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3. Atmospheric and Oceanic Boundary Layer Physics

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3.1 Introduction

The globe of the earth is surrounded by a gaseous atmosphere which is always in motion. When in contact with the land or the water surface of the earth the flow is reduced to zero, relative to the underlying surface, and it is this boundary flow that interests us here. As well as the planetary boundary layer in the air, also known as the Ekman layer, there is an oceanic boundary layer which interacts with the air above. An adequate description of physical processes and mechanisms that determine the structure of the interacting atmospheric and oceanic boundary layers as well as a theoretical background is needed for developing parameterization schemes. The more general features of this problem are treated in the monograph by Kraus & Businger (1994).

One of the most important problems is the parameterization of the turbulent fluxes of momentum, latent and sensible heat at the sea surface. The oceans are the major source of atmospheric water and a major contributor to the heat content of the atmosphere. Most of the solar energy is absorbed by the oceans, and this energy becomes available to maintain the atmospheric circulation only through turbulent fluxes of latent and sensible heat. Radiative, sensible and latent fluxes determine the ocean surface energy flux and, consequently, the vertical structure of the upper ocean. On average, surface buoyancy fluxes are stabilizing over vast oceanic areas (Gargett 1989). Momentum fluxes act as a drag on the atmospheric motions and induce the so-called "wind-driven" component of the ocean currents which produce considerable horizontal transport of energy and momentum. When surface buoyancy fluxes are destabilizing, ie. a heavier fluid overlies a lighter one, convection in the upper part of the ocean can occur. This process can lead to the sinking of near-surface water to relatively deep depths.

The turbulent exchange processes at the air-sea interface are strongly influenced by the state of the sea surface which is varying in time. However, the sea surface temperature changes little over a diurnal cycle because water has a large heat capacity. The roughness of the sea surface depends on the atmospheric surface layer parameters and consequently, on the processes in the whole atmospheric boundary layer. It is believed that the sea surface roughness governs the degree of wind drag time variability (Kitaigorodskii 1973). It is very important that the atmospheric motions generate surface waves which also contribute to a turbulent mixing of the atmosphere and ocean.

Thus, one can recognize that the wind stress over the ocean, which is the main interest of this monograph, is not an isolated characteristic of the sea surface state but rather an indicator of the coupled atmospheric and oceanic boundary layer dynamics. It is outside the scope of this chapter to consider the dynamic interaction of the atmospheric boundary layer with the troposphere, and oceanic boundary layer with deep ocean, and so it is assumed that all the necessary characteristics are known at the outer boundaries of the corresponding boundary layers.

3.2 Marine Atmospheric Boundary Layer

3.2.1 Vertical structure

In general, the atmospheric boundary layer may be more or less arbitrarily split into two regions: the region immediately adjacent to the air-sea interface which is called the constant-flux layer (Monin & Yaglom 1971), and a free-atmosphere-topped interfacial (or "transition") layer over it. Formally, the former is defined as the layer where vertical variations of the turbulent fluxes do not exceed, say, 10 % of their surface values. Typically, this layer is about 10 - 100 m thick which makes approximately 10 - 20 % of the overall boundary-layer thickness. Vertical distributions of meteorological parameters show here the logarithmic asymptotics when approaching the ocean surface and depend on the air density stratification. Small scale turbulent eddies with the sizes restricted by the distance to the underlying surface are mainly responsible here for the momentum, moisture and heat transport. In this layer also waves dominate the air flow and become important for the transport of momentum. Very close to the surface one can detect a thin microlayer with the height of order 1 cm (the so-called viscous sublayer), in which molecular processes dominate.

Over the subtropical and tropical parts of oceans the atmospheric boundary layer is convective throughout nearly the whole year. The surface density flux due to heating and moistening is directed downwards; the potential temperature and the specific humidity decrease across the constant flux layer. Based on experimental and model studies one may schematically describe the structure of the interfacial layer in this conditions as consisting of four layers, each governed by different physics (Augstein 1976). Above the constant flux layer there is a mixed layer with the thickness of order 1 km. The change of potential temperature with height is small here, and mixing is dominated by convectively-driven organized motions (large eddies). Atop the mixed layer one can detect the so-called entrainment zone which is 100 - 500 m thick. In this layer turbulence is intermittent, the air stratification is stable with regard to potential temperature, and internal waves and sometimes small clouds are observed. A cloud-topped boundary layer includes an additional layer of broken or uniform clouds. This cloud layer connects to the free atmosphere via an inversion layer.

The most variable part of this idealized structure is the cloud layer. If no clouds are formed, then the atmospheric boundary layer terminates at the entrainment zone. When very deep clouds are developed and extended through the whole troposphere (for example, in cyclonic conditions of surface convergence and large-scale upward motion), the top of the boundary layer is not defined. However, in many cases the free atmosphere is turbulently decoupled from the boundary layer, typically at the ocean on the rear side of depressions and in the ITCZ.

In middle and high latitudes, the balance between the pressure gradient, Coriolis and turbulent stress divergence determines mainly the structure of the boundary layer. The special but not too rarely observed case of a steady state, horizontally homogeneous, neutrally stratified, barotropic atmosphere with the stress represented by the independent of height eddy diffusivity results in the well-known Ekman spiral wind profile (Brown, 1974). The characteristic feature of the Ekman profile is that due to friction, winds in the boundary layer cross the isobars from high towards low pressure. In the case of clow or high pressure system, the cross-isobaric component of flow induces upward or downward vertical motions, respectively. Such a process known as Ekman pumping (Stull 1988) is very important for linking the boundary layer with the free atmosphere.

The one - dimensional representation of the boundary layer structure was found to be useful in many cases, but it may become incorrect when horizontal advection is dominant, in particular, in the vicinity of oceanic fronts. Observations have shown (Guymer *et al.* 1983; Khalsa & Greenhut 1989; Rogers 1989) that the spatial variability of the sea surface temperature on a scale of 200 km or less causes horizontal variability on similar scales in the atmosphere. It was also found that the turbulent structure of the marine atmospheric boundary layer has different scales on opposite sides of a sea surface temperature discontinuity. Effects of oceanic fronts on the wind stress are discussed by Gulev & Tonkacheev (1995).

3.2.2 Turbulence

Since Reynolds numbers for atmospheric motions are very large (of order 10^7), turbulence in the boundary layer is fully developed and three-dimensional. The vertical turbulent transport of momentum, heat and moisture is the main process which links the large scale motions in free atmosphere with the surface. Small-scale turbulence consists a set of disturbances with scales which do not exceed the distance to the surface. Besides the classical descriptions of turbulent flows which are connected with the names Reynolds and Taylor, a new approach was introduced by the discovery of ordered motions in many turbulent shear flows. It has been understood that traditional parameterization schemes do not describe some essential features of the atmospheric boundary layer, in particular, non-local nature of turbulence caused by the presence of coherent structures (large eddies). Following Stull (1991), one can define large eddies as turbulent structures with the size of the same order as that of planetary boundary layer, or of the same order as that of mean flow. Spectral decomposition of turbulent fields (see, for example, Pennel & LeMone 1974) indicated that large eddies contain most of the turbulent kinetic energy. Examples of these coherent structures (convective thermals of the same 1- to 2-km diameter as the mixed-layer depth, well-ordered roll vortices, mechanical eddies of the same size as the 100-m-thick shear region of the surface layer and convective plumes) are given in reviews by Mikhailova & Ordanovich (1991), Etling & Brown (1993) and Stull (1993). Byzova *et al.* (1989) have also discussed the turbulent cell convection with quasi-ordered convective structures of the 3 to 10 km size and small-scale cell convection with coherent motions of a few hundred meters. All these structures coexist with small-scale turbulence.

In particular, it was discovered (Wyngaard & Brost 1984) that the vertical diffusion of a dynamically passive conservative scalar through the convective boundary layer is the superposition of two processes. These processes are driven by the surface flux, "bottom-up" diffusion, and by the entrainment flux, "top-down" diffusion. It was also found that the vertical asymmetry in the buoyant production of the turbulent kinetic energy caused the top-down and bottom-up eddy diffusivities to differ.

At the very small scale, other coherent motions play an important role in the surface layer near the wall. Many laboratory experiments (e.g. Kline *et al.* 1967; Corino & Brodkey 1969; Narahari Rao *et al.* 1971; Brown & Roshko 1974) have demonstrated the significance of so-called bursting processes for turbulence production. Following Narasimha (1988), bursting processes may be considered as coherent, quasi-periodic cycles of events. These events can be described as a retardation of the near-surface fluid in forms of streaks, a build up of a shear layer leading to a violent ejection of fluid from the surface and a sweep of faster fluid from the outer layer towards the surface. It was found by Corino & Brodkey (1969) that nearly 70 % of the shear stress could be due to the ejection process. Sweep events play a major role in the bedload transport in rivers and oceans (Heathershaw 1974; Drake *et al.* 1988). It was shown (Wallace *et al.* 1972) that an interaction between ejections and sweeps accounts for a substantial part of the momentum flux to balance its excess produced by motions of these two categories. Intermittent coherent motions have been detected in the turbulence measurements as periodic, large-amplitude excursions of turbulent quantities from their means. It was discovered that this process strongly influences the turbulent transport through cycles of ejections and sweeps also within atmospheric surface layer (Narasimha 1988; Mahrt & Gibson 1992; Collineau & Brunet 1993) and, in particular, in the marine boundary layer (Antonia & Chambers 1978).

