A. Statistical and Spectral Analysis of Turbulence

A.1 Turbulence Properties

Flows qualified as "turbulent" can be found in most fields that make use of fluid mechanics. These flows posses a very complex dynamics whose intimate mechanisms and repercussions on some of their characteristics of interest to the engineer should be understood in order to be able to control them. The criteria for defining a turbulent flow are varied and nebulous because there is no true definition of turbulence. Among the criteria most often retained, we may mention [150]:

– the random character of the spatial and time fluctuations of the velocities, which reflect the existence of finite characteristic scales of statistical correlation (in space and time);
– the velocity field is three-dimensional and rotational;
– the various modes are strongly coupled, which is reflected in the non-linearity of the mathematical model retained (Navier–Stokes equations);
– the large mixing capacity due to the agitation induced by the various scales;
– the chaotic character of the solution, which exhibits a very strong dependency on the initial condition and boundary conditions.

A.2 Foundations of the Statistical Analysis of Turbulence

A.2.1 Motivations

The very great dynamical complexity of turbulent flows makes for a very lengthy deterministic description of them. To analyze and model them, we usually refer to a statistical representation of the fluctuations. This reduces the description to that of the various statistical moments in the solution, which sharply reduces the volume of information. Moreover, the random character of the fluctuations make this approach natural.

A.2.2 Statistical Average: Definition and Properties

We use $\langle \phi \rangle$ to denote the stochastic mean (or statistical average, or mathematical expectation, or ensemble average) of a random variable ϕ calculated from n independent realizations of the same phenomenon $\{\phi_l\}$:

$$\langle \phi \rangle = \lim_{n \to \infty} \frac{1}{n} \sum_{l=1}^{n} \phi_l \quad . \tag{A.1}$$

The turbulent fluctuation ϕ'_l associated with the realization ϕ_l is defined as its deviation from the mathematical expectation:

$$\phi'_l = \phi_l - \langle \phi \rangle \quad . \tag{A.2}$$

By construction, we have the property:

$$\langle \phi' \rangle \equiv 0 \quad . \tag{A.3}$$

On the other hand, fluctuation moments of second or higher order are not necessarily zero. The standard deviation σ can be defined as:

$$\sigma^2 = \langle {\phi'}^2 \rangle \quad . \tag{A.4}$$

We define the *turbulence intensity* as $\sigma / \langle \phi \rangle$.

The correlation at two points in space and two times, $(\boldsymbol{x}, \boldsymbol{x}')$ and (t, t') of the two random variables ϕ and ψ, denoted $R_{\phi\psi}(\boldsymbol{x}, \boldsymbol{x}', t, t')$ is:

$$R_{\phi\psi}(\boldsymbol{x}, \boldsymbol{x}', t, t') = \langle \phi(\boldsymbol{x}, t) \psi(\boldsymbol{x}', t') \rangle \quad . \tag{A.5}$$

A.2.3 Ergodicity Principle

When ϕ is a random steady function in time (i.e. its probability density function is independent of time), we can apply the ergodicity principle according to which it is equivalent, statistically speaking, to consider indefinitely repeated experiments with a single drawing or a single experiment with an infinite number of drawings. We will therefore admit that a single experiment of infinite duration can be considered as representative of all possible scenarios.

The theorem of ergodicity says that the quadratic mean of the random function $\phi_T(t)$ defined by:

$$\phi_T(t) = \frac{1}{T} \int_t^{t+T} \phi(t') dt' \quad , \tag{A.6}$$

converges to a non-random limit equal to the stochastic mean $\langle \phi \rangle$ as $T \to \infty$ only on the condition that:

$$\lim_{T\to\infty} \frac{1}{T} \int_0^T R_{\phi'\phi'}(t)dt = 0 \quad , \tag{A.7}$$

where $R_{\phi'\phi'}(t)$ is the time autocorrelation (or covariance) of the fluctuations of ϕ over time interval t:

$$R_{\phi'\phi'}(t) = \langle (\phi(t') - \langle \phi \rangle)(\phi(t'+t) - \langle \phi \rangle) \rangle \quad . \tag{A.8}$$

For turbulent fluctuations, the random character reflects the fact that $R_{\phi'\phi'}(t) \to 0$ as $t \to \infty$. So if we define the mean in time $\overline{\phi}$ as the limit of ϕ_T as $T \to \infty$, i.e.:

$$\overline{\phi} = \lim_{T\to\infty} \frac{1}{T} \int_0^T \phi(t)dt \quad , \tag{A.9}$$

we get the equality:

$$\overline{\phi} = \langle \phi \rangle \quad . \tag{A.10}$$

We establish that the standard error varies as $1/\sqrt{T}$ for sufficiently large T. Another way of estimating $\langle \phi \rangle$ is to construct the "experimental" average ϕ_n defined as the arithmetic mean from experiments:

$$\phi_n(t) = \frac{1}{n} \sum_{i=1}^{n} \phi_i(t) \quad , \tag{A.11}$$

where the time t is arbitrary since the flow is assumed to be statistically steady. We show that the standard error decreases as $1/\sqrt{n}$ if the experiments ϕ_l are independent.

