J. Fluid Mech. (1993), *vol.* 248, *pp.* 493–511

The sweeping decorrelation hypothesis and energy–inertial scale interaction in high Reynolds number flows

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(Received 29 June 1992)

The random sweeping decorrelation hypothesis was analysed theoretically and experimentally in terms of the higher-order velocity structure functions $D_{u_i}^{(m)}(r) = \langle [u_i^m(x+r) - u_i^m(x)]^2 \rangle$. Measurements in two high Reynolds number laboratory shear flows were used: in the return channel ($R_\lambda \approx 3.2 \times 10^3$) and in the mixing layer ($R_\lambda \approx 2.0 \times 10^3$) of a large wind tunnel. Two velocity components (in the direction of the mean flow, u_1, and in the direction of the mean shear, u_2) were processed for $m = 1$–4. The effect of using Taylor's hypothesis was estimated by a specially developed method, and found to be insignificant. It was found that all the higher-order structure functions scale, in the inertial subrange, as $r^{\frac{2}{3}}$. Such a scaling has been argued as supporting evidence for the sweeping hypothesis. However, our experiments also established a strong correlation between energy- and inertial-range excitation. This finding leads to the conclusion that the sweeping decorrelation hypothesis cannot be exactly valid.

The hypothesis of statistical independence of large- and small-scale excitation was directly checked with conditionally averaged moments of the velocity difference $\langle [u_i(x+r) - u_i(x)]^l \rangle_{u_i^*}$, $l = 2$–4, at a fixed value of the large-scale parameter u_i^*. Clear dependence of the conditionally averaged moments on the level of averaging was found. In spite of a strong correlation between the energy-containing and the inertial-scale excitation, universality of the intrinsic structure of the inertial subrange was shown.

1. Introduction

The random sweeping decorrelation hypothesis was first investigated by Tennekes (1975) who assumed that the small-scale eddies (with scales at least one order of magnitude less than those of energy-containing eddies) are advected past the Eulerian observer by the energy-containing eddies without dynamical distortion. Tennekes showed that if large-scale advection gives the dominant contribution to kinetic energy in the inertial subrange, the Eulerian frequency spectra of any velocity component square u_i^2 may be described as

$$\Phi_{u_i}^{(2)}(f) = \beta_{u_i} \epsilon^{\frac{2}{3}} \langle u_i^2 \rangle f^{-\frac{5}{3}}, \tag{1}$$

where β_{u_i} is an unknown constant, f is a frequency, ϵ is a mean energy dissipation rate. Hereafter u_i $(i = 1, 2, 3)$ denote velocity fluctuation components, i.e. $\langle u_i \rangle = 0$, and $\int \Phi_{u_i}^{(2)}(f)\, \mathrm{d}f = \langle u_i^4 \rangle$.

The same problem was studied by Van Atta & Wyngaard (1975) who analysed one of the earliest attempts, by Dutton & Deaven (1972), to extend Kolmogorov's (1941*a, b*) scaling to the higher-order spectra. Van Atta & Wyngaard assumed Gaussian statistics for velocity components u_i and found that, for the inertial subrange,

$$E^{(n)}_{u_i}(k) = n^2 \times 3 \times 5 \times \ldots (2n-3)\langle u_i^2\rangle^{n-1} E^{(1)}_{u_i}(k), \quad n \geqslant 2, \tag{2}$$

where $k = 2\pi f/U_1$ is a wavenumber, U_1 is a mean longitudinal velocity, and $\int E^{(n)}_{u_i}(k)\,\mathrm{d}k = \langle u_i^{2n}\rangle$. Now according to the Kolmogorov (1941*a, b*) and Oboukhov (1941) theory, the energy spectrum scales as

$$E^{(1)}(k) \sim k^{-\frac{5}{3}},$$

so that (2) shows that

$$E^{(n)}_{u_i}(k) \sim k^{-\frac{5}{3}}, \tag{3}$$

i.e. the higher-order spectra in the inertial subrange have the same power-law exponents as that of the energy spectrum. This result was strongly supported with experimental data by Van Atta & Wyngaard (1975) for $n \leqslant 9$.

The interest in high-order moments of the velocity field (especially the fourth-order one) was renewed recently in connection with renormalization group (RNG) results by Yakhot, Orszag & She (1989) who obtained

$$E^{(2)}_{u_i}(k) \sim k^{-\frac{7}{3}}, \tag{4}$$

i.e. the spectra of both kinetic energy and pressure fluctuations scale as $k^{-\frac{7}{3}}$ in the inertial subrange. This conclusion can be also obtained from a straightforward extension of Kolmogorov's dimensional analysis to higher-order spectra, if the mean dissipation rate ϵ is assumed as the only governing parameter.

The question was further examined by Nelkin & Tabor (1990). They showed that the random sweeping is dominant if the kinetic energy spectrum scales as $k^{-\frac{5}{3}}$ in the inertial subrange. On the other hand, if the kinetic energy spectrum scales as (4), i.e. has the same exponent as the pressure-fluctuation spectrum, the RNG prediction of no sweeping is recovered. Based on Van Atta & Wyngaard's (1975) experimental results, Nelkin & Tabor concluded that the random sweeping hypothesis is valid at high Reynolds numbers.

Chen & Kraichnan (1989) clearly showed that the RNG approach discards sweeping effects at the outset and so does not prove the unimportance of sweeping. The main conclusion of Chen & Kraichnan is that a 'precise coherence between energy-range and inertial-range excitation is needed to inhibit sweeping effects'. This conclusion has important implications for the theory of fully developed turbulence.

To clarify this statement let us first cite an explanation from Chen & Kraichnan (1989). 'The Kolmogorov (1941) theory appeals to effective statistical independence of the one-time probability distributions of energy-range and inertial-range excitation: at any instant, the two ranges know about each other only through ϵ. Tennekes (1975) has pointed out that this implies a statistical form of Taylor's hypothesis so that inertial-range components of velocity field suffer advective sweeping by the energy-range excitation. Consequently, the many-time distribution of the inertial-range excitation involves the magnitude of the sweeping as well as ϵ.' A major contradiction can be seen from the citation. Indeed, Kolmogorov's theory is directly applicable only to the velocity differences (or to the appropriate energy spectra) because the characteristic time of these quantities is essentially smaller than that of a large-scale motion. Since Kolmogorov's hypotheses cannot be simply extended to more complex quantities (e.g. the higher-order spectra) which may be governed by other timescales, special generalization of the theory is necessary for a given complex case. As an example, Kraichnan (1965) generalized the Kolmogorov theory for hydromagnetic turbulence. The sweeping decorrelation hypothesis not only relies on the Kolmogorov (1941 *a*, *b*) theory but also implies some generalization. An additional large-scale governing parameter $\langle u_i^2 \rangle$ is added to the inertial-range description. If the hypothesis is valid, this modification is correct for the particular case of the higher-order spectra and, perhaps, may be further extended. On the other hand, the RNG result of Yakhot *et al.* (1989), (4), reduces the theory to the original form, that is the same as the straightforward extension of the Kolmogorov theory.

The above considerations are our motivation to revisit issues related to the random sweeping hypothesis. The highest possible Reynolds number laboratory measurements were used for analysis. We have investigated not only the consequences of the different approaches, (1)–(4), but also the basis of the hypothesis, i.e. the assumption of statistical independence of energy-range and inertial-range excitation. This general idea was converted to two different 'measurable' forms. The analysis was accomplished in terms of structure functions instead of the more usual spectral approach. As will be seen later, the structure function approach allowed us to examine this complicated question step by step.

2. The higher-order velocity structure functions

The term 'higher-order structure functions' is used hereafter in the same way that the term 'higher-order spectra' was used by Van Atta & Wyngaard (1975), i.e. it denotes the quantities:

$$D_{u_i}^{(m)}(r) = \langle [u_i^m(x_1+r) - u_i^m(x_1)]^2 \rangle, \quad i = 1,2,3, \quad m \geqslant 1, \tag{5}$$

where u_i denote velocity component fluctuations ($\langle u_i \rangle = 0$), x_i are coordinates, subscript 1 hereafter corresponds to the direction of a mean flow, and r is a distance, within the inertial subrange, in the x_1 direction.