However, it seems that turbulence production close to the surface is a process which is nearly independent of large-scale processes in the outer layer (Kline & Robinson 1989). Most of the turbulence and most of the momentum flux is produced by the near-wall process, and this phenomen is usually called "active turbulence" (Townsend 1961). He has also suggested the concept of "inactive turbulence", according to which turbulence is of the boundary layer size scale and does not produce the momentum flux at the surface. Such frictional decoupling has often been observed at the sea surface (e.g. Volkov 1970; Makova 1975; Chambers & Antonia 1981; Smedman *et al.* 1994). It was derived from these investigations that momentum transfer from the decaying surface waves to the atmosphere ("high wave age " conditions) can be suggested as the possible mechanism causing the frictional decoupling. If by that time the surface buoyancy flux is relatively small, and the relatively large wind shear is maintained in the upper part of the boundary layer (for instance, due to the development of a low-level jet), then turbulent energy produced in this region can be brought down, including the surface layer, by pressure transport (Smedman *et al.* 1994). This imported from above turbulence is an example of turbulence of "inactive" kind. It was found that the turbulence statistics of the boundary layer resemble in these conditions those of a convective boundary layer but with different scaling, since the buoyancy production is small. Even if these conditions are not very typical for the ocean, the frictional decoupling mechanism must be taken into account when the wind stress parameterizations are developed for the use, for example, in global climate models.

Additionally, drag reduction in flows with suspended particles is another well known phenomenon (Toms 1948). It is recognized now that a major dynamical effect of suspended fine particles is the stabilization of the secondary inflexional instability, the suppression of the intense small-scale turbulence and the decrease of the turbulence production (Lumley & Kubo, 1984; Aubry *et al.* 1988). It was also found that particles influence the flow turbulence in two contrary ways - by the expense of the energy of turbulent fluctuations to suspend particles, and by destabilization of the flow, when a significant slip between the two phases exists. In the case of coarse particles such a flow destabilization is the main factor, leading to the additional turbulence production, since each particle sheds a wake disturbance to the flow like a turbulence-generating grid (Hinze 1972; Tsuji & Morikawa 1982).

Over oceans, at wind speeds above, say, 15 ms^{-1} intensive spray is detected near the air-sea interface. Spray droplets, detached from the sea surface, carry their instantaneous momentum with them. Biggest of these droplets fall down into the sea and return their momentum back to the sea surface. On their way through the air, they may interact with the air and exchange momentum with the air. This process seems to be not very important, since spray is generated mainly from the larger waves which have an orbital velocity near to the wind velocity. However, the small bubble-derived water droplets and salt particles can be suspended in the air flow, and carried by turbulence higher up to cloud heights (de Leeuw 1986). If their concentration is large enough, the density stratification can be remarkably altered and hence this can influence the momentum transfer. At the present time, it is not clear, how important this mechanism could be for the observed wind stress over the ocean. Due to additional evaporation from the spray droplets, a modification of the temperature and humidity gradients might be also expected.

There are studies (e.g. Lumley 1967; Shaw & Businger 1985; Narasimha & Kailas 1987; Narasimha 1988; Mahrt & Gibson 1992; Collineau & Brunet 1993), in which the intermittent nature of the near-surface turbulence is discussed with respect to the atmospheric boundary layer. Its possible connection with drag reduction phenomena due to the presence of suspended particles in the air (in particular, sea spray) seems to be also important. An expansion of the Monin-Obukhov similarity theory for the case of a two-phase flow (e.g. Barenblatt & Golitsyn 1974; Wamser & Lykossov, 1995) allows to describe generally the drag reduction due to more stable density stratification, but without regarding details caused by the intermittent nature of turbulence and by the non-regular loading of particles into air.

3.2.3 Turbulence closure

Let a be any meterological variable: horizontal components of the wind velocity (u and v for the alongwind and crosswind directions, respectively), potential temperature (θ), specific humidity (q), etc. The Reynolds equation for conservative statistically averaged quantity reads

$$\frac{\partial a}{\partial t} = -\frac{\partial \overline{a'w'}}{\partial z} + (\cdots), \tag{3.1}$$

where t is time, z is the vertical coordinate, and w is the vertical component of the wind velocity. The overbar denotes the average values, primes stand for the turbulent fluctuations, and the terms responsible for the contribution of the other (non-turbulent) processes into dynamics of the boundary layer are marked by dots. For simplicity in notation, the bar signifying mean values is omitted everywhere except for the notation of the turbulent covariances.

Under the assumption of horizontal homogeneity, the budget equation for the turbulent kinetic energy (e) may be written as follows (Monin & Yaglom 1971):

$$\frac{\partial e}{\partial t} = -\left(\overline{u'w'}\frac{\partial u}{\partial z} + \overline{v'w'}\frac{\partial v}{\partial z}\right) - \frac{g}{\rho}\overline{\rho' w'} - \frac{\partial \overline{w'e}}{\partial z} - \frac{1}{\rho}\frac{\partial \overline{p'w'}}{\partial z} - \epsilon, \tag{3.2}$$

where g is the acceleration due to gravity, ρ is the air density, p is the pressure, ϵ is the dissipation rate, and $e = \overline{(u'^2 + v'^2 + w'^2)}/2$. The terms on the right-hand side of this equation describe, in order, the shear production, the buoyancy production/destruction, the vertical turbulent transport, the pressure transport, and the viscous dissipation. Mainly, there are two sources of turbulence in the atmospheric boundary layer: the wind shear and the buoyancy flux near the surface, and the wind shear at the top of the boundary layer. The frequent presence of clouds in the marine boundary layer leads to strong change of the radiation budget at the surface. In this case there is also an additional production of turbulence due to long-wave cooling at the top of the cloud layer, short-wave heating of the inner part of clouds and phase changes of the water.

Most of parameterization schemes are based on the K-theory stability-dependent eddy diffusivity closure, which assumes that fluxes are associated with small-size eddies only in a manner similar to molecular diffusion and that the static stability is estimated on the basis of the local lapse rate. The turbulent fluxes in the boundary layer are calculated in this case as suggested by the Boussinesq (1877) hypothesis [1]

$$\overline{a'w'} = -K_a \frac{\partial a}{\partial z}, \tag{3.3}$$

where the eddy diffusivity K_a, having a positive value by its physical sense, is evaluated by means of the mixing length theory (Monin & Yaglom 1971):

$$K_a = \alpha_a l^2 \left|\frac{\partial \vec{V}}{\partial z}\right| F_a(Ri). \tag{3.4}$$

Here α_a is an "universal" constant, l is the integral turbulence scale, and F_a is an "universal" non-dimensional function, depending on the gradient Richardson number,

[1] Strongly speeking, Boussinesq has applied this approach to the turbulent transport of momentum. However, it was found that such a closure can be also used for the description of the turbulent transport of any passive scalar.

$$Ri = \lambda \frac{\partial \theta_v / \partial z}{|\partial \vec{V} / \partial z|^2},$$

where $\theta_v = (1 + 0.61q)\theta$ is the virtual potential temperature, $\vec{V}$ is the horizontal wind velocity vector with the components u and v, $\lambda = g/\theta_v$ is the buoyancy parameter. To calculate l, the following formula, suggested by Blackadar (1962), can be used

$$l = \frac{\kappa z}{1 + \kappa z / l_\infty}, \tag{3.5}$$

where $\kappa = 0.4$ is von Kármán's constant, and l_∞ is some function of the external parameters (see Holt and SethuRaman 1988; for the review). The eddy diffusivity coefficients K_a are also often related to the (mean) turbulent kinetic energy and dissipation rate (Monin & Yaglom 1971)

$$K = le^{1/2} = Ce^2/\epsilon, \qquad K_a = \alpha_a K, \tag{3.6}$$

where C is an universal constant.

At the same time numerous experimental data (e.g. Budyko & Yudin 1946; Priestley & Swinbank 1947; Deardorff 1966) showed that sometimes the atmospheric boundary layer is neutrally or weakly stably stratified (for example, the convective mixed layer), but the heat flux is directed upward. This corresponds to the negative heat diffusivity which is not consistent with the basic molecular diffusion analogy. Moreover, the verical transport of the turbulent kinetic energy is directed upward throughout the whole depth of the convective layer (see, for example, Andre 1976; Kurbatskii 1988).

The K-theory has also essential shortcomings in application to the jet-like flows. For example, in the case of a jet which was observed in the fair-weather trade wind boundary layer (Pennel and LeMone, 1974) a countergradient transport of momentum was found. Other examples of such phenomenon for the channel flows are given by Narasimha (1984), Yoshizawa (1984) and Kurbatskii (1988). In such flows the point of maximum velocity and that of zero momentum flux do not coincide. Thus, there is a region where positive momentum flux is accompanied by positive velocity gradient. This means that momentum is transported up from this region, but not down to the surface, as it is expected for the situation without jet. Persistent countergradient fluxes of momentum and heat (density) have also been observed in homogeneous turbulence forced by shear and stratification: at large scales when stratification is strong, and at small scales, independently of stratification (Gerz & Schumann 1996).