Let ϕ and ψ be two random variables. The operator $\langle\ \rangle$ thus defined verifies the following properties, sometimes called Reynolds rules:

$$\langle \phi + \psi \rangle = \langle \phi \rangle + \langle \psi \rangle \quad , \tag{A.12}$$

$$\langle a\phi \rangle = a\langle \phi \rangle \quad a = \text{const.} \quad , \tag{A.13}$$

$$\langle \langle \phi \rangle \psi \rangle = \langle \phi \rangle \langle \psi \rangle \quad , \tag{A.14}$$

$$\left\langle \frac{\partial \phi}{\partial s} \right\rangle = \frac{\partial \langle \phi \rangle}{\partial s} \qquad s = \boldsymbol{x}, t \quad , \tag{A.15}$$

$$\left\langle \int \phi(\boldsymbol{x},t) d^3\boldsymbol{x} dt \right\rangle = \int \langle \phi(\boldsymbol{x},t) \rangle d^3\boldsymbol{x} dt \quad . \tag{A.16}$$

Any operator that verifies these properties is called a Reynolds operator. We deduce from these relations the properties:

$$\langle \langle \phi \rangle \rangle = \langle \phi \rangle \quad , \tag{A.17}$$

$$\langle \phi' \rangle = 0 \quad . \tag{A.18}$$

A.2.4 Decomposition of a Turbulent Field

Decomposition Principle. One technique very commonly used for describing a turbulent field is statistical representation. The velocity field at time t and position $\boldsymbol{x}$ splits into:

$$\boldsymbol{u}(\boldsymbol{x},t) = \langle \boldsymbol{u}(\boldsymbol{x},t) \rangle + \boldsymbol{u}'(\boldsymbol{x},t) \quad . \tag{A.19}$$

Using this decomposition and the stochastic mean, we define an evolution equation for the quantity $\langle \boldsymbol{u}(\boldsymbol{x},t) \rangle$. To recover all the information contained in the $\boldsymbol{u}(\boldsymbol{x},t)$ field, we have to handle an infinite set of equations for the statistical moments of it. The quadratic non-linearity of the Navier–Stokes equations induces an intrinsic coupling among the various moments of the solution: the evolution equation of the moment of order n in the solution uses the moment of order $(n+1)$. To recover all the information in the exact solution, it is thus necessary to solve an infinite hierarchy of coupled equations. As this is impossible in practice, this hierarchy is truncated at an arbitrarily chosen level so as to obtain a finite number of equations. This truncation brings out an unknown term that will be modeled using closure hypotheses. If the degree of precision of the information obtained theoretically increases with the number of equations retained, the consequences of the truncation and of the hypotheses used are difficult to predict.

Equations of the Stochastic Moments. The evolution equations of the mean field are obtained by applying the averaging operator to the Navier–Stokes equations. By applying the rules of commutation with the derivation in the case of an incompressible Newtonian fluid and with no external forces, we get

$$\frac{\partial \langle u_i \rangle}{\partial t} + \frac{\partial}{\partial x_j} \langle u_i u_j \rangle = -\frac{\partial \langle p \rangle}{\partial x_i} + \nu \frac{\partial^2 \langle u \rangle}{\partial x_j \partial x_j} \quad , \tag{A.20}$$

$$\frac{\partial \langle u_i \rangle}{\partial x_i} = 0 \quad , \tag{A.21}$$

where ν is the kinematic viscosity. The non-linear term $\langle u_i u_j \rangle$ is unknown and has to be decomposed as a function of $\langle \boldsymbol{u} \rangle$ and $\boldsymbol{u}'$. By introducing relation (A.19) and considering the properties (A.12) to (A.18), we get:

$$\langle u_i u_j \rangle = \langle u_i \rangle \langle u_j \rangle + \langle u_i' u_j' \rangle \quad . \tag{A.22}$$

The last term of the right-hand side, called the Reynolds tensor, is unknown and has to be evaluated. It represents the coupling between the fluctuations and the mean field. This evaluation can be made by solving the corresponding evolution equation, either by employing a model, called closure or turbulence model.

A.2.5 Isotropic Homogeneous Turbulence

Definitions. A field is said to be statistically homogeneous along the parameter x, or imprecisely just "homogeneous", if its statistical moments are independent of the value of x where the measurements are made. This is expressed:

$$\frac{\partial}{\partial x}\langle \phi_1....\phi_n\rangle = 0 \quad . \tag{A.23}$$

A homogeneous field is said to be statistically isotropic (in the Taylor sense), or more simply "isotropic", if all statistical moments relative to a set of points $(x_1, ..., x_n)$ at times $(t_1, ..., t_n)$ remains invariant when the set of n points and the coordinate axis are rotated, and if there is statistical invariance for symmetry about an arbitrary plane.

We may note that there exists an idea of quasi-isotropy introduced by Moffat, which does not require the invariance by symmetry.

A Few Properties. A turbulent field is said to be homogeneous (resp. homogeneous isotropic) if its velocity fluctuation $\boldsymbol{u}'$ is homogeneous (resp. homogeneous isotropic). One necessary condition for achieving homogeneity is that the mean velocity gradient be constant in space:

$$\frac{\partial \langle u_i\rangle}{\partial x_j} = \text{const.} \tag{A.24}$$

Isotropy requires that the mean field $\langle \boldsymbol{u}\rangle$ be zero. When the turbulence is isotropic, only the diagonal elements of the Reynolds tensor are non-zero. Moreover, these are mutually equal:

$$\langle u_i' u_j'\rangle = \frac{2}{3} K \delta_{ij} \quad , \tag{A.25}$$

where K is the turbulent kinetic energy.