Let us define $u_i(x_1) \equiv u$, $u_i(x_1+r) - u_i(x_1) \equiv \Delta u$ and $D_{u_i}^{(m)}(r) \equiv D_m$. Equation (5) can be represented in the form

$$\begin{aligned} D_m &= \langle \{[u+\Delta u]^m - u^m\}^2 \rangle \\ &= \langle [mu^{m-1}\Delta u + \tfrac{1}{2}m(m-1)\,u^{m-2}\Delta u^2 + \tfrac{1}{6}m(m-1)\,(m-2)\,u^{m-3}\Delta u^3 + \ldots]^2 \rangle. \end{aligned} \tag{6}$$

From (6) it follows that

$$\begin{aligned} D_m = m^2 \langle u^{2m-2}\Delta u^2 \rangle + m^2(m-1)\,\langle u^{2m-3}\Delta u^3 \rangle \\ + m^2(m-1)\,\tfrac{1}{12}(7m-11)\,\langle u^{2m-4}\Delta u^4 \rangle + \ldots. \end{aligned} \tag{7}$$

In particular,

$$D_2 = 4\langle u^2 \Delta u^2 \rangle + 4\langle u \Delta u^3 \rangle + \langle \Delta u^4 \rangle. \tag{8}$$

Equation (7) contains the terms $\langle u^k \Delta u^l \rangle$ with $k \geqslant 1$ and $l \geqslant 2$. As will be seen later, analysis of these terms gives a better understanding of the sweeping decorrelation hypothesis and allows us to find the direct one-time large- and small-scale eddy interaction (if any). Two methods will be used.

In the first method, the correlation coefficients between u and Δu are introduced as

$$\rho_{k,l}(r) = \frac{\langle (u^k - \langle u^k \rangle)\,(\Delta u^l - \langle \Delta u^l \rangle) \rangle}{\sigma_{(u^k)}\,\sigma_{(\Delta u^l)}}, \quad k \geqslant 1, \quad l \geqslant 2, \tag{9}$$

where $\sigma_{(\chi)}$ denotes the r.m.s. value of any quantity χ, i.e. $\sigma_{(\chi)} = [\langle (\chi - \langle \chi \rangle)^2 \rangle]^{\frac{1}{2}}$. From (9)

$$\langle u^k \Delta u^l \rangle = \langle u^k \rangle \langle \Delta u^l \rangle + \rho_{k,l}\,\sigma_{(u^k)}\,\sigma_{(\Delta u^l)}. \tag{10}$$

The sweeping hypothesis is based on the assumption of statistical independence of large- and small-scale motion. For this assumption to be valid, it is necessary that the large-scale characteristics u^k and the inertial-scale ones Δu^l be completely uncorrelated, i.e.

$$\rho_{k,l} = 0 \quad \text{for all} \quad k \geqslant 1, \quad l \geqslant 2. \tag{11}$$

From (10) and (11) follows:

$$\langle u^k \Delta u^l \rangle = \langle u^k \rangle \langle \Delta u^l \rangle \quad \text{for all} \quad k \geqslant 1, \quad l \geqslant 2. \tag{12}$$

It will be shown later that (12) leads directly to the sweeping decorrelation relations (1)–(3). The advantage of this approach is that the condition (11) may be easily checked experimentally because the value of $\rho_{k,l}$ can be measured from its definition, (9).

The second method is a direct check of large- and small-eddy interaction. In accordance with the definition

$$\begin{aligned}\langle u^k \Delta u^l \rangle &= \int u^k \Delta u^l P(u, \Delta u)\, \mathrm{d}u\, \mathrm{d}\Delta u \\ &= \int u^k \left[\int \Delta u^l P(\Delta u \,|\, u)\, \mathrm{d}\Delta u \right] P(u)\, \mathrm{d}u = \int u^k \langle \Delta u^l \rangle_u P(u)\, \mathrm{d}u,\end{aligned} \tag{13}$$

where $P(u, \Delta u)$ is a joint probability function of u and Δu; $P(\Delta u \,|\, u)$ is a conditional probability function of Δu at fixed value of u; and $\langle \Delta u^l \rangle_u$ are the l-order moments of velocity difference Δu conditionally averaged at fixed value of u. Here the relation $P(u, \Delta u) \equiv P(\Delta u \,|\, u) P(u)$ is used.

If the large, energy-containing fluctuations, characterized by u, and the small, inertial-range ones, characterized by Δu, are statistically independent, the $P(\Delta u \,|\, u) = P(\Delta u)$ and

$$\langle \Delta u^l \rangle_u = \langle \Delta u^l \rangle \quad \text{for all} \quad l \geqslant 2. \tag{14}$$

Equation (14) is one of the possible mathematical (and 'measurable'!) formulations of the assumption of statistical independence of large- and small-scale motion. It leads directly to (12).

Thus, (12) may be derived via two different methods. The first one is the assumption of complete statistical independence between energy-containing and inertial-range excitation. This assumption can be expressed in a 'measurable' form, (14). The second assumption is that excitations in the large and small scales are uncorrelated. The appropriate 'measurable' expressions are (9) and (11).

Let us adopt the assumption that: *velocity and velocity difference fluctuations are uncorrelated.* As a result, (11) and, consequently, (12) are valid. Using (12), the exact expansion (7) becomes

$$\begin{aligned}D_m = m^2 \langle u^{2m-2} \rangle \langle \Delta u^2 \rangle + m^2(m-1) \langle u^{2m-3} \rangle \langle \Delta u^3 \rangle \\ + m^2(m-1) \tfrac{1}{12}(7m-11) \langle u^{2m-4} \rangle \langle \Delta u^4 \rangle + \dots.\end{aligned}$$

It can be rewritten as

$$\begin{aligned}D_m = m^2 \langle u^{2m-2} \rangle \langle \Delta u^2 \rangle \Bigg[1 + (m-1) \frac{\langle u^{2m-3} \rangle \langle \Delta u^3 \rangle}{\langle u^{2m-2} \rangle \langle \Delta u^2 \rangle} \\ + (m-1) \frac{7m-11}{12} \frac{\langle u^{2m-4} \rangle \langle \Delta u^4 \rangle}{\langle u^{2m-2} \rangle \langle \Delta u^2 \rangle} + \dots \Bigg], \quad m \geqslant 2.\end{aligned} \tag{15}$$

To estimate the terms in the square brackets (15), the Kolmogorov (1941*a*, *b*) theory may be adopted. In this case

$$\langle \Delta u^l \rangle = C_l(\epsilon r)^{l/3}, \tag{16}$$

where C_l are the constants. In particular, $C_2 = C$ is the well-known Kolmogorov constant. Let us denote

$$K_k = \frac{\langle u^k \rangle}{\langle u^2 \rangle^{k/2}}, \quad \phi(r) = \frac{\langle \Delta u^2 \rangle}{\langle u^2 \rangle}. \tag{17}$$

It is obvious that $\phi \to 2$ if $r \to \infty$ and $\phi \to 0$ if $r \to 0$. Using (16) and (17), (15) may be written as

$$D_m = m^2 < u^{2m-2} \rangle \langle \Delta u^2 \rangle \left[1 + (m-1) \frac{K_{2m-3}}{K_{2m-2}} \frac{C_3}{C_2^{\frac{3}{2}}} \phi^{\frac{1}{2}} + (m-1) \frac{7m-11}{12} \frac{K_{2m-4}}{K_{2m-2}} \frac{C_4}{C_2^2} \phi + \dots \right], \quad m \geqslant 2. \tag{18}$$

To estimate $\phi(r)$, the local isotropy relation

$$\epsilon = 15\nu \left\langle \left(\frac{\partial u_1}{\partial x_1} \right)^2 \right\rangle \tag{19}$$

may be used (ν is kinematic viscosity). The Taylor microscale and appropriate Reynolds number are defined as

$$\lambda = \left(\frac{\langle u_1^2 \rangle}{\langle (\partial u_1 / \partial x_1)^2 \rangle} \right)^{\frac{1}{2}}, \quad R_\lambda = \frac{\langle u_1^2 \rangle^{\frac{1}{2}} \lambda}{\nu}, \tag{20}$$

and from (16), (17), (19) and (20) it follows that

$$\phi(r) = C_2 \left(\frac{15}{R_\lambda} \frac{r}{\lambda} \right)^{\frac{2}{3}}. \tag{21}$$

It is well known from numerous measurements that in high Reynolds number free shear flows the large-scale characteristics are of the order of $|K_3| < 0.5$, $|K_5| < 3$, $K_4 \geqslant 3$, $K_6 \geqslant 15$. The inertial-range constants are approximately $C_2 \approx 2$, $C_3 = -0.8$, $C_4 \approx 25$ so the factor $C_2(15)^{\frac{2}{3}}$ in (21) is about 12. The ratio r/λ usually does not exceed 10 in the inertial subrange. It is seen from (21) that the value of ϕ is small for sufficiently large values of R_λ and small r. So all terms with $\phi(r)$ in the square brackets in (18) can be omitted if fine-scale turbulence (with small r) at sufficiently high Reynolds number R_λ is analysed. In this case (18) reduces to

$$D_{u_i}^{(m)}(r) = m^2 \langle u_i^{2m-2} \rangle \langle [u_i(x+r) - u_i(x)]^2 \rangle, \quad i = 1,2,3, \quad m \geqslant 2. \tag{22}$$

From (22) and (16), one finds

$$D_{u_i}^{(m)}(r) \sim r^{\frac{2}{3}}, \quad i = 1,2,3, \quad m \geqslant 2. \tag{23}$$

It is necessary to emphasize that the above estimate is valid for (15) but may be inapplicable to the exact equation (7) if the assumption (11) is violated.