At the present time, a number of reviews of the present state-of-the-art in the turbulence closure problem for the boundary layer with coherent structures is published (e.g. Stull 1993, 1994; Lykossov 1995). An hierarchical description is usually used to classify numereous approaches presented in the literature. From a most general point of view, these approaches can be subdivided into two broad classes which describe the local and non-local closures, respectively. The former is based on the assumption that the turbulent fluxes depend only on the mean quantities. The latter means that the turbulent fluxes are described as more or less arbitrary functionals of the mean flow parameters. There is no strong separation between these two classes, since some characteristics of non-locality can be also found in the local closures.

In order to account for the "countergradient" heat transport in the boundary layer, the Boussinesq hypothesis may be generalized as follows (e.g. Budyko & Yudin 1946; Priestley & Swinbank 1947; Deardorff 1966):

$$\overline{\theta' w'} = -K_\theta \left(\frac{\partial \theta}{\partial z} - \gamma_\theta \right), \tag{3.7}$$

where K_θ is, as before, the eddy diffusivity coefficient, and the term γ_θ is a "countergradient correction" term. The expressions for this term can be derived from the equations for higher order moments and from the large eddy simulation data.

For example, the Deardorff (1972) formula reads

$$\gamma_\theta = \lambda \frac{\overline{\theta'^2}}{\overline{w'^2}}. \tag{3.8}$$

Assuming, in particular, that $\overline{\theta'^2}$ and $\overline{w'^2}$ are constant throughout the convective boundary layer, Deardorff (1973) suggested

$$\gamma_\theta = \lambda \frac{\theta_*^2}{w_*^2} = \frac{\overline{w'\theta'}_0}{w_* h}, \tag{3.9}$$

where $\overline{w'\theta'}_0$ is the surface kinematic heat flux, $w_* = (\lambda \overline{w'\theta'}_0 h)^{1/3}$ is the convective velocity scale, h is the height of the convective boundary layer, and $\theta_* = \overline{w'\theta'}_0 / w_*$ stands for the convective temperature scale. Another modifications of the formula (missing?) can be found in the overview by Lykossov (1995).

Since turbulence in the convective boundary layer of the atmosphere is usually characterized by the narrow, intense, rising plumes and by the broad, low-intensity subsiding motions between the plumes, the vertical velocity field is strongly skewed. Wyngaard (1987) suggested that this skewness is responsible for the difference between the top-down and bottom-up scalar (e.g. heat) diffusion. This dependence of the scalar diffusivity on the location of the source was termed by Wyngaard & Weil (1991) as "transport assymetry". They found that the interaction between the skewness and the gradient of the transported scalar flux can induce this assymetry. Using the kinematic approach, Wyngaard & Weil (1991) derived an expression for the scalar flux which in the case of heat flux has the following form:

$$\overline{w'\theta'} = -K_\theta \left(\frac{\partial \theta}{\partial z} - \frac{1}{2} A S \sigma_w T_L \frac{\partial^2 \theta}{\partial z^2} \right), \tag{3.10}$$

where $S = \overline{w'^3 / (\overline{w'^2})^{3/2}}$ is the skewness of w, $\sigma_w = (\overline{w'^2})^{1/2}$, T_L is the Lagrangian integral time scale, A is a constant. Thus, in order to describe the scalar flux, the term proportional to the second derivative of the mean quantity is added in the theory, suggested by Wyngaard & Weil (1991), to the eddy diffusivity term.

A non-local generalization of the Boussinesq hypothesis can be written in the following form:

$$\overline{a'w'} = -\int_0^\infty K(z,z')\frac{\partial a}{\partial z'}dz'. \tag{3.11}$$

This idea was first suggested by Berkowicz & Prahm (1979) with the application to the air pollution studies. Their generalization of the diffusivity theory is based on the so-called spectral turbulent diffusivity concept, according to which the eddy diffusivity coefficient of single Fourier component of the passive scalar field is treated separately as a function of the wave number k. For each individual mode the Boussinesq closure (3.3) is used. In the case of spectral diffusivity $K(k) = K_0$ for all k, it follows that $K(z,z') = K_0\delta(z-z')$, and Eq. (3.11) coincides with formula (3.3). The integral closure of the type (3.11) was also used by Fiedler (1984) and Hamba (1995). A similar closure can be derived from a set of equations for the second and third moments applied for the modelling non-local turbulent transport of momentum in jet-like flows (Lykossov 1993).

3.2.4 The atmospheric Ekman layer

It is experimentally known that outside of tropics the mean wind changes direction with height in the transition layer (above the surface layer) and nearly coincides with the free-atmosphere velocity at heights far enough from the underlying surface. At the same time, the momentum fluxes $\overline{u'w'}$ and $\overline{v'w'}$ (as well as the stress components $\rho\overline{u'w'}$ and $\rho\overline{v'w'}$) decrease with height. When horizontal homogeneity is assumed and the viscous stress is neglected, the momentum equations can be written as follows (e.g. Brown 1974):

$$\frac{\partial u}{\partial t} + \frac{\partial \overline{u'w'}}{\partial z} = -\frac{1}{\rho}\frac{\partial p}{\partial x} + fv, \tag{3.12}$$

$$\frac{\partial v}{\partial t} + \frac{\partial \overline{v'w'}}{\partial z} = -\frac{1}{\rho}\frac{\partial p}{\partial y} - fu. \tag{3.13}$$

Assuming the steady-state conditions, the geostrophic balance between Coriolis and pressure-gradient forces in the free atmosphere can be written

$$\frac{1}{\rho}\frac{\partial p}{\partial x} = fv_g, \qquad \frac{1}{\rho}\frac{\partial p}{\partial y} = -fu_g, \tag{3.14}$$

where subscript g indicates the geostrophic wind, which is often taken as the upper boundary condition for the boundary flow.

Let us now consider the steady-state version of Eqs. (3.12) and (3.13)

$$\frac{d\overline{u'w'}}{dz} = f(v - v_g), \tag{3.15}$$

$$\frac{d\overline{v'w'}}{dz} = -f(u - u_g), \tag{3.16}$$

where u_g and v_g are substituted for the constant pressure gradients from Eq. (3.14). It is seen from Eqs. (3.15) and (3.16) that momentum is generated in the boundary layer by ageostrophic components of the wind velocity and transported down to the surface.

The most idealized model of the wind structure can be derived for this case with the help of the K-theory turbulence closure based on Eq. (3.3). The resulting equations with the constant eddy diffusivity coefficients $K_u = K_v \equiv K$ are known as the Ekman layer equations. When the x-axis is aligned with the geostrophic wind, these equations are written as follows:

$$K\frac{d^2u}{dz^2} + fv = 0,$$
$$K\frac{d^2v}{dz^2} - f(u - G) = 0, \tag{3.17}$$

where $G = \sqrt{u_g^2 + v_g^2}$ is the geostrophic wind speed. An analytical solution of these equations for the ocean was derived by Ekman (1905), and for the atmosphere, by Akerblom (1908). In complex notation, these equations take the form

$$\frac{d^2W}{dz^2} - \mathbf{i}\frac{f}{K}(W - G) = 0, \tag{3.18}$$

where $W = u + \mathbf{i}v$, and $\mathbf{i} = \sqrt{-1}$. The solution to Eq. (3.18), subject to the boundary conditions

$$W = 0 \quad \text{at} \quad z = 0, \tag{3.19}$$

$$W \to G \quad \text{as} \quad z \to \infty, \tag{3.20}$$

is written as follows:

$$W - G = -G\exp\left(-\frac{z}{h_E}\right)\left[\cos\left(\frac{z}{h_E}\right) - \mathbf{i}\sin\left(\frac{z}{h_E}\right)\text{sign}f\right], \tag{3.21}$$

where $h_E = \sqrt{2K/|f|}$. For $K = 12.5$ m^{2}s^{-1} and $f = 10^{-4}$ s^{-1}, the value of $h_E = 500$ m. Note that the quantity πh_E is the lowest height (the Ekman layer depth) where the boundary layer wind is parallel to the geostrophic wind. The $u-$ and $v-$component of the solution (3.21) are expressed in the form

$$u = G\left[1 - \exp\left(-\frac{z}{h_E}\right)\cos\left(\frac{z}{h_E}\right)\right],$$
$$v = G\exp\left(-\frac{z}{h_E}\right)\sin\left(\frac{z}{h_E}\right)\text{sign}f. \tag{3.22}$$

It seen from this solution that the wind veers with height, giving the so-called Ekman spiral wind profile, and slightly overshoots the geostrophic value. Since

$$\tan\alpha = \lim_{z\to 0}\frac{v}{u} = \lim_{z\to 0}\frac{dv/dz}{du/dz} = \text{sign}f,$$

the surface wind is parallel to the stress and directed to the left (right) of the free-atmosphere wind in the northern (southern) hemisphere. In this idealized model, the angle α between the surface and geostrophic wind is equal to 45° and does not depend on geographical location and meteorological situation. However, this is not true for the real atmospheric boundary layer. Observed angles between the geostrophic and surface winds may vary considerably due to fact that ageostrophic components of the wind and, consequently, the momentum transfer may be influenced by various physical processes. For example, the wind profile is sometimes characterized by the presence of a low level jet in the upper part of the Ekman layer (e.g. Pennel & LeMone 1974; Smedman *et al.* 1995). A baroclinicity of the free-atmosphere flow, which can be expressed, using the thermal wind relationships, in terms of the height-dependent geostrophic wind, may significantly alter the Ekman profile. The constant K assumption is also not valid since K is linearly increasing with height, at least, near the surface. A lot of proposed theoretical distributions of the eddy diffusivity coefficient $K(z)$ is presented in the literature (see, for example, Brown 1974; Holt & SethuRaman 1988 for the review). Nevertheless, it is widely recognized that the modelling of the more or less real dynamics of the boundary layer requires more sophisticated approaches, a brief review of which is given in Section 3.2.2.