A.3 Introduction to Spectral Analysis of the Isotropic Turbulent Fields

A.3.1 Definitions

The tensor of correlations at two points $R_{\alpha\beta}(\boldsymbol{r})$ of a statistically homogeneous vector field $\boldsymbol{u}$ defined as:

$$R_{\alpha\beta}(\boldsymbol{r}) = \langle u_\alpha(\boldsymbol{x}+\boldsymbol{r}) u_\beta(\boldsymbol{x})\rangle \tag{A.26}$$

can be related to a spectral tensor $\Phi_{\alpha\beta}(\boldsymbol{k})$ by the following two relations:

$$R_{\alpha\beta}(\boldsymbol{r}) = \int \Phi_{\alpha\beta}(\boldsymbol{k}) e^{\mathrm{i}k_j r_j} d^3\boldsymbol{k} \quad , \tag{A.27}$$

$$\Phi_{\alpha\beta}(\boldsymbol{k}) = \frac{1}{(2\pi)^3}\int R_{\alpha\beta}(\boldsymbol{r}) e^{-\mathrm{i}k_j r_j} d^3\boldsymbol{k} \quad , \tag{A.28}$$

where $\mathrm{i}^2 = -1$. The tensor at the origin, $R_{\alpha\beta}(0)$, is the Reynolds tensor. In the case of an isotropic field, the general form of the correlation tensor becomes:

$$R_{\alpha\beta}(r) = K\left([f(r) - g(r)]\frac{r_\alpha r_\beta}{r^2} + g(r)\delta_{\alpha\beta}\right) \quad , \tag{A.29}$$

where $f(r)$ and $g(r)$ are two real scalar functions. When the velocity field is solenoidal, these two functions are related by:

$$g(r) = f(r) + \frac{r}{2}\frac{\partial f(r)}{\partial r} \quad . \tag{A.30}$$

The incompressibility constraint also allows us to establish the following relation for the tensor $\Phi_{\alpha\beta}(\boldsymbol{k})$:

$$\Phi_{\alpha\beta}(\boldsymbol{k}) = \frac{E(\boldsymbol{k})}{4\pi k^2}\left(\delta_{\alpha\beta} - \frac{k_\alpha k_\beta}{k^2}\right) \quad , \tag{A.31}$$

where the scalar function $E(\boldsymbol{k})$ is called a three-dimensional spectrum. It represents the contribution of the wave vectors of k to the turbulent kinetic energy, i.e. wave vectors whose tips are included in the region located between two spheres of radius k and $k + dk$. The spectral energy density, denoted $A(\boldsymbol{k})$, is therefore equal to $E(\boldsymbol{k})/4\pi k^2$. The three-dimensional spectrum is computed from the spectral tensor by integration over the sphere of radius k:

$$E(k) = \frac{1}{2}\int \Phi_{ii}(\boldsymbol{k})dS(\boldsymbol{k}) \quad , \tag{A.32}$$

where $dS(\boldsymbol{k})$ is the integration element on the sphere of radius k. This quantity can also be related to the function $f(r)$ by the relation:

$$E(k) = \frac{K}{\pi}\int_0^\infty kr\left(\sin(kr) - kr\cos(kr)\right)f(r)dr \quad . \tag{A.33}$$

The turbulent kinetic energy, K, is found by summation over the entire spectrum:

$$K \equiv \frac{\langle u_i' u_i' \rangle}{2} = \int_0^\infty E(\boldsymbol{k})d^3\boldsymbol{k} \quad . \tag{A.34}$$

By construction, the spectral tensor has the property:

$$\Phi_{ij}(-\boldsymbol{k}) = \Phi_{ij}^*(\boldsymbol{k}) \quad , \tag{A.35}$$

where the asterisk indicates the complex conjugate number. The homogeneity property of the turbulent field implies:

$$\Phi_{ij}(\boldsymbol{k}) = \Phi_{ji}^*(\boldsymbol{k}) \quad . \tag{A.36}$$

The spectral tensor can also be related to the velocity fluctuation $\boldsymbol{u}'$ and to its Fourier transform $\widehat{\boldsymbol{u}}'$ defined as:

$$\widehat{\boldsymbol{u}}_i'(\boldsymbol{k}) = \frac{1}{(2\pi)^3}\int \boldsymbol{u}'(\boldsymbol{x})e^{-\imath k_j x_j}d^3\boldsymbol{x} \quad . \tag{A.37}$$

Simple expansions lead to the equality:

$$\langle \widehat{u}_i'(\boldsymbol{k}')\widehat{u}_j'(\boldsymbol{k})\rangle = \delta(\boldsymbol{k}+\boldsymbol{k}')\Phi_{ij}(\boldsymbol{k}) \quad . \tag{A.38}$$

So we see that the two modes are correlated statistically only if $\boldsymbol{k}+\boldsymbol{k}'=0$. An equivalent definition of the spectral tensor is:

$$\Phi_{ij}(\boldsymbol{k}) = \int \langle \widehat{u}_i'^*(\boldsymbol{k})\widehat{u}_j'(\boldsymbol{k}')\rangle d^3\boldsymbol{k}' \quad . \tag{A.39}$$

A.3.2 Modal Interactions

The nature of the interactions among the various modes can be brought out by analyzing the non-linear term that appears in the evolution equation associated with them. This equation, for the mode associated with the wave vector $\boldsymbol{k}$ (the dependency on $\boldsymbol{k}$ is not expressed, for the sake of simplicity) is:

$$\frac{\partial \widehat{u}_i}{\partial t} + \imath k_j \widehat{a}_{ij} = -\imath k_i \widehat{p} - \nu k^2 \widehat{u}_i \quad . \tag{A.40}$$

The two quantities $\widehat{a}_{ij}$ and $\widehat{p}$ are related to $u_i u_j$ and the pressure p by the relations:

$$u_i(\boldsymbol{x})u_j(\boldsymbol{x}) = \int \widehat{a}_{ij}(\boldsymbol{k})e^{\imath k_l x_l}d^3\boldsymbol{k} \quad , \tag{A.41}$$

$$\frac{1}{\rho}p(\boldsymbol{x}) = \int \widehat{p}(\boldsymbol{k})e^{\imath k_l x_l}d^3\boldsymbol{k} \quad . \tag{A.42}$$