Equations (22) and (23) express the sweeping decorrelation hypothesis in terms of high-order structure functions. Indeed, the Fourier transformation of (23) coincides with (3). Moreover, if a Gaussian distribution for u_i is assumed, the Fourier transformation of (22) reproduces (2), proposed by Van Atta & Wyngaard (1975). The structure function approach has two advantages over the spectral one. First, the derivation of the final equation (22) is thoroughly clear and all assumptions, i.e. (11), the hypothesis of statistical independence of large- and small-scale eddies, (14), etc., are completely transparent. Second, all of these assumptions may be directly measured and checked. These detailed measurements were the purpose of the experimental investigation presented below.

3. Apparatus and measurement techniques

Two velocity time series in high Reynolds number laboratory shear flows were analysed. The first one was obtained in the large wind tunnel of the Central Aerohydrodynamic Institute (Moscow). The mixing layer between a jet issuing from an elliptical nozzle (14×24 m^2) and ambient air was studied. The wind tunnel had an open 24 m long working section. Measurements were performed on the line, which continued the nozzle wall, at a distance $x_1 = 20$ m downstream of the nozzle. The free jet velocity was equal to $U_0 = 11.8$ m/s.

The second time series was obtained in the return channel of the same wind tunnel. The channel was 175 m long and 22 m wide. Its height rose linearly from 20 m to 32 m. Measurements were done in the plane of symmetry from a tower 5 m above floor level. Flow velocity at the measuring station was 10.8 m/s.

Standard thermoanemometers were used. An X-wire probe with perpendicular wires was operated at an overheat ratio of 0.8. Wires were made of platinum-plated tungsten. The length of each wire was 0.5 mm, its diameter was 2.5 μm. The distance between wires was 0.5 mm.

Signals from both wires were filtered to reduce noise level, digitized and processed on a computer. The low-pass filter cutoff frequency was $f_c = 1.7$ kHz. The sampling frequency f_s and one-channel time series length N were 8 kHz and 2000000 samples respectively. These values were chosen to investigate the inertial subrange. To measure the energy dissipation rate ϵ, f_c, f_s and N were doubled.

Descriptions of the experiments and analysis of measurement errors (temporal and spatial resolution, statistical convergence, nonlinearity of the hot-wire response etc.) can be found in Karyakin, Kuznetsov & Praskovsky (1991) (and a detailed version in English: Kuznetsov, Praskovsky & Karyakin 1992).

The main flow characteristics at the measurement points are listed in table 1. Two velocity fluctuation components were analysed at every point: in the direction of the mean flow (subscript 1) and in the direction of the mean shear (subscript 2). The mean energy dissipation rate ϵ was estimated using the local isotropy relation (19). The local isotropy hypothesis was used only for presentation of the results and did not influence the conclusions. Taylor's hypothesis was used to convert the temporal into the spatial coordinate. The influence of Taylor's hypothesis will be discussed in the next section. The integral lengthscale L and the Kolmogorov scale η were estimated with standard formulae:

$$L = \frac{U_1}{\langle u_1^2\rangle}\int_0^\infty \langle u_1(t+\tau)\,u_1(t)\rangle\,\mathrm{d}\tau, \quad \eta = (\nu^3/\epsilon)^{\frac{1}{4}}.$$

where t is a time and τ is a time delay.

It is seen from table 1 that at Reynolds numbers R_λ up to 3.2×10^3, which were achieved in the laboratory measurements, the energy–viscous eddy separation, i.e.

Flow	Return channel	Mixing layer
U_1 (m/s)	10.8	7.87
$\sigma_{(u_1)}$ (m/s)	1.03	1.67
$\sigma_{(u_1)}/U_1 \times 100\%$	9.54	21.2
$\sigma_{(u_1^2)}/\langle u_1^2 \rangle$	1.52	1.35
$\sigma_{(u_1^4)}/\langle u_1^4 \rangle$	3.61	3.08
$\sigma_{(u_1^6)}/\langle u_1^6 \rangle$	7.59	6.74
$\sigma_{(u_2)}$ (m/s)	0.828	1.27
$\sigma_{(u_2^2)}/\langle u_2^2 \rangle$	1.58	1.39
$\sigma_{(u_2^4)}/\langle u_2^4 \rangle$	3.98	3.36
$\sigma_{(u_2^6)}/\langle u_2^6 \rangle$	8.21	7.42
ϵ (m^2/s^3)	0.11	1.9
λ (mm)	46	18
$R_\lambda \times 10^{-3}$	3.2	2.0
L (m)	4.8	1.3
η (mm)	0.41	0.21
$L/\eta \times 10^{-3}$	12	6.2

TABLE 1. Main turbulence characteristics at analysed points.

the ratio L/η, exceeded 6000 which was big enough for the existence of a substantial inertial subrange.

4. Results

4.1. *Measurements of the higher-order structure functions*

The measured higher-order structure functions are presented in Figure 1 (in all plots solid and open symbols correspond to the u_1 and u_2 velocity components respectively). The quantities $D_{u_i}^{(m)}(r)$ are normalized with $r^{\frac{2}{3}}$ in accordance with (23). First, it is necessary to determine the inertial subrange bounds. The well-known determination $\eta \ll r \ll L$ is not really precise so the choice is always somewhat arbitrary. It is seen in figure 1 that the interval $20\eta \leqslant r \leqslant \frac{1}{4}L$ may be treated as the inertial one in both flows. To exclude any doubts, it is assumed that the distance r belongs to the inertial subrange if $30\eta \leqslant r \leqslant \frac{1}{8}L$. Vertical arrows in all plots correspond to the latter bounds. It is seen in figure 1 that the choice is quite reasonable. At least, the 'two-thirds law' is valid within these bounds with high accuracy.

It is worthwhile to explain here the chosen form of data presentation. In fact, all results from the two flows studied are qualitatively similar. So either series of plots, (*a*) or (*b*), may be sufficient to illustrate the results obtained. However, we believe that it is beneficial to present the data for both flows: these data are unique owing to their high (laboratory) Reynolds numbers as well as the high data reliability; also the experimental results may be useful to other workers in studies of other turbulence problems.

It is seen from figure 1 that measured values of $D_{u_i}^{(m)}(r)$ are in agreement with (23) for both velocity components. This agrees with previous measurements of higher-order spectra by Dutton & Deaven (1972) and Van Atta & Wyngaard (1975). However, it is somewhat premature to draw any conclusion at this stage. Yakhot *et al.* (1989) pointed out that the use of Taylor's hypothesis may contaminate experimental results. Let us try to analyse this important objection.

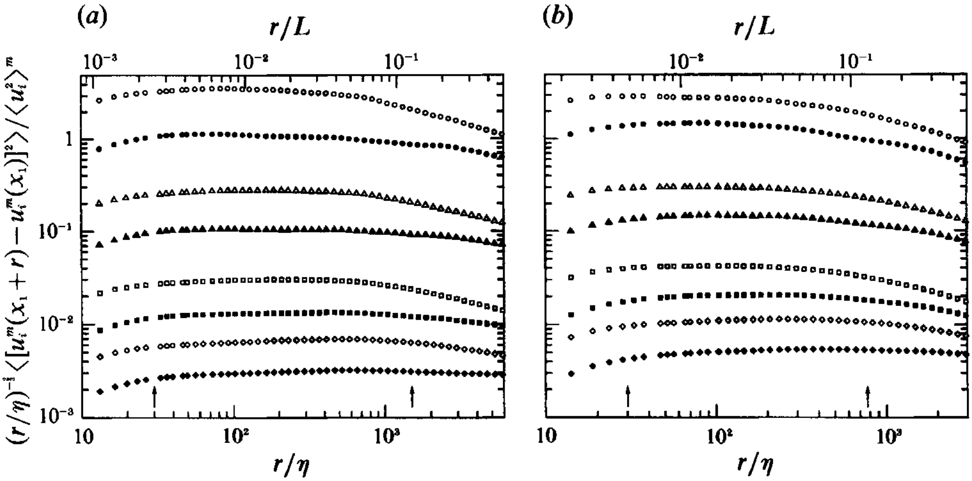

FIGURE 1. Higher-order structure functions $D_{u_i}^{(m)}(r)$: (a) return channel; (b) mixing layer. ◇, $m = 1$; □, 2; △, 3; ○, 4. In this and subsequent figures: solid symbols, u_1; open symbols, u_2; vertical arrows correspond to the inertial-range bounds.