3.2.5 Surface layer

In the atmospheric surface layer, typically the lower 10 % of the boundary layer, the turbulent fluxes of momentum, water vapor and sensible heat are nearly constant with height. In a surface microlayer less than 1 cm thick, the molecular transport dominates over turbulent transport. Over oceans, diabatic processes and wave motion are dominating in their influence on the wind shear. Customarily, these two effects are treated as independent (e.g. Hasse & Smith 1996).

Orientating the $x-$axis in the surface stress direction, and following to Brown (1974), one can transform Eq. (3.15) to the nondimensional form by the use of an arbitrary characteristic velocity scale V_0 and vertical scale H together with the surface stress τ_0. This produces

$$E\frac{d\tilde{\tau}}{d\tilde{z}} + \tilde{v} - \tilde{v_g} = 0, \tag{3.23}$$

where the tilde indicates nondimensional variables, and $E = \tau_0/(\rho f V_0 H)$ is the Ekman number. For $H \to 0$, the parameter $E \to \infty$, and Eq. (3.23) yields

$$d\tilde{\tau}/d\tilde{z} = 0,$$

subject to the boundary condition $\tilde{\tau}(0) = 1$. Integration gives $\tau \equiv \tau_0$. Observations show that close to the surface, where E is large, the layer of nearly constant stress can be really detected. Similarly, it can be shown that for steady state conditions the buoyancy flux $-g/\rho\overline{\rho' w'}$ in the atmospheric surface layer is nearly constant.

Assuming that for neutral conditions in this layer the eddy diffusivity coefficient $K = \kappa u_* z$ (Prandtl 1932), where $u_* = \sqrt{\tau_0/\rho}$ is the so-called friction velocity, one can obtain

$$\frac{\partial u}{\partial z} = \frac{u_*}{\kappa z}, \tag{3.24}$$

and on integration the wind profile

$$u(z) = \frac{u_*}{\kappa} \ln\left(\frac{z}{z_0}\right). \tag{3.25}$$

Here z_0 is an integration constant called the roughness length. In some sense, this parameter is an artificial quantity which results from extrapolation of the wind profile to zero wind speed. However, it was experimentally found that z_0 is a parameter, which "in the whole" characterizes the geometrical properties of the solid underlying surface.

Contrary to the land surface, where the roughness is determined by the roughness elements of fixed geometry, the sea surface should be considered as the interface of two fluids of different density, that both are in motion and may generate waves. In this case z_0 will not reflect "topology" of the sea surface but must be obtained, if possible, as a parameter, characterizing dynamics of the interfacial layer. The experimentally derived Charnock (1955) formula for the calculation of z_0 from the friction velocity u_* (see Chapter 2) can be considered as an example of such parameterization. It is necessary to point out that the logarithmic wind profile (3.25) is valid only well above the height z_0 (see, e.g. Monin & Yaglom 1971). Close to the surface, the viscous stress, which is neglected in Eq. (3.24) should determine the wind profile.

Using an aerodynamic approach, from the surface stress τ_0, mean wind speed u, and density of air ρ can be derived the drag coefficient C_d as

$$\tau_0/\rho = C_d u^2. \tag{3.26}$$

Note that since u is a function of z, this coefficient depends on height. Typically, it is defined for a reference height of 10 to 25 m above the sea level. This approach is commonly used as a parameterization to describe the air-sea fluxes, the only practicable tool that is available to apply results of empirical investigations. In the neutrally stratified momentum constant flux layer, from the logarithmic profile, z_0 is related to the neutral drag coefficient C_{dN} by

$$C_{dN} = \left(\frac{u_*}{u}\right)^2 = \left[\kappa/\ln\left(\frac{z}{z_0}\right)\right]^2. \tag{3.27}$$

Eq. (3.27) shows that there is a one-to-one relation between the drag coefficient and the roughness length so that C_d grows when z_0 increases. At very low winds, no waves are generated and the stress should not be less than for an aerodynamically smooth flow, for which $z_0 = 0.111\nu/u_*$ (Schlichting 1951) where $\nu = 0.14 \times 10^{-4}$ m^{2}s^{-1} is the kinematic viscosity of air. This smooth flow constraint causes the drag coefficient to increase with decreasing wind speed below about 3 ms^{-1} (e.g. Zilitinkevich 1970; Wippermann 1972; Smith 1988).

For the non-neutral conditions, the Monin-Obukhov (1954) similarity theory predicts that the dimensionless gradient of the wind velocity can be expressed by an universal function of dimensionless stability z/L only

$$\frac{\kappa z}{u_*}\frac{\partial u}{\partial z} = \phi_M(z/L), \tag{3.28}$$

where $\phi_M(0) = 1$, and the stability parameter L is the Monin - Obukhov length scale

$$L = \frac{\rho u_*^3}{\kappa g \overline{\rho' w'}}. \tag{3.29}$$

To calculate the buoyancy flux, the following relation is usually used:

$$\frac{g}{\rho}\overline{\rho' w'} = -\frac{g}{\theta_v}\overline{\theta_v' w'}. \tag{3.30}$$

Note that a positive buoyancy flux $\overline{\rho' w'}$ is away from the surface. The flux would be expected to be positive when the virtual potential temperature θ_v increases with height. This is known as stable conditions.

Given the surface roughness, integration of Eq. (3.28) from z_0 to z results in a diabatic wind profile (see e.g. Monin and Yaglom 1971)

$$u(z) = \frac{u_*}{\kappa}\left[\ln\left(\frac{z}{z_0}\right) - \psi_M\left(\frac{z}{L}\right)\right] \tag{3.31}$$

where

$$\psi_M\left(\frac{z}{L}\right) = \int_0^{z/L} \frac{1-\phi_M(\zeta)}{\zeta} d\zeta. \tag{3.32}$$

is the integrated universal function for velocity. The drag coefficient C_d can be now formally expressed as

$$C_d = \left(\frac{u_*}{u}\right)^2 = \kappa^2 / \left[\ln\left(\frac{z}{z_0}\right) - \psi_M\left(\frac{z}{L}\right)\right]^2. \tag{3.33}$$

Eq. (3.33) shows that in the non-neutral constant flux layer the drag coefficient depends not only on the roughness length z_0 but on the stability parameter L too, which also reflects dynamics of the air-sea interfacial layer. This is especially important for those ocean regions where highly unstable thermal stratification can be formed in the atmospheric surface layer. On the other hand, the z/L- dependence of universal functions means that the boundary layer is nearly neutral close to the surface and becomes more non-neutral when the height increases. Note also that when the winds are strong, the u_* value is generally high, and $\psi_M(z/L) << \ln(z/z_0$ so that the surface layer can be again considered as neutrally stratified.

Since the stability parameter includes the virtual potential temperature flux, which can be treated as a linear combination of the potential temperature flux and water vapour flux

$$\overline{\theta_v' w'} \approx (1 + 0.61q)\overline{\theta' w'} + 0.61\theta\overline{q' w'}, \tag{3.34}$$

it is advisable to give here the corresponding dimensionless fluxes in the form of the Monin-Obukhov universal functions

$$\frac{\kappa z}{\theta_*}\frac{\partial \theta}{\partial z} = \phi_H\left(z/L\right), \qquad \frac{\kappa z}{q_*}\frac{\partial q}{\partial z} = \phi_E\left(z/L\right), \tag{3.35}$$

and their aerodynamic representations

$$\overline{\theta' w'} \equiv u_*\theta_* = C_H u(\theta_s - \theta), \qquad \overline{q' w'} \equiv u_* q_* = C_E u(q_s - q), \tag{3.36}$$

where $u_*\theta_* = \overline{\theta' w'}_s, u_* q_* = \overline{q' w'}_s$, and subscript s refers to the values at the surface. As above for the drag coefficient C_d, the bulk transfer coefficients C_H (Stanton number) and C_E (Dalton number) are expressed with the help of integrated universal functions

$$\begin{aligned} C_H &= \alpha_H \kappa^2 \left[\ln\left(\frac{z}{z_0}\right) - \psi_M(z/L)\right]^{-1} \left[\ln\left(\frac{z}{z_H}\right) - \psi_H(z/L)\right]^{-1}, \\ C_E &= \alpha_E \kappa^2 \left[\ln\left(\frac{z}{z_0}\right) - \psi_M(z/L)\right]^{-1} \left[\ln\left(\frac{z}{z_E}\right) - \psi_E(z/L)\right]^{-1}, \end{aligned} \tag{3.37}$$

where α_H and α_E stand for the ratios of the eddy diffusivities of sensible heat and water vapor to that of momentum, z_H and z_E are the roughness lengths for temperature and specific humidity, respectively.