By introducing the spectral decompositions:

$$u_i(\boldsymbol{x}) = \int \widehat{u}_i(\boldsymbol{k}')e^{\imath k_l' x_l}d^3\boldsymbol{k}' \quad , \tag{A.43}$$

$$u_j(\boldsymbol{x}) = \int \widehat{u}_j(\boldsymbol{k}'')e^{\imath k_l'' x_l}d^3\boldsymbol{k}'' \quad , \tag{A.44}$$

the non-linear term becomes:

$$u_i(\boldsymbol{x})u_j(\boldsymbol{x}) = \int \underbrace{\int \widehat{u}_i(\boldsymbol{k}')\widehat{u}_j(\boldsymbol{k}-\boldsymbol{k}')d^3\boldsymbol{k}'}_{\widehat{a}_{ij}(\boldsymbol{k})} e^{\imath k_l x_l}d^3\boldsymbol{k} \quad , \tag{A.45}$$

where we have performed the variable change $\boldsymbol{k}=\boldsymbol{k}'+\boldsymbol{k}''$. The pressure term is computed by the Poisson equation:

$$\frac{1}{\rho}\frac{\partial^2 p}{\partial x_i \partial x_i} = -\frac{\partial^2 u_i u_j}{\partial x_i \partial x_j} \quad , \tag{A.46}$$

or, in the spectral space:

$$k^2 \widehat{p} = -k_l k_m \widehat{a}_{lm} \quad . \tag{A.47}$$

The momentum equation therefore takes the form:

$$\left[\frac{\partial}{\partial t} + \nu k^2\right] \widehat{u}_i(\boldsymbol{k}) = M_{ijm}(\boldsymbol{k}) \int \widehat{u}_m(\boldsymbol{k'}) \widehat{u}_j(\boldsymbol{k} - \boldsymbol{k'}) d^3\boldsymbol{k'} \quad , \tag{A.48}$$

in which

$$M_{ijm}(\boldsymbol{k}) = -\frac{\imath}{2} \left(k_m P_{ij}(\boldsymbol{k}) + k_j P_{im}(\boldsymbol{k})\right) \quad , \tag{A.49}$$

where $P_{ij}(\boldsymbol{k})$ is the projection operator on the plane orthogonal to the vector $\boldsymbol{k}$. This operator is expressed:

$$P_{ij}(\boldsymbol{k}) = \left(\delta_{ij} - \frac{k_i k_j}{k^2}\right) \quad . \tag{A.50}$$

The linear terms are grouped into the left-hand side and the non-linear terms in the right. The first linear term represents the time dependency and the second the viscous effects. The non-linear term represents the effect of convection and pressure. We can see that the mode $\boldsymbol{k}$ interacts with the modes $\boldsymbol{p} = \boldsymbol{k'}$ and $\boldsymbol{q} = (\boldsymbol{k} - \boldsymbol{k'})$ such that $\boldsymbol{k} + \boldsymbol{p} = \boldsymbol{q}$. This triadic nature of the non-linear interactions is intrinsically related to the mathematical structure of the Navier–Stokes equations.

A.3.3 Spectral Equations

The equations for the spectral tensor components Φ_{ij} are obtained by applying an inverse Fourier transform to the transport equations of the two-point double correlations. After computation, we get:

$$\begin{aligned}\frac{\partial \Phi_{ij}}{\partial t} - \lambda_{lm} k_l \frac{\partial \Phi_{ij}}{\partial k_m} \quad + \quad & \lambda_{il}\Phi_{lj} + \lambda_{jl}\Phi_{il} + 2\nu k^2 \Phi_{ij} = \\ & \left(k_l \Theta_{ilj} + k_l \Theta^*_{jli}\right) + \left(k_i \Sigma_j + k_j \Sigma^*_j\right) \quad , \end{aligned} \tag{A.51}$$

where:

$$\Theta_{ilj} = \frac{\imath}{(2\pi)^3} \int \langle u'_i(\boldsymbol{x}) u'_l(\boldsymbol{x}) u'_j(\boldsymbol{x} + \boldsymbol{r})\rangle e^{-\imath k_n r_n} d^3\boldsymbol{r} \quad , \tag{A.52}$$

$$\Sigma_j = \frac{\imath}{(2\pi)^3} \int \frac{1}{\rho} \langle p'(\boldsymbol{x}) u'_j(\boldsymbol{x} + \boldsymbol{r})\rangle e^{-\imath k_n r_n} d^3\boldsymbol{r} \tag{A.53}$$

$$= 2\lambda_{lm} \frac{k_l}{k^2} \Phi_{mj} - \frac{k_l k_m}{k^2} \Theta_{mlj} \quad , \tag{A.54}$$

$$\lambda_{ij} = \frac{\partial \langle u \rangle_i}{\partial x_j} \quad . \tag{A.55}$$

By expanding the terms (A.52) and (A.54), equation (A.51) takes the form:

$$\begin{aligned}\left(\frac{\partial}{\partial t}+2\nu k^2\right)\Phi_{ij}(\boldsymbol{k}) &+ \frac{\partial\langle u_i\rangle}{\partial x_l}\Phi_{jl}(\boldsymbol{k})+\frac{\partial\langle u_j\rangle}{\partial x_l}\Phi_{il}(\boldsymbol{k})\\ &- 2\frac{\partial\langle u_l\rangle}{\partial x_m}\left(k_i\Phi_{jm}(\boldsymbol{k})+k_j\Phi_{mi}(\boldsymbol{k})\right)\\ &- \frac{\partial\langle u_l\rangle}{\partial x_m}\frac{\partial}{\partial k_m}\left(k_l\Phi_{ij}(\boldsymbol{k})\right)\\ &= P_{il}(\boldsymbol{k})T_{lj}(\boldsymbol{k})+P_{jl}(\boldsymbol{k})T^*_{li}(\boldsymbol{k}) \quad , \end{aligned} \qquad (A.56)$$

where

$$T_{ij}(\boldsymbol{k}) = k_l \int\int\int \langle u_i(\boldsymbol{k})u_l(\boldsymbol{p})u_j(-\boldsymbol{k}-\boldsymbol{p})\rangle d^3\boldsymbol{p} \quad . \qquad (A.57)$$