At the moment a direct test of this question is impossible because only one-point measurements of small-scale fluctuations at high Reynolds numbers are available, so we devised an indirect approach. The sweeping hypothesis was expressed in two different measurable forms which will be assumed to be affected by the use of Taylor's hypothesis (if it all) in different ways. Any difference in the results is expected to be a test of the effect.

Equation (22) can be rewritten in the form

$$d_{u_i}^{(m)}(r) = \frac{D_{u_i}^{(m)}(r)}{m^2\langle u_i^{2m-2}\rangle\langle[u_i(x+r)-u_i(x)]^2\rangle} = 1, \quad i = 1,2,3, \quad m \geqslant 2. \tag{24}$$

These non-dimensional ratios include similarly constructed velocity differences both in the numerator and denominator, so the influence of Taylor's hypothesis on the values of $d_{u_i}^{(m)}(r)$ must be weaker than on $D_{u_i}^{(m)}(r)$. However, this influence (if any) may be different at different m because of independent averaging in the numerator and denominator.

The second test was based on the quantities

$$G_{u_i}^{(m)}(r) = \left\langle \frac{[u_i^m(x+r)-u_i^m(x)]^2}{[u_i(x+r)-u_i(x)]^2} \right\rangle, \quad m \geqslant 2. \tag{25}$$

These quantities include velocity differences both in the numerator and denominator at the same locations (or instances). Thus the influence of Taylor hypothesis on the values of $g_{u_i}^{(m)}(r)$ may be expected to be either negligible or at least quite different from that on the $d_{u_i}^{(m)}(r)$. With a derivation completely similar to that of (24), one can easily obtain

$$g_{u_i}^{(m)}(r) = \frac{G_{u_i}^{(m)}(r)}{m^2\langle u_i^{2m-2}\rangle} = 1, \quad i = 1,2,3, \quad m \geqslant 2. \tag{26}$$

Equation (26) cannot be directly connected with the spectral expressions (1)–(3). Nevertheless, it is another form of the sweeping hypothesis because it is based on the same assumptions and derived in the same way. Experimental data are presented in

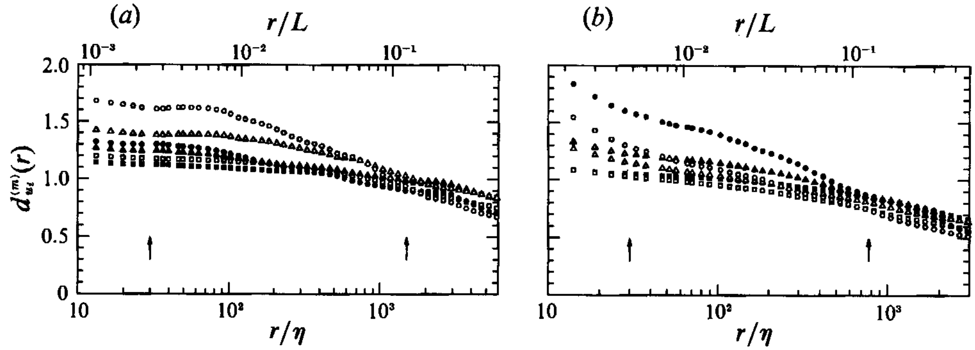

FIGURE 2. Non-dimensional higher-order structure functions $d_{u_i}^{(m)}(r)$. For symbols see figure 1.

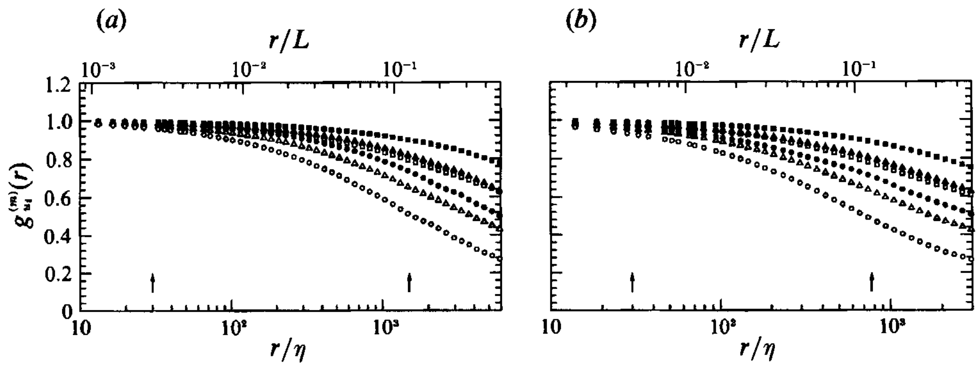

FIGURE 3. Non-dimensional higher-order quantities $g_{u_i}^{(m)}(r)$. For symbols see figure 1.

figures 2 and 3. Linear scales for the ordinate were chosen to underline departures from unity in (24) and (26). In both cases the distance r was estimated as $r = \tau U_1$. It is seen that the variations of $d_{u_i}^{(m)}(r)$ and $g_{u_i}^{(m)}(r)$ over the inertial subrange are approximately of the same magnitude. It means that the error due to using Taylor's hypothesis is negligibly small in this case. For further analysis of such an influence, an additional experiment was carried out. The two-point distances r in (25) were estimated as (see Heskestad 1965)

$$r = \tau \overline{U_1(\tau)}, \tag{27}$$

where τ_1 was a variable parameter and the local large-scale velocity $\overline{U_1(\tau_l)}$ at the instant t_j was calculated as

$$\overline{U_1(\tau_l)} = U_1 + \frac{1}{\tau_l}\int_{-\tau_l/2}^{\tau_l/2} u_1(t_j+\zeta)\,\mathrm{d}\zeta. \tag{28}$$

Results presented in figure 3 correspond to $\tau_l = T$, where T is the whole duration of the time series. The value of τ_l was changed from $0.3L/U_1$ up to $3L/U_1$ and no dependence of $g_{u_i}^{(m)}(r)$ on τ_l was found. The results for $m = 4$, where the dependence on τ_l was found to be the most significant one, are presented in figure 4. Thus it can be concluded from figures 2–4 that the influence of Taylor's hypothesis on the investigated parameters can be neglected both in the mixing layer, where the turbulence intensity was over 21 %, and in the return channel where it was less than 10 %.

It is seen in figures 2 and 3 that within the inertial subrange for both velocity components the departures of $d_{u_i}^{(m)}(r)$ and $g_{u_i}^{(m)}(r)$ from unity are sufficiently small for

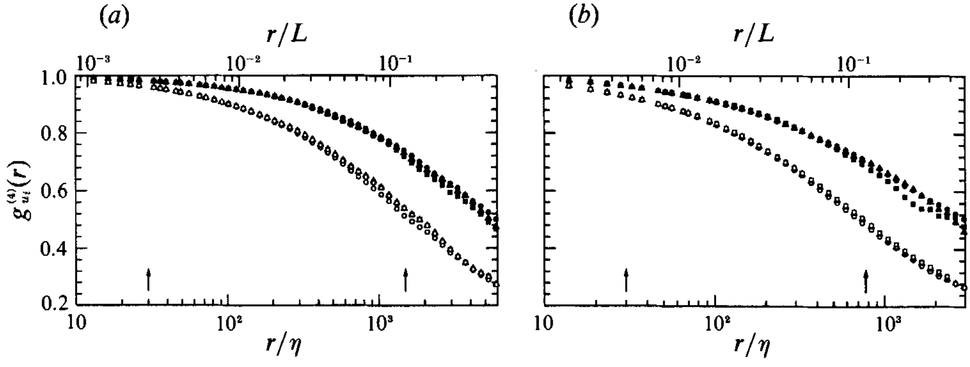

FIGURE 4. The influence of the averaging interval τ_l on the values of $g_{u_i}^{(4)}(r)$: (*a*) return channel; (*b*) mixing layer. ○, $\tau_l = T$; △, L/U_1; △, $0.3L/U_1$.