Similar to the aerodynamic roughness z_0, these quantities are associated with a logarithmic profile and the surface fluxes. The magnitude of z_H and z_E is controlled by transport mechanisms very close to the surface where molecular processes dominate. This is especially important under low-wind, unstable conditions over water (Godfrey & Beljaars 1991). As an example of parameterization for z_H, the following approximation of an experimental data obtained for natural and artificial surfaces (Garratt & Hicks 1973; Hicks 1975; Garratt 1977) is suggested (Kazakov & Lykossov 1982):

$$\ln(z_0/z_H) = \begin{cases} -2.43 & \text{for } Re_* \le 0.111, \\ 0.83\ln(Re_*) - 0.6 & \text{for } 0.111 \le Re_* \le 16.3, \\ 0.49 Re_*^{0.45} & \text{for } Re_* \ge 16.3, \end{cases} \tag{3.38}$$

where $Re_* = u_* z_0/\nu$ is the roughness Reynolds number.

The use of the diabatic wind profile requires the knowledge of stability functions $\phi_M(z/L), \phi_H(z/L)$ and $\phi_E(z/L)$. Variations of stability are more pronounced over land than over sea. Hence it has been found suitable to adopt stability functions determined over land for the use over sea, too. It is usually assumed that $\phi_H = \phi_E$. In the case of stable stratification, the linear type functions are theoretically derived and experimentally supported (see e.g. Monin & Yaglom 1971)

$$\begin{aligned} \phi_M = \phi_H &= 1 + \beta\zeta, \\ \psi_M = \psi_H &= -\beta\zeta, \end{aligned} \tag{3.39}$$

where $\zeta = z/L$, and the parameter β varies, according to observations, from 4.7 to 5.2 (Panofsky & Dutton 1984). For the regions with moderate unstable stratification ($-2\zeta \le 0$), the Businger-Dyer formulations are widely used (Businger *et al.* 1971; Dyer 1974)

$$\phi_M = (1 - \alpha\zeta)^{-1/4}, \quad \phi_H = (1 - \alpha\zeta)^{-1/2}, \tag{3.40}$$

where values of α ranging from 16 to 28 fitted the data derived from the measurements over oceans (Edson *et al.* 1991). The corresponding integrated universal functions have the following form (Paulson 1970)

$$\psi_M(\zeta) = \ln\left[\frac{1}{8}\left(1+\phi_M^{-2}\right)\left(1+\phi_M^{-1}\right)^2\right] - 2\arctan\phi_M^{-1} + \pi/2,$$

$$\psi_H(\zeta) = 2\ln\left[\frac{1}{2}\left(1+\phi_H^{-1}\right)\right], \tag{3.41}$$

When convection dominates so that ζ are large and negative (in particular, in the case of light winds), the universal functions should vary as $(-\zeta)^{-1/3}$, a relation called the free-convection condition (Panofsky and Dutton 1984). Carl *et al.* (1973) suggested for momentum that

$$\phi_M = (1-16\zeta)^{-1/3}, \tag{3.42}$$

which satisfy to this condition when $-\zeta$ becomes large. The corresponding integrated universal function can be written as follows:

$$\psi_M = \frac{3}{2}\ln\left[\frac{1}{3}(X^2+X+1)\right] - \sqrt{3}\left(\arctan\frac{2X+1}{\sqrt{3}} - \frac{\pi}{3}\right), \tag{3.43}$$

where $X = (1-16\zeta)^{1/3}$. A similar -1/3 power law dependence is also required for ϕ_H in order to satisfy the theoretical prediction. To combine the Businger-Dyer expressions and free-convection limit, one can use (Kazakov & Lykossov 1982; Large *et al.* 1994)

$$\phi_a = (b_a - c_a\zeta)^{-1/3} \quad \text{for} \quad \zeta < \zeta_a, \tag{3.44}$$

where a stands for M or H, and the b_a and c_a are chosen so that both ϕ_a and its first derivative are continuous across the matching value $\zeta = \zeta_a$.

In order to show an importance of the air-sea temperature difference, the drag coefficient C_d and transfer coefficients C_H and C_E are sometimes presented in the form that depends on the bulk Richardson number Ri_b

$$Ri_b = \frac{gz(\theta_v - \theta_{vs})}{\theta_v u^2}. \tag{3.45}$$

The variables Ri_b and z/L can be converted into each other, when the stability functions are known (see, for example, Launiainen 1995).

The above presented consideration does not include effects of the sea surface waves. In a wave boundary layer, part of the shear stress is replaced by momentum flux carried by pressure fluctuations to surface waves (Hasse and Smith 1996). On the other hand, if the air flow is modulated by the wavy surface, the mean wind profile may be distorted. For example, Dittmer (1977) derived two average diabatic wind profiles from GATE - for wave heights below 25 cm and between 25 and 75 cm - and found that there was a more pronounced deformation of the wind profile with increased wave height. Moreover, the measurements of wind profiles, carried out during the JONSWAP experiment (Hasselmann *et al.* 1973), showed that 1) the profile slope depends on wave energy, but not on mean wind speed, and 2) the wave-influence on the profile is confined mainly to the lower heights which are comparable to the wave height (Kruegermeyer *et al.* 1977; Hasse *et al.* 1978).

To explicitly separate the relative influences of mean, wave, and turbulence components of the wind field, one can decompose the instantaneous horizontal and vertical wind velocity as

$$a = \overline{a} + \tilde{a} + a', \qquad a = (u, w), \tag{3.46}$$

where $\overline{a}$ is the time-average component, $\tilde{a}$ stands for the periodic wave-induced component, and a' is the turbulence component of the motion (e.g. Anis & Moum 1995). Time averages must be performed over time scales much larger than the characteristic wave period. Assuming that the mean, the periodic wave-induced, and the turbulence components of the motion are non-correlated, one can formulate the constant stress approximation as follows:

$$-\rho(\overline{u'w'} + \overline{\tilde{u}\tilde{w}}) = const = \tau_0. \tag{3.47}$$

The problems related to the wave-induced momentum flux $\overline{\tilde{u}\tilde{w}}$ are considered in Chapter 4.

3.2.6 Clouds in the boundary layer

Marine low-level clouds cover a large area of the World ocean surface. They are very important for the radiation budget and play a significant role in the surface energy budget and in the water balance of the atmosphere. These clouds determine also the vertical structure of the turbulent fluxes of momentum, moisture and heat. There are two major forms of the boundary layer clouds: stratocumulus clouds and cumulus clouds (Jonas 1993). Solid stratocumulus cover extensive areas of the oceans over the cool water; cumulus clouds are mainly observed in trade-wind regions. A strong correlation between regions of strong radiative cooling and regions of extended stratocumulus cloud cover in the marine boundary layer was derived from the climotological data (Bretherton 1993). Stratocumulus and boundary-layer cumulus have many features in common and very often the transition from one form to another takes place. It is known from observations that there is also a climatological transition from nearly solid relatively shallow subtropical stratocumulus to trade cumulus clouds with lower fractional cloud cover and a deeper boundary layer (Bretherton 1993).

Marine stratocumulus clouds have intensively been studed within frame of recent field programs FIRE (First International Satellite Cloud Climatology Regional Experiment, 1987) and ASTEX (the Atlantic Stratocumulus Transition Experiment, 1992). Selected FIRE and ASTEX results relevant for the development of cloudy boundary layer models are presented by Albrecht *et al.* (1988), Albrecht (1993), Bretherton (1993) and Tjernström & Rogers (1996). It was found from experimental and model investigations that the list of important physical processes which control the structure and type of stratocumulus includes, in particular, cloud top entrainment instability, diurnal turbulent decoupling, microphysics and drizzle.

Cloud-topped boundary layers are usually capped by warm, dry air and there is the tendency for the cloud to dissipate due to the entrainment of this air into the boundary layer. Negatively buoyant downdrafts produce additional turbulent kinetic energy (TKE) which can enhance mixing and entrainment

(Stull 1988). The additionally entrained air can then become unstable and again produce more TKE and cause more entrainment. This positive feedback process can lead to the rapid breakup and evaporation of the cloud. On the other hand, one could expect that such a mechanism of the top-down convection, which causes enhanced turbulence, might lead also to the enhanced momentum transfer.