The evolution equation for the energy spectrum $E(k)$, derived from (A.51) by integration over the sphere of radius k, is:

$$\frac{\partial E(k)}{\partial t} = P(k)+T(k)+D(k) \quad , \qquad (A.58)$$

where the kinetic energy production term $P(k)$ by interaction with the mean field, the transfer term $T(k)$ and the dissipation term $D(k)$ are given by:

$$P(k) = -\lambda_{ij}\phi_{ij}(k) \quad , \qquad (A.59)$$

$$T(k) = \frac{1}{2}\int\left(k_l(\Theta_{ili}+\Theta^*_{ili})+\lambda_{lm}\frac{\partial(k_l\phi_{ii})}{\partial k_m}\right)dS(\boldsymbol{k}) \quad , \qquad (A.60)$$

$$D(k) = -2\nu k^2 E(k) \quad , \qquad (A.61)$$

where the tensor $\phi_{ij}(k)$ is defined as the integral of $\Phi_{ij}(\boldsymbol{k})$ over the sphere of radius k:

$$\phi_{ij}(k) = \int \Phi_{ij}(\boldsymbol{k})dS(\boldsymbol{k}) \quad . \qquad (A.62)$$

The kinetic energy conservation property for ideal fluid is expressed by:

$$\int_0^\infty T(k)dk = 0 \quad . \qquad (A.63)$$

We come up with the kinetic energy evolution equation in the physical space by integrating (A.58) over the entire spectrum:

$$\frac{\partial K}{\partial t} = \int_0^\infty \frac{\partial E(k)}{\partial t}dk = \int_0^\infty P(k)dk+\int_0^\infty T(k)dk+\int_0^\infty D(k)dk \quad . \qquad (A.64)$$

In the isotropic homogeneous case, production is zero and we get:

$$\frac{\partial K}{\partial t} = -\varepsilon \quad , \qquad (A.65)$$

where the kinetic energy dissipation rate ε is given by:

$$\varepsilon = \int_0^\infty 2\nu k^2 E(k)dk \quad . \qquad (A.66)$$

A.4 Characteristic Scales of Turbulence

Several characteristic scales of turbulence can be defined. We define the integral scale L_{ij}^l as:

$$L_{ij}^l = \int_{-\infty}^{+\infty} R_{ij}(r) dr \quad . \tag{A.67}$$

This scale is representative of the length over which the turbulent fluctuations are mutually correlated, so it is directly related to the size of the structures that form the turbulent field. Another scale, called the Taylor microscale, is denoted λ_τ and is defined as:

$$\frac{\langle u'^2 \rangle}{\lambda_\tau^2} = \left\langle \left(\frac{\partial u'}{\partial x} \right)^2 \right\rangle \quad . \tag{A.68}$$

While the first scale is associated with all the turbulent structures, this latter scale is related directly to the small scales of the turbulence. Considering that the dissipation , can be written as:

$$\varepsilon = 2\nu \left\langle \left(\frac{\partial u'}{\partial x} \right)^2 \right\rangle \quad , \tag{A.69}$$

we get the relation:

$$\varepsilon = 2\nu \frac{\langle u'^2 \rangle}{\lambda_\tau^2} \quad . \tag{A.70}$$

The Taylor microscale thus appears as characteristic of the dissipative phenomena.

A.5 Spectral Dynamics of Isotropic Homogeneous Turbulence

A.5.1 Energy Cascade and Local Isotropy

Analyses on the basis of (A.58) show that the dynamical mechanism associated with the term $T(k)$ is a kinetic energy transfer from the small wave numbers to the large. This process is called the energy cascade. It is relatively local in frequency: the transfers are negligible among wave numbers separated by more than two decades. It repeats itself until such time as the structures are so small that the viscous mechanisms, represented by the $D(k)$ term, become preponderant.

The local isotropy hypothesis formalized by Kolmogorov assumes statistical homogeneity and isotropy in a small space–time domain and not throughout the flow. This is equivalent to the hypothesis that the flow scales, while sufficiently small, are governed by a dynamics similar to that of isotropic homogeneous turbulence. They are thus independent of the large scales and their statistical structure acquires a universal character.

Kolmogorov's first hypothesis is that the statistical moments of the scales located in such a domain depend only on the separation distance r, the total dissipation by viscosity per unit mass ε, and the viscosity ν.

The second hypothesis is that the statistical moments for large separation distances become independent of the viscosity and are no longer a function of r and ε.

A.5.2 Equilibrium Spectrum

With the hypothesis of local isotropy, we find three distinct regions in the energy spectrum $E(k)$:

– the large scale region where the turbulence associated with the $P(k)$ term is produced. These scales are coupled with the mean field and are affected by the boundary conditions, so they possess no universal character. However, following arguments related to the finite character of the energy spectral density $A(k)$, we can say that:

$$E(k) \simeq k^4 \text{ or } E(k) \simeq k^2 \text{ for } k \ll 1 \tag{A.71}$$

– the inertial range, associated with the intermediate scales, in which the energy is transferred by non-linear interaction with no action by viscosity or production. The energy spectrum depends only on k and ε. Since the energy is transferred without loss, ε remains constant. Assuming that there exists a self-similar form of the power-law spectrum, by dimensional arguments we get:

$$E(k) = K_0 \varepsilon^{2/3} k^{-5/3} \quad , \tag{A.72}$$

where the constant K_0, called the Kolmogorov constant, is close to 1.5.