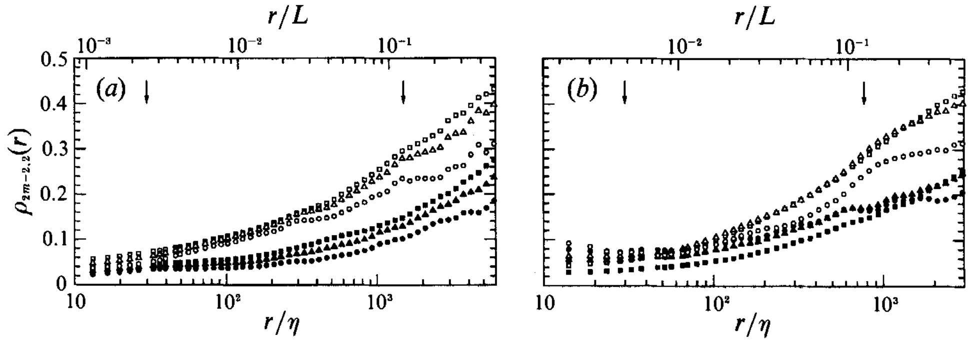

FIGURE 5. Correlation coefficients $\rho_{2m-2,2}(r)$. For symbols see figure 1.

$m = 2$. These discrepancies increase with m and the values of $d_{u_i}^{(4)}(r)$ and $g_{u_i}^{(4)}(r)$ vary by a factor of about two over the inertial subrange. Nevertheless, it appears that the data agree with (23), at least qualitatively, and confirm the sweeping decorrelation hypothesis because the variation of the parameters $g_{u_i}^{(m)}(r)$ and $d_{u_i}^{(m)}(r)$ could be considered as small (see figure 1). But an alarming observation can be seen in figures 2 and 3. Indeed, it can be easily found from (15) and (24) that if the assumption (11) is even approximately valid, the quantity $d_{u_i}^{(2)}(r)$ cannot decrease as the distance r increases. This result is valid for $g_{u_i}^{(2)}(r)$ even more rigorously. The tendency obtained in figures 2 and 3 may be attributed either to measurement errors or to some more deep physical reasons. As it will be shown in the next subsection, the latter alternative is the correct one. The fact is that the basic assumption (11) is not valid in the flows studied.

4.2. *Correlation of energy- and inertial-range excitation*

The first right-hand-side terms in (15) are connected with correlation coefficients $\rho_{2m-2,2}(r)$ in (9). The measured values of these coefficients are presented in figure 5. Two features are seen there. First, $\rho_{2m-2,2}(r)$ decreases with decrease of r, as expected. Second, these values are surprisingly large throughout the whole inertial subrange. Indeed, $\rho_{2m-2,2} \approx 0.04-0.08$ for $r/\eta = 30$. To illustrate the significance of the correlations obtained, let us return to (10). It can be rewritten as

$$\langle u^{2m-2}\Delta u^2\rangle = \langle u^{2m-2}\rangle\langle \Delta u^2\rangle\left[1+\rho_{2m-2,2}\frac{\sigma_{(u^{2m-2})}}{\langle u^{2m-2}\rangle}\frac{\sigma_{(\Delta u^2)}}{\langle \Delta u^2\rangle}\right], \quad m \geqslant 2. \tag{29}$$

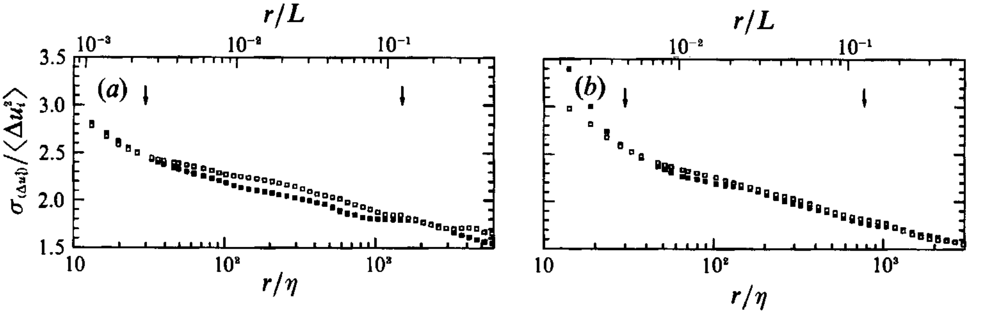

FIGURE 6. Ratio of r.m.s.-to-average values for Δu_i^2: (*a*) return channel; (*b*) mixing layer.

So the real influence of correlation coefficients depends on the ratios of r.m.s. values to average ones for moments both of velocities and of the inertial-subrange velocity differences. The latter values characterize the variability of energy at different r through the cascade and are interesting in themselves. The measured values of $\sigma_{(\Delta u^2)}/\langle \Delta u^2 \rangle$ are presented in figure 6. The ratios $\sigma_{(u^{2m-2})}/\langle u^{2m-2}\rangle$ are listed in table 1. It is seen that multipliers for $\rho_{2m-2,2}$ in (29) are of the order of 10 for $r/\eta = 30$ so the second term on the right-hand side is comparable with unity and therefore a significant one. The results clearly show that the sweeping decorrelation hypothesis cannot be exactly valid even at very high (up to $R_\lambda = 3.2\times 10^3$) but finite Reynolds numbers.

On the other hand, two contradictions seem to be present in our results reported above. The first one is that the measured higher-order structure functions are in agreement with the hypothesis (figures 1–3). The second one was mentioned above and concerns the behaviour of the quantity $d^{(2)}_{u_i}(r)$. Because coefficient $\rho_{2,2}$ is significant and positive, $d^{(2)}_{u_i}(r)$ must be an increasing function of r. To resolve these seeming contradictions, it is necessary to realize that derivation of the final equation (22) was based on two assumptions. The first one was formalized by (11) as a decorrelation of energy-range and inertial-range excitation. But this assumption leads only to (15). The second assumption was a consequence of the first but may be treated as an additional independent assumption. It was used when the right-hand-side terms in (15) were neglected. As was emphasized in §2, the estimation may be invalid for exact (7) if the assumption (11) is violated.

To demonstrate this point, we write the exact equation (8) as

$$D_2(r) = 4\langle u^2\Delta u^2\rangle\,[1+\delta_1(r)+\delta_2(r)], \tag{30}$$

$$\delta_1 = \langle u\Delta u^3\rangle/\langle u^2\Delta u^2\rangle, \quad \delta_2 = \langle \Delta u^4\rangle/4\langle u^2\Delta u^2\rangle.$$

As was shown in §2, $\delta_1 = 0$ and $\delta_2 \ll 1$ if (11) is valid. The measured values of δ_1 and δ_2 are presented in figure 7. It is seen that the negative quantity δ_1 becomes dominant with increasing r and overcomes the increase of δ_2. This is additional evidence of interaction between large- and small-scale motions. Thus there is no intrinsic contradiction in the results presented, i.e. they are self-consistent.

We have encountered a surprising situation here. The experimental results, figures 1–3, are in agreement with theoretical predictions, (1)–(3) and (22)–(26), which follow from the sweeping decorrelation hypothesis. Such agreement is usually treated as a confirmation of a hypothesis. But a direct check of the assumption (11), which is the basis of the hypothesis, has shown that this agreement is a mere coincidence. The

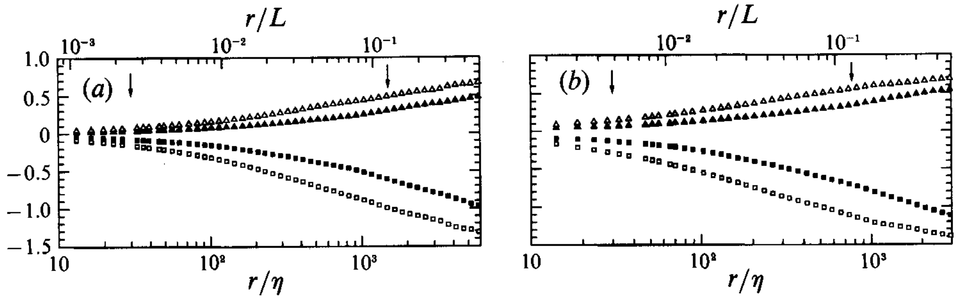

FIGURE 7. Non-dimensional right-hand-side terms in (30): (*a*) return channel; (*b*), mixing layer. □, δ_1; △, δ_2.

violation of this assumption leads to some self-compensating consequences. Let us illustrate this with (8) for $m = 2$: the quality $D_2 = 4\langle u^2 \Delta u^2\rangle$ becomes invalid, but another consequence of assumption (11): $\delta_1 = 0$ and $\delta_2 \ll 1$, becomes invalid too. Because they act in opposite directions, the final result looks like a confirmation of the hypothesis.

