties of the interaction of granular convection and unipolar magnetic fields. However, the numerical resolution (50 km vertically, 190 km horizontally) is insufficient to realistically model the
steep walls of the flux concentrations, and the numerical diffusion of plasma through the magnetic field is a significant factor in the determination of the internal temperature structure of the flux concentrations in these simulations.

5.2. Magnetic Field Topology

Numerical simulations, although still with insufficient numerical resolution, provide an important basis for a qualitative understanding of the interaction of convection, magnetic fields, and radiation. For one thing, they serve as a reminder of the complicated topology of the magnetic field lines in the solar photosphere. This may be a crucial factor, both in the determination of (quasi-)equilibrium field strengths in the flux concentrations, and for the role of the magnetic field as a 'channel' for mechanical energy to the solar chromosphere and corona.

Torsional disturbances due to the 'bath-tub' effect mentioned above, or 'buffeting' by the granulation, are possible sources for 'sudden disturbances' in the photosphere which, according to Hollweg et al., (1982) and Hollweg (1982), may produce spicule-like ejections into the chromosphere. In this connection it should be noted that the horizontal velocities and, especially, the relative pressure fluctuations obtained in the numerical simulations of solar granulation are significantly larger than the values assumed by Spruit (1984a) in his discussion of tube mode generation ($u_{\text{rms}} \approx 3 \text{ km s}^{-1}$ and $\delta P/P \approx 0.5$, rather than $2 \text{ km s}^{-1}$ and 0.05, respectively).

The finite lifetime of the individual flux concentrations, and the strong horizontal velocity field present, cause the field lines in the photosphere to wander around in a random-walk fashion. Such field line motions have been discussed by Parker (1983) as a possible heating mechanism ('topological heating') for the quiet solar corona. The actual energy dissipation due to the foot point motion depends on details of the dissipation mechanism and is hard to estimate quantitatively, as discussed by Spruit (1984b).

5.3. Solar Dynamo Models

The failure of even the best self-consistent models so far to produce the observed equatorward drift of magnetic field with cycle phase (Gilman, 1983; Glatzmaier, 1984, 1985a, b; Gilman and Miller, 1985; cf. also the discussion by Gilman, 1984), is telling us that something essential is still missing in the way these models describe the interaction of convection and magnetic fields. (The apparent success of kinematic models (e.g., Yoshimura 1978a, b, 1983) results from the abundance of free parameters controlling the behavior in these models; e.g., the freedom to choose the depth dependence of the rotation profile.) Since the same physical situation (magneto-convection in a rotating spherical shell) has been modeled with several independent numerical methods, with largely the same results, we can be quite confident that – within the given assumptions of these numerical models – the basic magnetohydrodynamics has been correctly modeled, and that the remaining disagreement with the behavior of the solar dynamo is not due to any peculiarity of the numerical methods used.

Glatzmaier's models were calculated with a semi-spectral method, where the ($\theta, \phi$) dependence is represented with spherical harmonics $Y_{ijm}(\theta, \phi)$ and the radial dependence
is represented by Chebyshev polynomials $T_n(r)$. These calculations (with $17 \times 1024$ degrees of freedom per variable) are close to the limits of what is practical with present day computers. Nevertheless, it would be highly desirable to either increase the number of degrees of freedom per variable horizontally, or else use the available degrees of freedom in such a way as to increase the horizontal resolution. This may perhaps be achieved by using a representation where the $\phi$ dependence has a period equal to $2\pi/M$, where $M$ is larger than 1. The need for increased horizontal resolution is clear if we consider the ‘aspect ratio’

$$
\frac{2\pi R/L}{0.43R/N} \approx 15N/L \approx 7.5.
$$

(24)

Of course, this is not a true aspect ratio (the $\Delta r$ spacing varies with $r$, from a maximum of the order of $(\pi/2)0.43\, R/N$, to a minimum which is much smaller), but the ratio (24) serves to illustrate the relatively poor horizontal resolution compared to the vertical one. The broad maximum in the kinetic energy spectrum around $l = 15$ has only a factor 2 margin to the maximum horizontal wavenumber (31); i.e., the dominant scales modeled are only twice as big as the smallest ones that could be represented. Thus, for example, up/down asymmetries of the energy carrying cells can not be well represented in the model, since the harmonics of the dominating modes lie outside the wavenumber limits of the model. As discussed in Section 2.4, up/down asymmetries are a prominent feature of resolved compressible convection, and the failure to resolve such asymmetries (and corresponding asymmetries in the magnetic field) may be one reason for the shortcomings of present models of global convection and dynamo action. The separation of convective motions and magnetic fields (‘flux expulsion’), e.g., Galloway and Wiess, 1981) found in well resolved models of magnetoconvection may be an important factor in determining the period of the solar cycle, since such a separation would be expected to decrease the coupling between the convective motions and the magnetic field. Such a separation can only be modeled if convective cells are well resolved.