It is also known (Johnson 1993) that boundary layers capped by stratocumulus clouds demonstrate a large diurnal variations. Long wave cooling from the top of the cloud produces TKE that forms large eddies which transport water vapour up from the sea surface to the cloud layer. During the day, solar absorption by the cloud reduces the effect of the long wave cooling and consequently, the size of the vertical eddies. If the surface heat and water vapor fluxes are not sufficiently strong, as it takes place in near-neutral surface conditions, the heating of the lower parts of the cloud may produce a secondary inversion between the cloud base and surface (Tjernström & Rogers 1996). Thus, the surface layer becomes decoupled from the cloud and sub-cloud layer (in the sense that the turbulent fluxes are severed), and this cuts off the moisture supply to the cloud. The entrainment of dry air will tend to thin the cloud and even break up. On the other hand, moisture build up in the well mixed surface layer can lead to condititional instability which produces small cumulus clouds at the top of the surface layer that can grow and penetrate the stratocumulus layer. If the boundary layer becomes deep enough, then it may remain decoupled all the time. In this case, the cumulus clouds will be produced even at night and will then restore a recoupling between the surface layer and the cloud. Such a mechanism may transport enough moisture from the surface layer to the top of the boundary layer to maintain or thicken the stratocumulus. The field experiments, carried out over quite different regions of the ocean, such as FIRE (Moyer & Young 1993), ASTEX (Rogers *et al.* 1995) and the North Sea stratocumulus study project (Nicholls 1989) demonstrated that decoupling of the surface layer from the cloud and sub-cloud layer is not a very frequent feature of the marine atmospheric boundary layer.

The aerosol characteristics of the boundary layer play an important role in microphysics and radiative transfer of the cloud layer. The cloud condensation nuclei govern the size and number of cloud drops. Drizzle production can be an effective decoupling mechanism due to evaporative cooling and can also limit the cloud liquid water. The presence of drizzle is highly correlated with cumulus - stratocumulus interactions which has a modifying effect on the reflectivity of the stratocumulus (Johnson 1993).

3.3 Oceanic Upper Boundary Layer

3.3.1 Vertical structure

The structure of the oceanic upper boundary layer can consist at least of three layers (Anis & Moum 1992). The oceanic analogue to the atmospheric viscous sublayer is the cool skin which has an average thickness of a few millimeters and very large temperature gradients (Khundzhua *et al.* 1977). Dynamics of this layer is mainly driven by the surface wave breaking, Langmuir circulations and shear due to the wind stress and wave drift. Turbulence here is an intermediary in the transfer of momentum, heat and salt between ocean and atmosphere. Below the oceanic surface layer two essentially different layers are observed. First, the upper quasi-homogeneous well mixed layer includes large scale convective eddies with the size of an order of the whole layer. The presence of these coherent structures causes the existence of certain differences in the modelling upper ocean dynamics. In particular, the effect of "negative viscosity" may be encountered (Muraviev & Ozmidov 1994). Second, the underlying thermocline is characterized by an abrupt increase of density with depth and, consequently, by a very stable stratification. Turbulence exists in the thermocline and not fully developed but intermittent (Kraus 1977; Monin & Ozmidov 1981; Gargett 1989). The internal waves are very often observed here.

3.3.2 Turbulence

One can note the following major mechanisms of the oceanic turbulence production: shear instability, internal wave breaking, double diffusion, and deep convection (Monin & Ozmidov 1981; Large *et al.* 1994). To calculate the vertical turbulent fluxes of momentum, heat and salt, corresponding to the first three listed processes, the K-theory closure (3.3) is widely applied. It is assumed (Large *et al.* 1994) that the eddy diffusivity coefficient can be parameterized as a sum of coefficients characterizing a separate process

$$K_a = K_a^s + K_a^w + K_a^d, \tag{3.48}$$

where a stands for momentum, temperature or salinity. The deep convection mechanism requires a more complicated approach.

1. *An instability caused by the vertical gradients of the drift currents velocity.* This shear-generated turbulence develops in the whole upper layer of the ocean as the result of direct action of the wind on the sea surface. Since Reynolds numbers for the drift currents are very large (of order 10^7) and significantly exceed the critical Reynolds number (of order 10^3) this kind of turbulence is produced practically everywhere in the World ocean. The turbulent mixing of the oceanic upper layer takes place in the density-stratified sea water when the vertical velocity shear overcomes the stabilizing effect of the buoyancy gradient. This process is characterized by the gradient Richardson number

$$Ri = -\frac{g}{\rho}\frac{\partial\rho/\partial z}{|\partial\vec{V}/\partial z|^2},$$

where the vertical coordinate z is positive up, ρ is the water density, and $\vec{V}$ is the horizontal current velocity vector with the components u and v. Shear instability occurs when Ri is below some critical value Ri_0. According to oceanic field measurements, Ri_0 is generally higher than the theoretical value of 0.25 and vary from 0.4 to 1 (Large *et al.* 1994). The eddy diffusivity coefficients K_a^s are often chosen as depending on the gradient Richardson number Ri and are the same for momentum, heat and salt. For example, the parameterization suggested by Large *et al.* (1994) reads

$$K_a^s = \begin{cases} K_{max} & \text{for } Ri \le 0, \\ K_{max}[1-(Ri/Ri_0)^2]^3 & \text{for } 0 < Ri < Ri_0, \\ 0 & \text{for } Ri_0 < Ri, \end{cases} \tag{3.49}$$

where $K_{max} = 50 \times 10^{-4}$ m^{2}s^{-1} and $Ri_0 = 0.7$.

2. *A breaking of surface waves and an hydrodynamical instability of the wave motions in the oceanic upper layer.* The superposition of internal waves increases shear and consequently, the Richardson number decreases. The eddy diffusivity coefficient K_a^w, describing the effect of mixing due to internal wave breaking, seems to be small and depends mainly on the internal wave energy. It is found that the internal wave diffusivity for heat and salt is about 0.1×10^{-4} m^{2}s^{-1} (Ledwell *et al.* 1993), and for $Ri_0 < Ri$ the wave momentum transfer is expected to be from 7 to 10 times more effective (Peters *et al.* 1988).

3.*Double-diffusive convection.* This is a very important process by which the heating (or cooling) and the salting (or freshening) at the sea surface become distributed in the oceanic boundary layer. Such ocean-mixing mechanism is associated with the fact that the density of sea water ρ is determined by the temperature T and the salinity S with the large (of two orders) differences in their molecular diffusivities. During this process a statically unstable vertical distribution of one property can be balanced by a distribution of the other. While the resulting density distribution is stable, small-scale instabilities can lead to release of gravitational potential energy from the unstable component. The mixing of momentum from the double-diffusive convection is found (e.g. Large *et al.* 1994) to be the same as for salt ($K_M^d = K_S^d$), but the temperature diffusivity K_T^d is quite different. Depending upon whether the resulting motions are driven by energy stored in the component of the higher (T) or lower (S) diffusivity, two basic types of convective instabilities occur (the "diffusive" and "finger" forms) which differ in the relative efficiency of heat and salt transport (Monin & Ozmidov 1981; Turner 1985; Gargett 1989). To quantify these two regimes, the following ratio of the density flux due to heat to that due to salt is used:

$$R_f = \frac{\beta F_T}{\alpha F_S}, \tag{3.50}$$

where F_T and F_S are the heat flux and the salt flux, respectively; $\alpha = -\rho^{-1}(\partial\rho/\partial T), \beta = \rho^{-1}(\partial\rho/\partial S)$. It was found that $R_f < 1$ for the diffusive case and $R_f > 1$ for the fingering case (Gargett 1989). The instability growth rate anf the flux ratio R_f are functions of the stability parameter

$$R_\rho = \frac{\beta \partial S/\partial z}{\alpha \partial T/\partial z}. \tag{3.51}$$

The diffusive instability occurs in regions where cold, dilute water lies above warm salty water. As the result of this kind of instability, weakly stirred convective layers separated by much thinner interfaces of strong molecular transport are formed. A part of the potential energy released from the heat field during the thermal convection is converted to the kinetic energy to transport salt upward. It was found that the flux ratio R_f decreases with increase of the stability parameter and becomes nearly constant ($R_f \approx 0.15$) for $R_\rho > 2$. Layered structures in T and S are relatively rare in the World ocean. They are mainly observed in polar regions where circumstances are favorable for the diffusive instability (Gargett 1989). For example, the following parameterization is suggested (Fedorov 1988; Large *et al.* 1994):

$$K_T^d = 0.909\nu_w \exp(4.6\exp[-0.54(R_\rho - 1)]),$$
$$K_S^d = \begin{cases} K_T^d(1.85R_\rho^{-1} - 0.85) & \text{for } 1 < R_\rho \le 2, \\ 0.15K_T^d R_\rho^{-1} & \text{for } R_{rho} > 2, \end{cases} \tag{3.52}$$

where ν_w is molecular viscosity of the water.

The salt-fingering instability occurs in regions where warm salty water lies above cold, dilute water. Extended thin columns of fluid moving vertically in both directions form the "salt finger" pattern of the resulting motion. Each upward-moving finger is surrounded by downward-moving fingers, and vice versa. The downgoing fingers lose heat and become more dense, whereas the upgoing fingers gain heat and become less dense. The potential energy is now derived from the salt field. It was experimentally discovered that over a wide range of conditions the flux ratio R_f is nearly constant and close to 0.56 (Turner 1985). Finger activity is strongest at $0.5 < R_\rho < 1$ (Gargett 1989). Large regions of the subtropical and tropical oceans are favorable for the salt fingering process. Mixing due to this process can be parameterized, for instance, in the form (Large *et al.* 1994)

$$K_S^d = \begin{cases} K_f\left[1 - \left(\frac{R_{\rho 0}}{R_\rho}\right)^2 \left(\frac{1-R_\rho}{1-R_{\rho 0}}\right)^2\right]^3 & \text{for } R_{\rho 0} < R_\rho < 1, \\ 0 & \text{for } R_\rho \le R_{\rho 0}, \end{cases}$$

$$K_T^d = 0.7K_S^d, \tag{3.53}$$

where $K_f = 10 \times 10^{-4}$ m^{2}s^{-1} and $R_{\rho 0} = 0.526$.