4.3. *Direct large- and small-scale eddy interaction*

As was pointed in §2, the non-zero correlations (9) directly lead to the violation of the hypothesis (14) on the statistical independence of large- and fine-scale fluctuations. This hypothesis is the basis of the recent theory of locally isotropic turbulence at high Reynolds numbers. Detailed analysis of the hypothesis can be found in the book by Batchelor (1953). The Kolmogorov (1941*a*, *b*, 1962) theory relies on this hypothesis too. The validity of (14) in the limit of $r \rightarrow 0$ was confirmed by Kuznetsov, Praskovsky & Sabelnikov (1992) where the equality $\langle(\partial u_1/\partial x_1)^{2l}\rangle_u = \langle(\partial u_1/\partial x_1)^{2l}\rangle$ was experimentally verified within turbulent fluid for $l = 1, 2, 3$. But the results presented above show that the hypothesis seems to be invalid in the inertial subrange. This question should be re-examined.

To calculate the moments of velocity differences, conditionally averaged over the fixed value of velocity u_i^*, the following procedure was used. The values $[u_i(t+\frac{1}{2}\tau) - u_i(t-\frac{1}{2}\tau)]^l$ were averaged over the instants t_j when $u_i^* - \frac{1}{2}\delta u_i \leqslant u_i(t_j) \leqslant u_i^* + \frac{1}{2}\delta u_i$. The velocity interval was equal to $\delta u_i \approx 0.1\sigma_{u_i}$. Then Taylor's hypothesis $r = \tau U_1$ was used to estimate the separation r.

The ratios $\langle \Delta u_i^l\rangle_{u_i^*}/\langle \Delta u_i^l\rangle$ for $l = 2$–4 were measured. In accordance with (14), these values have to be equal to unity for all levels u_i^* and orders l if large- and small-scale motions are statistically independent. For $l = 3$ the modulus of the velocity difference was calculated. It is obvious that (14) is valid both for Δu and for $|\Delta u|$ if the hypothesis analysed is valid. There were two reasons to use absolute values for third-order moment calculations. The first, well-known, one was that the statistical convergence of $\langle|\Delta u|^3\rangle$ is essentially better than of $\langle \Delta u^3\rangle$. The second reason was the following physical considerations. The quantity $\langle \Delta u^3\rangle$ is connected with the mean energy dissipation rate (Kolmogorov 1941*b*) by

$$\langle \Delta u_1^3\rangle = -\tfrac{4}{5}\epsilon r. \tag{31}$$

The quantity $|\Delta u|^3/r$ determines the instantaneous energy flux through eddies with characteristic scale r within the inertial subrange (Kolmogorov 1962, see also Kraichnan 1974). As will be seen in §5, it is a more important quantity for our purpose.

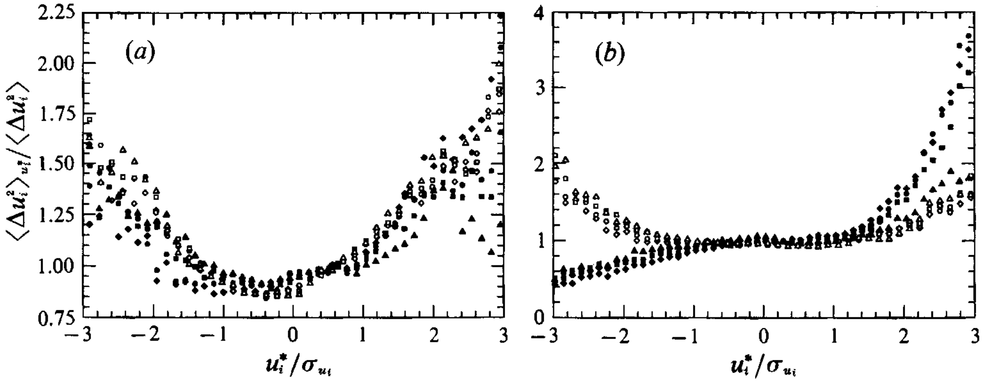

FIGURE 8. Second-order moments of velocity differences, conditionally averaged at a fixed value of the velocity. (*a*) Return channel: ◇, $r/\eta = 33$; ○, 115; □, 390; △, 1300; (*b*) mixing layer: ◇, $r/\eta = 37$; ○, 110; □, 320; △, 940.

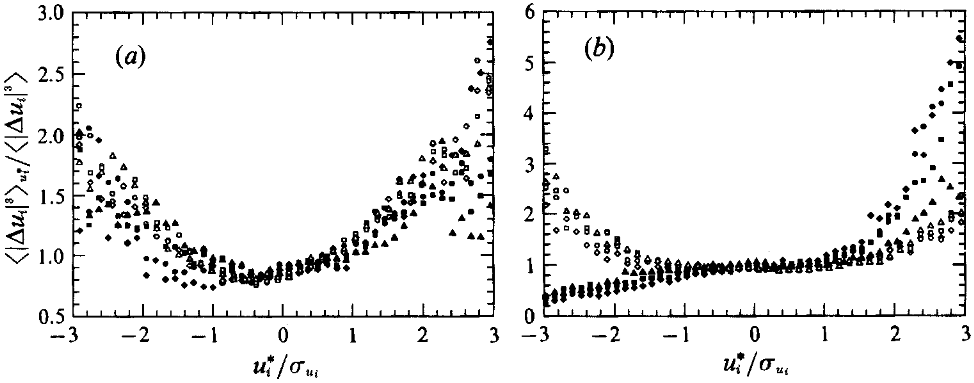

FIGURE 9. Third-order moments of velocity difference modulus, conditionally averaged at a fixed value of the velocity. For symbols see figure 8.

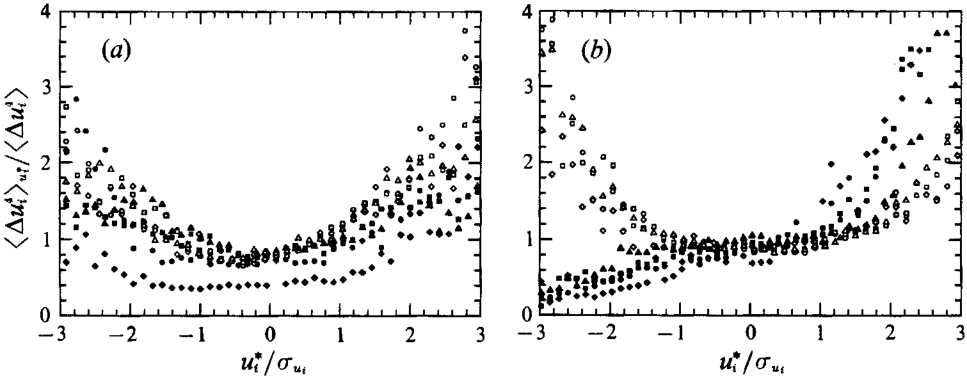

FIGURE 10. Fourth-order moments of velocity differences, conditionally averaged at a fixed value of the velocity. For symbols see figure 8.

The results obtained are presented in figures 8–10. Data scatter inevitably increases for the rare events at large u_i^*, so results are given only for $|u_i^*|/\sigma_{u_i} \leqslant 3$. Within these bounds, the number of independent samples for conditional averaging was not less than 1500 for any value of u_i^*. Nevertheless the data scatter is large enough especially for $l = 4$. In spite of the scatter, two main characteristic features are seen in these plots. The first one is strong dependence of the inertial-subrange fluctuations on large-scale parameters, in agreement with the measurements of correlations $\rho_{2m-2,2}$ presented above. Moreover, the magnitude of the interaction depends weakly on r within the inertial subrange. The second feature is that this dependence is just qualitatively different in different flows and for different velocity components (compare *a* and *b* in figures 8–10).

To check the reliability of the results, two additional experiments were performed. The first one† was directed to estimate the influence of Taylor's hypothesis. The numerical algorithm described above was modified in the following way (it was done for the u_1 component only). After the instant for conditional averaging t_j was fixed, i.e. the condition $u_i^* - \frac{1}{2}\delta u_1 \leqslant u_1(t_j) \leqslant u_i^* + \frac{1}{2}\delta u_1$ was recognized, the local sweeping velocity $\overline{U_1(\tau_l)}$ was calculated from (28). Then the two-point separation was estimated from (27). The value of τ_l was changed from $0.3L/U_1$ up to $3L/U_1$ and no dependence of $\langle \Delta u_1^2 \rangle_{u_i^*}$ on τ_l was found.