– The dissipation region, which comprises the smallest scales where the kinetic energy is dissipated by the viscous effects. In this area, the relaxation time τ_d associated with the viscous effects is at least equal to that of the non-linear transfers, denoted τ_c. For a length scale l, these two times are evaluated as:

$$\tau_\mathrm{d} \approx l^2/\nu, \qquad \tau_\mathrm{c} \approx \nu^2/\varepsilon \quad . \tag{A.73}$$

The dissipation region is characterized by the relation:

$$\tau_\mathrm{d} \leq \tau_\mathrm{c} \Longrightarrow l \leq \sqrt{\nu^3/\varepsilon} \quad . \tag{A.74}$$

We call the Kolmogorov scale, which is denoted η_K, the scale for which these two times are equal and which marks the beginning of the dissipation region:

$$\eta_\mathrm{K} = \sqrt{\nu^3/\varepsilon} \quad . \tag{A.75}$$

The characteristic velocity associated with this scale is:

$$v_{\mathrm{K}} = (\nu\varepsilon)^{1/4} \quad . \tag{A.76}$$

The energy spectrum depends explicitly only on k, ν and ε, or equivalently on k, the Kolmogorov scales η_{K}, and v_{K}. The dimensional arguments do not lead to a unique form $E(k)$, and several solutions have been proposed. Arguments concerning the regularity of the velocity field and these gradients suggest an exponential decay of $E(k)$ in this region.

B. EDQNM Modeling

The EDQNM model is briefly described here in its isotropic and anisotropic versions. For more details concerning the isotropic version, the reader may refer to the work of Lesieur [439].

B.1 Isotropic EDQNM Model

Starting with the Navier–Stokes equations written in symbolic form:

$$\left(\frac{\partial}{\partial t} + \nu k^2\right) u = uu \quad , \tag{B.1}$$

we derive an infinite hierarchical set of evolution equations as usual for the statistical moments of the velocity u:

$$\begin{aligned} \left(\frac{\partial}{\partial t} + \nu k^2\right) \langle uu \rangle &= \langle uuu \rangle \quad , & \text{(B.2)} \\ \left(\frac{\partial}{\partial t} + \nu (k^2 + p^2 + q^2)\right) \langle uuu \rangle &= \langle uuuu \rangle \quad , & \text{(B.3)} \\ \dots &= \dots \quad , \end{aligned}$$

where the symbol $\langle \rangle$ designates a statistical average. We then adopt the following hypothesis.

Hypothesis B.1 (Quasi-normality Hypothesis) *The velocity distribution law is close to the Gaussian bell curve and its fourth-order cumulant, denoted $\langle uuuu \rangle_c$ is zero.*

The evolution equation of the triple correlations then become:

$$\left(\frac{\partial}{\partial t} + \nu (k^2 + p^2 + q^2)\right) \langle uuu \rangle = \sum \langle uu \rangle \langle uu \rangle \quad . \tag{B.4}$$

The quasi-normal approximation does not provide the realizability condition, i.e. the spectrum $E(k)$ can take negative values. To recover this property, Orszag proposes introducing a triple-correlation damping term and then gets:

$$\left(\frac{\partial}{\partial t} + \nu (k^2 + p^2 + q^2)\right) \langle uuu \rangle = \sum \langle uu \rangle \langle uu \rangle - (\eta_k + \eta_p + \eta_q) \langle uuu \rangle \quad . \tag{B.5}$$

The solution to this equation is:

$$\langle uuu\rangle(t) = \int_0^t \sum \langle uu\rangle\langle uu\rangle e^{-(\mu_k+\mu_p+\mu_q)t} dt \quad , \tag{B.6}$$

with

$$\mu_k = \eta_k + \nu k^2 \quad . \tag{B.7}$$

To get a solution that is easier to calculate, we adopt the following hypothesis:

Hypothesis B.2 *The relaxation time of the triple correlations is small compared with the relaxation time of the double correlations (which is also that of the energy spectrum).*

With this hypothesis, we can "Markovianize" equation (B.6), which leads to:

$$\begin{aligned}\langle uuu\rangle(t) &= \sum \langle uu\rangle\langle uu\rangle \int_0^t e^{-(\mu_k+\mu_p+\mu_q)t} dt \\ &= \sum \langle uu\rangle\langle uu\rangle \frac{1 - e^{-(\mu_k+\mu_p+\mu_q)t}}{\mu_k + \mu_p + \mu_q} \quad .\end{aligned} \tag{B.8}$$

This relation closes the derivation equation of the second-order moments. This closure is equivalent to replacing the solution of the Navier–Stokes equations with that of the following Langevin-type stochastic model:

$$\left(\frac{\partial}{\partial t} + (\nu + \eta(k,t))k^2\right) u = f(k,t) \quad , \tag{B.9}$$

in which

$$\eta(k,t) = \frac{1}{2} \int\int \Theta_{kpq}(t) \frac{p}{kq} b_{kpq} E(q,t) dp dq \quad . \tag{B.10}$$