4. *Convection due to unstable density stratification.* Convection in the upper part of the ocean can be caused by cooling of the sea surface in cold seasons or by accumulation of salt due to intensive evaporation of the sea water. This process is mostly pronounced in the form of open-ocean (deep) convection. During deep convection localized bursts of violent vertical mixing transport and mix water over several hundred meters and short time periods (Aagard and Carmack 1989). Observations indicated that vertical velocities can exceed in this case 0.1 $ms{-}1$. There are identified three phases associated with deep convection: preconditioning, violent mixing, and sinking and spreading (Alves 1995).

During the preconditioning phase the vertical static stability is reduced so that the formation of convectively unstable layers can begin. The mechanisms which may trigger a deep convection are instabilities of a different kind, e.g. double diffusion, convective instability (heating, cooling, evaporation or brine release), thermobaric instability (caused by depth dependence of the equation of state), baroclinic instability, and Kelvin-Helmholtz instability (Chu 1991). The main areas of deep convection are reviewed by Alves (1995) and include the Labrador, Greenland and Baltic Sea, the eastern Mediterranean, the Adriatic Sea and the Gulf of Lions in the north-western Mediterranean, the Weddel Sea and the Bransfield Strait at the tip of Antarctica, the subpolar oceans near the ice edge and in polynyas.

The violent mixing phase is the period when dense plumes develop and gradually extend deeper into the ocean transporting and mixing the sea water in the vertical. These plumes have vertical scales of 1-2 km and horizontal scales of 100 - 1000 m. It was found from observations that the plumes form a nearly homogeneous patch of the sea water which is called a chimney. A density-driven rim current of the local Rossby radius of a few kilometers exists between the chimney and the surrounding fluid. Since the vertical static stability is low, the rim current is baroclinically unstable. This leads to the development of baroclinic eddies which advect relatively lighter water from the surroundings into the chimney and limit the convective activity (Madec & Crepon 1991). Thus, the eddies and plumes are responsible for the two vertical mixing processes which determine this stage of deep convection.

The sinking and spreading phase begins when surface forcing turns off. The dense water sinks under gravity, and water from outside the chimney restores the stratification of the near-surface layer (Jones & Marshall 1993). During a few days the chimney breaks up into geostrophically adjusted fragments ("cones") with a spatial size of several kilometers. The possible mechanisms for the breakdown of the chimney are also mixing by internal waves, topography, and shear in the large scale cyclonic circulation (Killworth 1983).

Oceanic deep convection remindes very much convection in the atmospheric boundary layer. In both media, the mean stratification is slightly stable, and a countergradient heat flux may exist. This means that the the collective effect of an ensemble of oceanic plumes might be also parameterized on the basis of a nonlocal closure theory.

3.3.3 The oceanic Ekman layer

By definition, the ocean drift current is driven only by the wind stress at the sea surface when the pressure gradient in the ocean is neglected. For steady state and horizontally homogeneous conditions, the equations of motion are of the Ekman type (3.17) but reduced to

$$\begin{aligned} K\frac{d^2u}{dz^2} + fv = 0, \\ K\frac{d^2v}{dz^2} - fu = 0, \end{aligned} \tag{3.54}$$

It is convenient to orientate the $x-$axis in the surface stress direction. Assuming that the current vanishes deep in the ocean, and the stress is continuous across the air-sea interface, the boundary conditions subject to Eq. (3.54) read

$$K\frac{du}{dz} = \sigma^{-1}u_*^2, \quad K\frac{dv}{dz} = 0 \quad \text{at} \quad z = 0, \tag{3.55}$$

$$u \to 0, \quad v \to 0 \quad \text{as} \quad z \to -\infty, \tag{3.56}$$

where u_* is the air friction velocity, and σ is the ratio of the water density to the air density. The solution to the problem (3.54), (3.56), obtained first by Ekman (1905), reads

$$u = \frac{u_*^2}{\sigma\sqrt{Kf}} \exp\left(\frac{z}{h_E}\right) \cos\left(\frac{z}{h_E} - \frac{\pi}{4}\right),$$
$$v = \frac{u_*^2}{\sigma\sqrt{Kf}} \exp\left(\frac{z}{h_E}\right) \sin\left(\frac{z}{h_E} - \frac{\pi}{4}\right) \mathrm{sign} f, \tag{3.57}$$

where, as for the atmosphere, $h_E = \sqrt{2K/|f|}$ but with K applied to ocean values. For $K = 12.5 \times 10^{-4}$ m^{2}s^{-1} and $f = 10^{-4}$s^{-1}, the value of $h_E = 5$ m.

As it is seen from (3.57), the surface current in the northern (southern) hemisphere is 45° to the right (left) of the surface stress, and consequently, it is parallel to the atmospheric geostrophic wind (compare (3.57) with (3.21) from Section 3.2.4). Taking Eq. (3.26) into account, $u_*^2 = C_d U_a^2$, where C_d and U_a are the drag coefficient and wind speed, say, at the height 10m, one can calculate from (3.57) the surface current speed $U_o = \sqrt{u^2 + v^2}_{z=0}$ as

$$U_o = \frac{C_d U_a^2}{\sigma\sqrt{Kf}}. \tag{3.58}$$

To estimate U_o, let us use $C_d = 10^{-3}$, $\sigma = 10^3$, $U_a = 10$ ms^{-1}, and $Kf = 9 \times 10^{-8}$ m^{2}s^{-2} as a typical values. Their substituting into the expression (3.58) leads to $U_o = 1/3$ ms^{-1}. Since $U_a \sim G$, where G is the geostrophic wind speed, this means that roughly the surface drift current 30 times weeker than the geostrophic wind.

Since

$$\tan\alpha = \frac{v}{u} = \tan\left(\frac{z}{h_E} - \frac{\pi}{4}\right) \mathrm{sign} f,$$

and z is negative, the drift current turns with depth to the right (left) of the surface current direction in the northern (southern) hemisphere. Thus, the large-scale horizontal divergence should be expected in the ocean under atmospheric regions of the large-scale horizontal convergence, and vice versa. This means that synoptic low (high) pressure systems in the atmosphere induce upwelling (downwelling) motion in the underlying oceanic domains (e.g. Stull 1988).

3.4 Processes at the Air-Sea Interface

Qualitatively new effects in the air-sea interaction are detected in the case of strong winds (nominally, above 15 $ms{-}1$). The development of the surface waves leads at high wind velocities to the significant change of the sea surface roughness length. The air-sea interface is disrupted under stormy conditions, during which intensive spray is detected in the atmospheric surface layer, air bubbles are found in the water, and as the result, that can be treated as a two-phase flow. The contribution of spray droplets to the sea surface heat and moisture budgets, as well as to the sea surface aerosol flux, is found to be important (Woodcock 1955; Toba 1965a, 1965b 1966; Borisenkov 1974; Wu 1979, 1990; Bortkovskii 1987; Ling 1993; Andreas *et al.* 1995). Shearing of wave crests by wind, aerodynamic suction at the crests of capillary waves, and bursting of air bubbles at the sea surface are principal mechanisms of the sea spray production (Wu 1979). It was shown that bubble bursting is the primary mechanism of producing spray due to a large volume of air entrained by breaking waves. The bubbles in the subsurface plume form a whitecap and produce two distinct types of droplets (film droplets and jet droplets) when they burst (Andreas *et al.* 1995). The film droplets are generated when the upper, protruding surface of a bubble thins by down-slope drainage and shatters. The jet droplets are produced by a microscopic column of water which forms in the center of the cavity resulted from the rupture of the bubble film. Jet droplets dominate the spectrum of spray droplet flux in the 1 - 100 μm radius range with initial vertical velocities of 5-20 ms^{-1}, whereas film droplets contribute mainly at radii below 0.1 - 10 μm (Smith *et al.* 1996). At high wind speeds spume drops, mainly larger than 40 μm, are blown off the crests of the spilling waves. It was found that the maximum ejection height for jet droplets is 18 cm (e.g. Wu 1979), but turbulence in the atmospheric boundary layer can carry the small bubble-derived droplets higher (de Leeuw, 1986) up to cloud heights, where they contribute to cloud condensation nuclei and form a salt inversion below cloud base (Blanchard & Woodcock 1980).