The second experiment was directed to clarify whether the instantaneous value of u_1 is really a representative characteristic of large-scale motion. To check this question, another kind of conditional averaging was done. The quantities $\langle \Delta u_1^2 \rangle_{\overline{u_1(\tau_l)}}$ were calculated, where $\overline{u_1(\tau_l)}$ was equal to the second term in the right-hand side of (28). The averaging interval was changed in the same range $0.3L/U_1 \leqslant \tau_l \leqslant 3L/U_1$ and no dependence of $\langle \Delta u_1^2 \rangle_{\overline{u_1(\tau_l)}}$ on τ_1 was found.

So it seems that the results presented in figures 5 and 8–10 reliably prove the existence of significant direct large- and inertial-scale eddy interaction. This interaction was established in two different types of high Reynolds number laboratory shear flows with $R_1 \approx 2.0 \times 10^3$ and 3.2×10^3. The results show that some revision of the recent models of small-scale turbulence at high Reynolds numbers is desirable.

5. Discussion

Almost all recent models of fine-scale turbulence structure in high Reynolds number flows are based on the hypotheses proposed by Kolmogorov (1941 *a*, *b*, 1962) and Oboukhov (1941, 1962). These hypotheses about locally homogeneous and isotropic turbulence significantly simplify the description of high Reynolds number turbulence and so they are very attractive and widely accepted. However, sometimes the Kolmogorov–Oboukhov results and even consequences from their hypothesis obtained by others are extended to cases for which they are inapplicable. The issue related to the sweeping decorrelation hypothesis is an example. Nelkin & Tabor (1990) wrote (p. 81): 'With this hypothesis of random sweeping, the power laws for the spectrum are still universal, but the 1941 Kolmogorov assumption breaks down since the velocity u, which is a property of the nonuniversal large scales, enters the result'. Indeed, it breaks down because Kolmogorov's hypotheses were extended to a case where they are inapplicable. Moreover, the result about direct interaction of large- and small-scale eddies, presented in the previous subsection, which appears as direct evidence against the hypotheses, does not really contradict them.

To clarify this matter, let us extract two main features of both (1941 and 1962) hypotheses. The first one is that they are based on a cascade mechanism of energy transfer and so they are valid only asymptotically, i.e. for $R_\lambda \rightarrow \infty$. As was pointed

† Proposed by Dr R. Rogallo.

out by Karyakin *et al.* (1991), the number of steps N in such a cascade cannot be very big even in the highest Reynolds number flows physically possible since $N \sim \log_2 L/\eta$. So at finite Reynolds numbers the hypotheses may be treated only as a first approximation and some deviations are possible. The second and most important feature is that the hypotheses are valid only for the 'locally averagable' values. This idea is clearly explained by Kolmogorov (1962). The small-scale turbulence knows about the energy-containing range only through the energy flux $\varepsilon(\boldsymbol{x}, t) \sim |\Delta u|^3/r$, but not through the mean rate of energy dissipation ϵ (see remarks by Kraichnan 1974). The quantity $\varepsilon(\boldsymbol{x}, t)$ is determined by the instantaneous energy flux from large-scale eddies $\varepsilon_{\mathrm{L}}(\boldsymbol{x}, t)$, which may be roughly estimated as

$$\varepsilon_{\mathrm{L}}(\boldsymbol{x}, t) \sim \frac{|\boldsymbol{u}(\boldsymbol{x}, t)|^3}{L(\boldsymbol{x}, t)}. \tag{32}$$

This relation has been reliably proved in numerous semi-empirical models. For very high Reynolds numbers and in a state of dynamical equilibrium between large- and small-scale motions, the average value of energy flux ε was assumed by Kolmogorov to be equal to the average value of energy dissipation. But averaging here is accomplished over regions with sizes of order of (or smaller than) the integral scale, and over times of order of (or smaller than) the interval of time correlation. Only in this case may the large-scale influence on fine-scale structure be assumed to be quasi-homogeneous and quasi-stationary so that the dynamical equilibrium can be achieved. So both the hypotheses and their consequences may be applied only to quantities for which such an averaging may be done. Velocity differences Δu_i^l are an example of such quantities and all Kolmogorov's results were obtained for Δu_i^l only. The 'two-thirds law' and (31) have been well verified by numerous experiments.

Let us come back to the results presented in figures 8–10. They become natural and expected if Kolmogorov's considerations above are really taken seriously. Indeed, the instantaneous energy flux from large-scale eddies to small-scale ones depends on $|\boldsymbol{u}|^3/L$, equation (32). If local dynamical equilibrium between the flux from energy-containing eddies and viscous dissipation in the smallest eddies is achieved, the instantaneous flux through the inertial subrange eddies with scale r, proportional to $|\Delta \boldsymbol{u}|^3/r$, must depend on $\varepsilon_{\mathrm{L}}(\boldsymbol{x}, t)$. To check this dependence quantitatively, $|\boldsymbol{u}|^3$ and $|\Delta u|^3$ have to be measured, i.e. simultaneous records of all three velocity components are necessary. In our experiments only two components were recorded simultaneously so surrogate quantities, (14), were measured and only qualitative results were obtained. It is obvious that the large-scale quantity u_i is connected with ε_{L}, and the inertial-quantity $|\Delta u_i|$ is connected with ε. The dependence of $\langle \Delta u_i^l \rangle_{u_i}$ on u_i obtained reflects the variability of the energy flux through the cascade in accordance with the variability of this flux due to inhomogeneity of large-scale motion. It is natural that this dependence is different in different flows and for different velocity components owing to different relations between $|\boldsymbol{u}|$ and u_i and $|\Delta \boldsymbol{u}|$ and Δu_i. Thus the type of direct interaction between large- and small-scale eddies we have found does not contradict the Kolmogorov's basic hypotheses.

Moreover, the results allow some further generalization of the classical hypotheses. To clarify it, the non-dimensional quantities

$$\alpha_l(u_i^*) = \frac{\langle \Delta u_i^l \rangle_{u_i^*}}{\langle |\Delta u_i|^3 \rangle_{u_i^*}^{l/3}}, \quad l = 2, 4$$

are presented in figures 11 and 12 respectively. It is obvious that for $\eta \ll r \ll L$ these

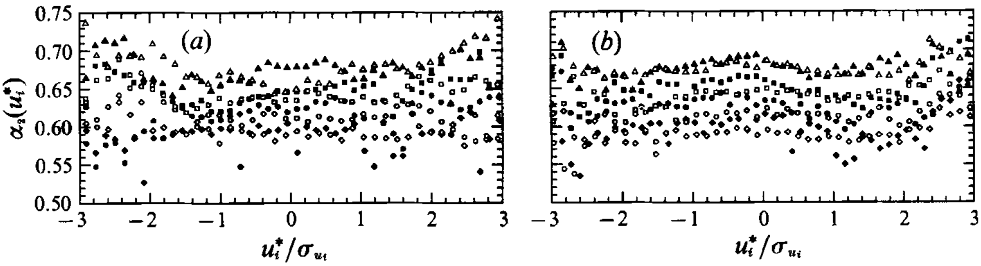

FIGURE 11. Non-dimensional second-order moments of velocity differences, conditionally averaged at a fixed value of the velocity. For symbols see figure 8.

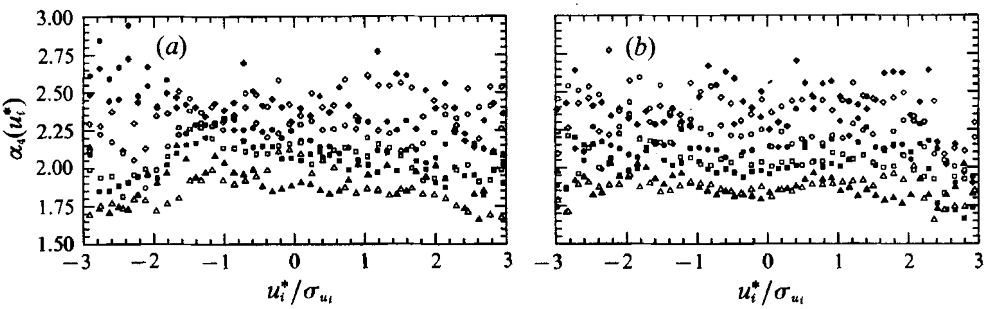

FIGURE 12. Non-dimensional fourth-order moments of velocity differences, conditionally averaged at a fixed value of the velocity. For symbols see figure 8.