The forcing term f is such that:

$$\begin{aligned}F(k,t) &\equiv 4\pi k^2 \int_0^t \langle f(k,t) f(k,s)\rangle_{|k|=cste} ds \\ &= \int\int \Theta_{kpq}(t) \frac{k^3}{pq} a_{kpq} E(p,t) E(q,t) dp dq \quad ,\end{aligned} \tag{B.11}$$

where a_{kpq} and b_{kpq} are coefficients linked to the geometry of the triad $(\boldsymbol{k}, \boldsymbol{p}, \boldsymbol{q})$, defined as:

$$a_{kpq} = \frac{1}{2}(1 - xyz - 2y^2 z^2), \; b_{kpq} = \frac{p}{k}(xy + z^3) \quad , \tag{B.12}$$

where x, y and z are the cosines of the angles of the triangle formed by the wave vectors $(\boldsymbol{k}, \boldsymbol{p}, \boldsymbol{q})$, opposed respectively to $\boldsymbol{k}$, $\boldsymbol{p}$ and $\boldsymbol{q}$. The relaxation time $\Theta_{kpq}(t)$ is evaluated as:

$$\Theta_{kpq}(t) = \frac{1 - e^{-(\mu_k + \mu_p + \mu_q)t}}{\mu_k + \mu_p + \mu_q} \quad , \tag{B.13}$$

where the damping factor μ_k is chosen as follows:

$$\mu_k = \nu k^2 + 0.36 \sqrt{\int_0^k p^2 E(p) dp} \quad . \tag{B.14}$$

B.2 Cambon's Anisotropic EDQNM Model

To study anisotropic homogeneous flows, we define the following two spectral tensors [96] (see Appendix A):

$$\Phi_{ij}(\boldsymbol{k}) \simeq \langle \widehat{u}_i'^*(\boldsymbol{k}) \widehat{u}_j'(\boldsymbol{k}) \rangle \quad , \tag{B.15}$$

$$\phi_{ij}(k) = \int \Phi_{ij}(\boldsymbol{k}) dA(\boldsymbol{k}) \quad . \tag{B.16}$$

The evolution equation for the quantities $\phi_{ij}(k)$ in the presence of mean velocity gradients is:

$$\begin{aligned} \left(\frac{\partial}{\partial t} + 2\nu k^2 \right) \phi_{ij}(k) &= -\frac{\partial \langle u \rangle_i}{\partial x_k} \phi_{jl}(k) - \frac{\partial \langle u \rangle_j}{\partial x_l} \phi_{il}(k) \\ &+ P_{ij}^{\mathrm{l}}(k) + S_{ij}^{\mathrm{l}}(k) \\ &+ P_{ij}^{\mathrm{nl}}(k) + S_{ij}^{\mathrm{nl}}(k) \quad , \end{aligned} \tag{B.17}$$

where

$$P_{ij}^{\mathrm{l}}(k) = 2 \frac{\partial \langle u \rangle_l}{\partial x_m} \int \frac{k_l}{k^2} \left[k_i \Phi_{mj}(\boldsymbol{k}) + k_j \Phi_{mi}(\boldsymbol{k}) \right] dA(\boldsymbol{k}) \quad , \tag{B.18}$$

$$S_{ij}^{\mathrm{l}}(k) = \frac{\partial \langle u \rangle_l}{\partial x_m} \int \frac{\partial}{\partial k_m} \left(k_l \Phi_{ij}(\boldsymbol{k}) \right) dA(\boldsymbol{k}) \quad , \tag{B.19}$$

$$P_{ij}^{\mathrm{nl}}(k) = - \int \frac{k_l}{k^2} \left[k_i T_{lj}(\boldsymbol{k}) + k_j T_{li}^*(\boldsymbol{k}) \right] dA(\boldsymbol{k}) \quad , \tag{B.20}$$

$$S_{ij}^{\mathrm{nl}}(k) = \int \left[T_{ij}(\boldsymbol{k}) + T_{ij}^*(\boldsymbol{k}) \right] dA(\boldsymbol{k}) \quad . \tag{B.21}$$

Equation (B.17) is closed by replacing $\Phi(\boldsymbol{k})$ with a modeled form as a function of $\phi(k)$, where the direction of k is controlled in the integrals (B.18) to (B.21):

$$\begin{aligned} P^{\mathrm{l}}_{ij}(k) &= 2E(k)\left[\frac{2}{5}\langle S\rangle_{ij}\right. \\ &- 3D\left(\langle S\rangle_{lj}H_{li}(k)+\langle S\rangle_{li}H_{lj}(k)-\frac{2}{3}\delta_{ij}\langle S\rangle_{lm}H_{lm}(k)\right) \\ &+ \left.\frac{14}{3}\left(D+\frac{4}{7}\right)\left(\langle \Omega\rangle_{il}H_{lj}(k)+\langle \Omega\rangle_{jl}H_{li}(k)\right)\right] \quad , \end{aligned} \tag{B.22}$$

$$\begin{aligned} S^{\mathrm{l}}_{ij}(k) &= -\frac{2}{15}\langle S\rangle_{ij}\frac{\partial}{\partial k}\left(kE(k)\right)+2\langle S\rangle_{il}\frac{\partial}{\partial k}\left(kDE(k)H_{jl}(k)\right) \\ &+ 2\langle S\rangle_{jl}\frac{\partial}{\partial k}\left(kDE(k)H_{il}(k)\right) \\ &- \frac{1}{3}\delta_{ij}\langle S\rangle_{lm}\frac{\partial}{\partial k}\left([2+11D]\,kE(k)H_{lm}(k)\right) \quad , \end{aligned} \tag{B.23}$$

$$\begin{aligned} P^{\mathrm{nl}}_{ij}(k) &= \int\!\!\int \Theta_{kpq}\frac{2}{pq}(x+yz)H_{ij}(q)\Big[k^1pE(p)E(q)\Big(y(z^2-y^2)(a(q)+3) \\ &+ (y+xz)\frac{a(q)}{5}\Big)-p^3E(k)E(q)y(z^2-x^2)(a(q)+3)\Big]dpdq\,, \end{aligned} \tag{B.24}$$