An analysis of the processes of heat and moisture transfer in spray clouds has resulted in the conclusion (Borisenkov 1974) that it is necessary to solve problems related to the formation and time evolution of the spray cloud and to determination of the thermal regime of an individual droplet. Initially, the droplets have the same properties as the sea surface. After ejection many of them quickly return to the sea but being in the air they attempt to adapt to the air conditions. All spray droplets reach thermal equilibrium within 1 s. The same time is required for the smallest droplets to reach moisture equilibrium, whereas the largest droplets need for this an hour or more (Andreas *et al.* 1995). The evaporation of droplets contributes to the water vapour flux, decreases liquid water flux, and cools a region where the droplet cloud is formed. In particular, it was estimated (Andreas *et al.* 1995) that in a 20 ms^{-1} wind, with an air temperature of 20°C, a water temperature of 22°C, and a relative humidity of 80%, the sea spray contributes 150 Wm^{-2} to the latent heat flux and 15 Wm^{-2} to the sensible heat flux. Additionally, the transfer of radiation between the droplets and the environment can significantly increase the consumption of sensible heat by the droplets. The resulting vertical distribution of droplets and sea salt is one of the factors determining the optical properties of the marine atmosphere. However, from HEXOS field and laboratory experiments it was derived that there is a negative feedback: in a "droplet evaporation layer" close to the surface the evaporating droplets modify the temperature and water vapor gradients and reduce the turbulent sensible and latent heat fluxes (Hasse & Smith 1996).

Effects of sea spray on the wind profile. The sea spray droplets can also influence the density stratification and consequently, the parameters of the surface layer. It was shown (Pielke & Lee 1991) that the water loading effect on the surface-layer wind profile during strong wind conditions can be significant in white-cap sea-spray situations. Thus, the turbulence statistics (especially, the friction velocity and drag coefficient) near the sea surface may be remarkably altered.

The water droplets are embedded into air flow, and if their concentration is large enough, the flow must be considered as multi-phase flow and the drag reduction effects must be included into a parameterization scheme. Assuming that air and droplets form a two-phase fluid, the density ρ of the mixture may be expressed by

$$\rho = \rho_a(1-S) + \rho_w S = \rho_a(1+\sigma S), \tag{3.59}$$

where $S(z)$ is the volume concentration of droplets, and σ indicates the relative excess of the droplet density over the air density:

$$\sigma = (\rho_w - \rho_a)/\rho_a$$

The following expression for the stability parameter L can be derived for these conditions (Wamser & Lykossov 1995):

$$L = \frac{\rho_a(1+\sigma S)u_*^3}{\kappa g[\overline{\rho_a' w'}(1-S) + \rho_a \sigma \overline{S'w'}]}. \tag{3.60}$$

In the absence of droplets ($S \equiv 0$ and $\overline{S'w'} \equiv 0$) Eqs. (3.30) and (3.60) lead to the expression for the stability parameter L derived by Monin & Obukhov (1954):

$$L = -\frac{\overline{\theta_v} u_*^3}{\kappa g \overline{\theta_v' w'}}$$

In the case of thermally neutral stratification ($\overline{\theta_v' w'} \equiv 0$) the parameter L in the presence of droplets has the form:

$$L = \frac{(1+\sigma S)u_*^3}{\kappa g \sigma \overline{S'w'}}. \tag{3.61}$$

Making an eddy diffusion assumption for the concentration flux of droplets (3.3) and using the balance relation of steady state sea spray

$$K_S \frac{\partial S}{\partial z} = -w_f S, \tag{3.62}$$

where $w_f > 0$ is the falling velocity of droplets, the expression (3.61) can be rewritten as follows:

$$L = \frac{(1+\sigma S)u_*^3}{\kappa g w_f \sigma S}, \tag{3.63}$$

It is seen from this equation that the stability parameter is not constant with height, and, consequently, the Monin-Obukhov similarity theory cannot be applied in its traditional form. A comprehensive review of the modified similarity theory based on the height-dependent stability parameter is presented by Stull (1988).

Since L is positive, the density stratification is stable, and one can use the log-linear universal function (Kondo *et al.* 1978). Assuming that $\sigma S << 1$, Eq. (3.28) reads in this case as

$$\frac{\kappa z}{u_*}\frac{\partial u}{\partial z} = 1 + \frac{\beta g \kappa w_f z \sigma S}{u_*^3}. \tag{3.64}$$

The value of the parameter β in this formula is different in this case from that found for scaling with the surface density flux. Kondo *et al.* (1978) found from observations of the atmospheric boundary layer that the most appropriate value of β for the height-depending L equals 7, while a value of $\beta = 4.7$ was derived by Businger *et al.* (1971). The equation for the droplet concentration (3.62) transforms to [2]

$$\frac{\partial S}{\partial z} + \frac{w_f S}{\kappa u_* z}\left(1 + \frac{\beta \kappa g w_f z \sigma S}{u_*^3}\right) = 0. \tag{3.65}$$

The solution to this equation, subject to the boundary condition $S = S_r$ at a reference height $z = z_r$, has the following form (Taylor & Dyer 1977):

$$S(z) = \frac{(1-\omega) S_r (z/z_r)^{-\omega}}{1 - \omega + \alpha\omega^2[(z/z_r)^{1-\omega} - 1]}, \tag{3.66}$$

where

$$\omega = \frac{w_f}{\kappa u_*}, \qquad \alpha = \frac{\beta g \kappa^2 z_r \sigma S_r}{u_*^2}. \tag{3.67}$$

For $\omega \to 1$, the solution (3.66) converges to

$$S(z) = \frac{S_r (z/z_r)^{-1}}{1 + \alpha \ln(z/z_r)}. \tag{3.68}$$

Given the concentration profile (3.66), (3.68), integration of Eq. (3.64) from z_r to z results in the following wind profile (Taylor & Dyer 1977):

$$u(z) = u_r + \frac{u_*}{\kappa}\ln\left(\frac{z}{z_r}\right) + \frac{u_*}{\kappa}\begin{cases} \omega^{-1}\ln\left(1 + \frac{\alpha\omega^2}{1-\omega}\left[\left(\frac{z}{z_r}\right)^{1-\omega} - 1\right]\right) & \text{for } \omega \neq 1, \\ \ln\left[1 + \alpha\ln\left(\frac{z}{z_r}\right)\right], & \text{for } \omega = 1. \end{cases} \tag{3.69}$$

If $z_r = z_0$, $u_r = 0$, and $\alpha = 0$, Eq. (3.69) coincides with Eq. (3.25). Comparing these equations, one can see that effects of the sea spray on the wind profile are described by the additional logarithmic term. Since α and ω are positive, the surface winds should be stronger, for the same friction velocity u_*, in the case of sea spray. Consequently, the drag coefficient should be lower than without spray.

To estimate this effect quantitatively, let us use $u_* = 0.4$ ms^{-1} as a typical value of the friction velocity, $r = 10\mu$m as a typical value of the droplet radius, and $\sigma = 10^3$. Applying the Stokes law, $w_f = 2\sigma g r^2/9\nu$, one can obtain from (3.67) $w_f = 0.016$ ms^{-1} and, consequently, $\omega = 0.1$. Taking $z_r = 0.18$ m, $z_0 = 10^{-4}$m, and assuming that the wind profile between z_0 and z_r is not disturbed by sea spray, one can calculate from (3.69) that for S_r ranging from 10^{-5} to 10^{-4} the wind at 10 m is accelerated about 4 - 35 % with regard to the neutral case without spray. It is nesessary to note, however, that the wave-induced stress, which has an opposite influence on the wind profile, may sufficiently compensate the possible drag reduction effect.

[2] Generally speaking, the eddy diffusivity for momentum and droplets may be different, $K \neq K_S$, but any differences are neglected here.

Effects of rain on the surface stress. There are two aspects of this problem. First, when the rain originates in the boundary layer and/or falls through, it interacts with turbulence and may affect the ageostrophic wind component. Second, the rain falling over the ocean may change the sea surface state and exchange by momentum with it. We consider here this aspect only.

It was found experimentally (Poon *et al.* 1992) that the low-frequency waves are attenuated by the rain, and most of the damping occurs in the frequency region of 2 - 5 Hz. Theoretical studies (e.g. Le Méhauté and Khangaonkar 1990) indicated that the damping rate depends on rain intensity, falling velocity, and inclination of the drops. In particular, observations showed an increase in the damping rate with rain intensity. Reynolds (1900) suggested that the falling raindrops cause the formation of vortex rings which enhance the vertical mixing of the surface water layer. At present time, it is found that under various rain conditions, a mixed layer with depth between 5 and 20 cm is formed with the effective eddy viscosity at least of an order of magnitude greater than the water molecular viscosity (e.g. Katsaros & Buettner 1969; Poon *et al.* 1992). Van Dorn (1953) found that during a period of moderately heavy rainfall (~ 0.5 cm/h) the magnitude of measured stress was significantly larger in comparision with the case without rainfall.

Raindrops move over an ocean at the same velocity as the wind, and when the drops fall and reach the sea surface, they contribute to a rain-induced surface stress. The surface stress τ_r produced by rainfall can be written as follows (Caldwell & Elliott 1971):

$$\tau_r = \rho_w U_s R, \tag{3.70}$$

where ρ_w is the density of rainwater, U_s is the horizontal speed of the rain drop at impact, and R is the rainfall rate. The ratio between the rain-induced and wind stresses at the sea surface was found to be about 7 - 25 % (Poon *et al.* 1992). The contribution of τ_r into the total stress is higher for lower winds and heavier rain conditions. Contrary to the damping mechanism, the rain-induced stress contributes to the wave growth. Thus, in the presence of wind, the increase in damping with the rain intensity is partially compensated by the increase in the surface stress due to rain-induced component.

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