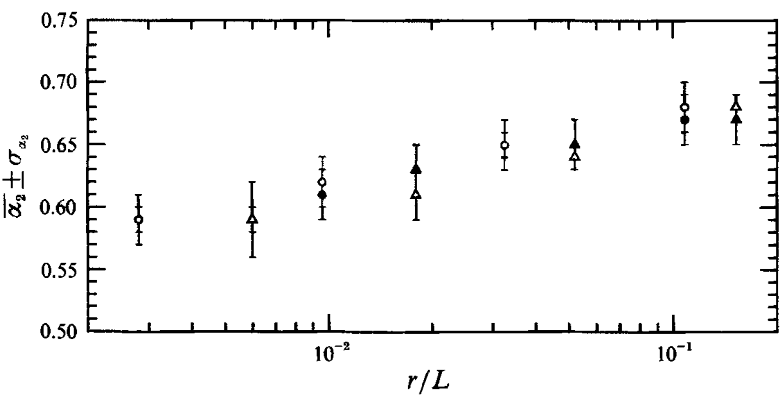

FIGURE 13. Non-dimensional second-order moments $\alpha_2(u_i^*)$, averaged over the range $-3 \leqslant u_i^*/\sigma_{u_i} \leqslant 3$: ○, return channel; △, mixing layer.

quantities characterize the intrinsic behaviour of the inertial subrange fluctuations at a fixed value of the energy flux from large-scale eddies. In spite of significant data scatter, it seems that the intrinsic structure of the inertial range does not depend on the value of u_i^*. To estimate the possible dependence of α_l on u_i^*, r, type of flow etc., let us introduce the quantities $\overline{\alpha_l}$ and $\sigma_{\alpha_l} = \overline{[(\alpha_l - \overline{\alpha}_l)^2]}^{\frac{1}{2}}$ averaging over the range $-3 \leqslant u_i^*/\sigma_{u_i} \leqslant 3$ (denoted by the overbar). As before this limited range is used simply to reduce scatter due to inadequate sampling times in our experiment, and is not a physically significant limitation. The value of σ_{α_l} may be treated here as an error of the assumption $\alpha_l(u_i^*) = \text{const}$. Measured values of $\overline{\alpha_l}$ for $l = 2, 4$ are presented in figures 13 and 14. The error bars in these figures correspond to the values of $\pm\sigma_{\alpha_l}$. It is seen that to a first approximation the inertial-subrange intrinsic

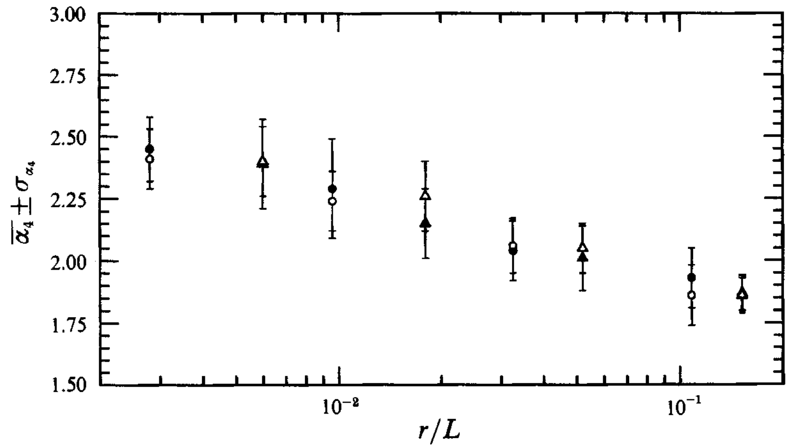

FIGURE 14. Non-dimensional fourth-order moments $\alpha_4(u_i^*)$, averaged over the range $-3 \leqslant u_i^*/\sigma_{u_i} \leqslant 3$. For symbols see figure 13.

structure depends weakly only on separation r and not on flow type, direction of u_i, or Reynolds number.

Returning to the sweeping decorrelation hypothesis. The first right-hand-side term in the exact equation (7), $\langle u^{2m-2}\Delta u^2\rangle$, includes the direct interaction of large- and small-scale quantities. In contrast to $\langle \Delta u^m \rangle$, it is not 'averagable' in the sense used above, i.e. over the regions smaller than the integral scale, and so it may depend on u^{2m-2} even in the inertial subrange. It appears that there is no real contradiction with Kolmogorov hypotheses if the higher-order spectra or structure functions in the inertial subrange depend on the large-scale parameters. The dependence reflects the spatial and temporal variability of the energy flux through the cascade, in full agreement with Kolmogorov's ideas. The only reason why this hypothesis cannot be exactly valid is the strong correlation described above between energy-containing and inertial-range excitation (see the conclusion by Chen & Kraichnan 1989). But even in this case the main result, (3) or (23), which follows from the hypothesis, is in acceptable agreement with experiments.

6. Summary

The validity of Tennekes's (1975) random sweeping hypothesis is one of the unresolved questions of high Reynolds number turbulence theory. The hypothesis is based on the assumption of complete decorrelation of energy- and inertial-range excitation. The latter assumption is itself extremely important because it is (either directly or indirectly) used in almost all recent models of fine-scale turbulence structure. The main consequence of the sweeping hypothesis, that the higher-order spectra $E^{(m)}$ of velocity powers $u_i^m, m \geqslant 2$, scale in the inertial subrange as $k^{-\frac{5}{3}}$, was strongly confirmed experimentally by Van Atta & Wyngaard (1975). But both the hypothesis and its confirmation were challenged by Yakhot *et al.* (1989). Using the RNG methods, they obtained the scaling $E^{(2)} \sim k^{-\frac{7}{3}}$ for the spectrum u^2 in the inertial subrange. In their paper an experimental result $E^{(2)} \sim k^{-\frac{5}{3}}$ was attributed to the effect of Taylor's hypothesis.

We have attempted to analyse again the random sweeping hypothesis, and some relevant questions, both theoretically and experimentally. Theoretical analysis was accomplished in terms of the higher-order structure functions $D_{u_i}^{(m)}(r) =$

$\langle [u_i^m(x+r) - u_i^m(x)]^2 \rangle$. It was shown that the structure function analysis has two important advantages over the spectral one. The first advantage is that the derivation of the scaling $D^{(m)} \sim r^{\frac{2}{3}}$ is thoroughly clear and all assumptions are completely transparent. The second one is that assumptions, in particular the assumption about the statistical independence of large- and small-scale excitation, can be converted to 'measurable' form.

To verify the formulae obtained experimentally, measurements in two high Reynolds number laboratory shear flows were used: in the return channel ($R_\lambda \approx 3.2 \times 10^3$) and in the mixing layer ($R_\lambda \approx 2.0 \times 10^3$) of the large wind-tunnel at the Central Aerohydrodynamic Institute (Moscow). Two velocity components (in the direction of the mean flow and in the direction of the mean shear) were processed. The extent of a clearly established inertial subrange, taken as $30\eta \leqslant r \leqslant \frac{1}{8}L$, was about two decades in all flows. The higher-order structure functions were measured for $m = 1$–4. Special methods were devised to estimate the influence of Taylor's hypothesis. It was found that this influence on the parameters studied is negligible.

Experiments have shown that higher-order structure functions scale in the inertial subrange as $r^{\frac{2}{3}}$ in agreement with the random hypothesis. However, strong correlation between large-scale parameters u_i^k, $k = 2, 4, 6$, and small-scale ones $\Delta u_i^2(r)$ for any distance r within the inertial subrange was established. Thus the sweeping decorrelation hypothesis cannot be exactly valid, in spite of the scaling prediction agreeing with experiments. It was found the large–small scale correlation has two relevant consequences: it violates the sweeping relation (22) and all right-hand-side terms in the exact equation (7) become significant. Both these consequences contradict the sweeping hypothesis but they act in opposite directions and the final result becomes close to the hypothesis's prediction.

The assumption about statistical independence of large- and small-scale excitation was checked with conditionally averaged moments of velocity difference $\langle \Delta u_i^l \rangle_{u_i^*}$, $l = 2$–4, at fixed value of the large-scale parameter u_i^*. Clear dependence of the conditionally averaged moments on the level of averaging u_i^* was found. It was shown that this dependence does not really contradict Kolmogorov's hypotheses and it reflects the spatial and temporal variability of the energy flux through a cascade. In spite of a strong correlation between the energy-containing and the small-scale excitation, the intrinsic structure of the inertial subrange was found to be universal, i.e. it depends weakly on separation r but not on the flow type, direction of u_i, or Reynolds number.

The authors are deeply grateful to Professors P. Bradshaw, V. R. Kuznetsov and V. A. Sabelnikov and Dr P. Durbin for comments on the draft and to Dr R. Rogallo for fruitful discussions.

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