$$\begin{aligned} S^{\mathrm{nl}}_{ij}(k) &= \int\!\!\int \Theta_{kpq}\frac{2}{pq}\left[(xy+z^3)\left(k^2pE(p)E(q)\left\{\frac{1}{3}\delta_{ij}+H_{ij}(p)\right.\right.\right. \\ &\left.\left. +H_{ij}(q)\right\}-p^3E(k)E(q)\left\{\frac{1}{3}\delta_{ij}+H_{ij}(k)+H_{ij}(q)\right\}\right) \\ &+ \left. H_{ij}(q)\ \left(k^2pE(p)E(q)c_{kpq}-p^3E(k)E(q)c_{pkq}\right)\right]dpdq\,, \end{aligned} \tag{B.25}$$

where

$$\langle S\rangle=\frac{1}{2}\left(\nabla\langle\boldsymbol{u}\rangle+\nabla^T\langle\boldsymbol{u}\rangle\right), \qquad \langle\Omega\rangle=\frac{1}{2}\left(\nabla\langle\boldsymbol{u}\rangle-\nabla^T\langle\boldsymbol{u}\rangle\right) \quad ,$$

and x, y and z are the cosines of the interior angles opposite the wave vectors $\boldsymbol{k}$, $\boldsymbol{p}$ and $\boldsymbol{q}$, respectively, in the triangle formed by these vectors. The anisotropy parameter $a(k)$ is optimized by the Rapid Distortion Theory. The factor D is defined as:

$$D=\frac{2}{7}\left(1+\frac{4}{5}a\right) \quad . \tag{B.26}$$

The energy and anisotropy spectra, denoted respectively $E(k)$ and $H_{ij}(k)$, are given by the relations:

$$E(k) = \frac{1}{2}\phi_{ll}(k) \quad , \tag{B.27}$$

$$H_{ij}(k) = \frac{\phi_{ij}(k)}{2E(k)} - \frac{1}{3}\delta_{ij} \quad . \tag{B.28}$$

The geometric factor c_{kpq} is defined as:

$$c_{kpq} = \frac{1}{2}(xy+z)\left[(y^2-z^2)(a(q)+3) + \frac{2}{5}a(q)(1+z^2)\right] \quad . \tag{B.29}$$

The relaxation time $\Theta_{kpq}(t)$ is evaluated as:

$$\Theta_{kpq}(t) = \frac{1-e^{-(\mu_k+\mu_p+\mu_q)t}}{\mu_k+\mu_p+\mu_q} \quad , \tag{B.30}$$

where the damping term μ_k is:

$$\mu_k = \nu k^2 + 0,36\left(\int_0^k p^2 E(p)dp + \langle\Omega\rangle_{ij}\langle\Omega\rangle_{ij}\right)^{1/2} \quad . \tag{B.31}$$

It should be noted that fully anisotropic versions, which call for no angular parameter setting, have been proposed and compared with simulations for the case of pure rotation and stable stratification [97].

B.3 EDQNM Model for Isotropic Passive Scalar

The transport equation for the passive scalar (15.1) can be rewritten in the Fourier space under the following symbolic form:

$$\left(\frac{\partial}{\partial t} + \kappa k^2\right)\theta = u\theta \quad , \tag{B.32}$$

leading to a hierarchy of symbolic equtions for the statistical moments:

$$\left(\frac{\partial}{\partial t} + \kappa(k^2+p^2)\right)\langle u\theta\rangle = \langle\theta\theta u\rangle \quad , \tag{B.33}$$

$$\left(\frac{\partial}{\partial t} + [\kappa(k^2+p^2)+\nu q^2]\right)\langle\theta\theta u\rangle = \langle\theta\theta uu\rangle \quad , \tag{B.34}$$

$$... = ... \tag{B.35}$$

The EDQNM procedure can be applied to this set of equations to obtain a simple, closed equation for the temperature spectrum $E_\theta(k)$. Details will not be given here, since the process that yields the closed EDQNM model involves the same steps as in the isotropic model for the velocity moments discussed above. The reader can refer to Lesieur's book [439] for a detailed presentation. The model is expressed as

$$\left(\frac{\partial}{\partial t} + 2\kappa k^2\right) E_\theta(k) = \int\int \Theta^\theta_{kpq} \frac{k}{pq}(1-y^2)E(q)[k^2 E_\theta(p) - p^2 E_\theta(k)] \times \delta(\boldsymbol{k}-\boldsymbol{p}-\boldsymbol{q})dpdq \quad , \tag{B.36}$$

where the triad geometrical parameter is the same as in above sections, and the characteristic time is defined as

$$\Theta^\theta_{kpq} = \frac{1 - e^{-[\kappa(k^2+p^2)+\nu q^2+\mu_1(k)+\mu_1(p)+\mu_2(q)]t}}{\kappa(k^2+p^2)+\nu q^2+\mu_1(k)+\mu_1(p)+\mu_2(q)} \quad , \tag{B.37}$$

where

$$\mu_i(k) = a_i \sqrt{\int_0^k p^2 E(p)dp}, \quad i = 1, 2 \quad . \tag{B.38}$$

The two constants a_1 and a_2 are not uniquely defined, and several choices are possible. Setting $a_1 = 0$ and $a_2 = 1.3$ enables the solve analytically the system and yields a satisfactory prediction from a physical point of view. Another possibility is to take $a_1 = a_2 = 0.218 K_0^{3/2}$.

Anisotropic passive scalar models [303] and models for stably stratified flows [266, 689, 267] are very complex and will not be discussed here. This is also the case of EDQNM models for mixed scalars [723].

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