The Chebyshev representation automatically produces an increase in the numerical resolution close to the boundaries, which is very desirable, since the scale height decreases rapidly at the upper boundary, and boundary effects near the transition from unstable to stable stratification at the bottom need to be resolved. This, however, requires a correspondingly increased horizontal resolution. Thus, ideally, the code should be able to (at least marginally) resolve supergranulation at wavenumbers of 100 to 200.

Several suggestions as to what else might be missing in the modeling have been put forward: the suggestion by Galloway and Weiss (1981) and by Golub et al. (1981) that the solar dynamo might be operating at the base of the convection zone, in the transition zone between the stable interior and the turbulent convection zone, was tested by Glatzmaier (1982b) by including a stable region below the convection zone, and by decoupling the velocity field and the magnetic field in the bulk of the convection zone, to simulate magnetic flux structures that are too small and too concentrated to be
affected by the helicity in the convection zone. Only part of a cycle has been simulated, and the results are not conclusive.

Glatzmaier also suggested that the solar dynamo could be operating in the outer 5% or so of the solar radius, where angular velocity is probably decreasing with depth, and where the helicity has the correct sign for producing equatorward drift of the magnetic field. Considerations of magnetic buoyancy (Parker, 1975) suggested that a magnetic cycle with a period as long as 22 years may be hard to achieve if the dynamo is not deep seated, but these considerations neglected the possibly important role of radiative cooling of flux structure interiors at the solar surface.

6. Conclusions and Recommendations for Future Work

Solar convection, and the interaction of convection and magnetic fields on the Sun serves as an important test for our understanding of the dynamics of the (magneto-) hydrodynamics of astrophysical plasmas in general. From the interaction of large-scale convection and magnetic fields in the solar interior, through the interaction of convection, radiation, and magnetic fields in the photosphere, to the plasma physics of the chromosphere and corona, there is a wealth of physical phenomena, likely to have counterparts in non-solar astrophysical systems we cannot resolve observationally. For this reason, improvements in our understanding of the dynamics and energy transport in the solar plasma is of great importance.

I would like to conclude this review with some hopes and recommendations as to where we should push for progress in research on solar convection in the near future:

- The numerical simulations of global convection and dynamo action in the solar convection zone need further development, to resolve the discrepancy in the temporal behavior of the magnetic fields generated by the convective dynamo. Part of such a development should go in the direction of improved numerical resolution, to allow a better representation of processes on small scales (close to the upper and lower boundaries). This may perhaps be feasible with further improvement of the numerical methods, or else must await the next generation of gigaflop super-computers.

- ‘Cartoon’ models of the solar dynamo/solar activity phenomenon (Schüssler, 1980, 1983; Van Ballegoijen, 1982a, b) should be developed further. ‘Cartoons’ may provide important inspiration for improved numerical models, and/or may suggest critical observational tests.

- There is a need for numerical simulations of convection on the scale of supergranulation, both to model supergranulation itself, and to model its interaction with magnetic fields on the scale of active regions. This is probably feasible with present day computers and numerical methods.

- Our understanding of granular convection, and its direct and indirect influence on the solar photosphere, chromosphere, and corona needs improvement both on the theoretical and observational sides. As regards the chromosphere and the corona, we need a better understanding of how the energy is transferred from motions in the convection zone (on all scales), through stresses in the magnetic field that extends out
into the chromosphere and corona, where dissipative processes and magnetic field instabilities convert the energy into heat and non-thermal particle energy. As regards the photosphere, we need a better understanding of the influence of spectral lines on the energy balance in the optically layers. Also, one would like to see fully compressible codes developed, ideally covering a sufficient depth range to study the interaction of the 5-minute oscillations with convection.

- Observationally, we may look forward to the Solar Optical Telescope, which will be able to resolve solar granular convection and its interaction with small scale magnetic fields in great detail (cf. Spruit and van Ballegooijen, 1985). Apart from the direct advantage of improved spatial resolution, a great advantage of observations from space as compared to ground based observations is the temporal stability of the picture degradation, and the consequent possibility to obtain long time series of observations with a well determined spread function.

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