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1 **Scale-by-scale turbulent energy budget in the intermediate wake of two-dimensional**  
2 **generators**

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It is first established, on the basis of new X-wire measurements, that the equilibrium similarity of the terms in the scale-by-scale (s-b-s) budget of the turbulent energy  $\overline{q^2}$  is reasonably well approximated on the axis of the intermediate wake of a circular cylinder. The similarity, which scales on the Taylor microscale  $\lambda$  and  $\overline{q^2}$ , is then used to determine s-b-s energy budgets from the data of Antonia, Zhou & Romano<sup>1</sup> for 5 different two-dimensional wake generators. In each case, the budget is reasonably well closed, using the locally isotropic value of the mean energy dissipation rate, except near separations comparable to the wavelength of the coherent motion (CM). The influence of the initial conditions is first felt at a separation  $L_c$  identified with the cross-over between the energy transfer and large scale terms of the s-b-s budget. When normalized by  $\overline{q^2}$  and  $L_c$ , the mean energy dissipation rate is found to be independent of the Taylor microscale Reynolds number. The CM enhances the maximum value of the energy transfer, the latter exceeding that predicted from models of decaying homogeneous isotropic turbulence.

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## 9 I. INTRODUCTION

10 The concept of self-preservation, which effectively assumes that the flow is governed by  
11 single length and velocity scales, has an important role in describing the spatial evolution  
12 of turbulent flows. Implicit in this concept is the notion of forgetfulness or disentanglement  
13 from initial conditions. There is now however abundant evidence to suggest that initial  
14 conditions can influence the type of self-preservation that can be achieved in several types  
15 of turbulent flows, in particular free shear flows such as jets and wakes. The latter flows are  
16 especially interesting in view of the invariably strong organisation, in the form of vortical  
17 coherent structures, which is first observed immediately downstream of the wake-generating  
18 obstacle and which subsequently influences the streamwise development of the flow (*e.g.*  
19 Refs. 1–6).

20 Although the conditions for self-preservation, or sometimes self-similarity, have been in-  
21 vestigated by considering the mean momentum and turbulent energy equations at one point  
22 in space, it is of course more powerful and general to consider the self-preservation of the two-  
23 point velocity correlation tensor. The precursor of this approach was outlined by Ref. 7 who  
24 obtained a similarity solution of the two-point velocity correlation for homogeneous isotropic  
25 turbulence when the Reynolds number is very large. George<sup>8</sup> proposed an equilibrium simi-  
26 larity solution of Lin’s spectral equation, equivalent to the Kármán-Howarth equation, which  
27 is valid at any Reynolds number. This solution was tested in grid turbulence by Antonia  
28 *et al.*<sup>9</sup>, in the context of the transport equation for the turbulent energy structure function  
29  $\overline{(\Delta q)^2} = \overline{(\Delta u)^2} + \overline{(\Delta v)^2} + \overline{(\Delta w)^2}$ , where  $\Delta u$ ,  $\Delta v$ ,  $\Delta w$  are the velocity increments for  $u$ ,  $v$ ,  $w$   
30 between two points separated by a distance  $r$  (*e.g.*  $\Delta u = u(x+r) - u(x)$ ), and the overbar  
31 denotes time-averaging. A plausible interpretation of the transport equation of  $\overline{(\Delta q)^2}$  is  
32 that it represents an energy budget at each scale, or a scale-by-scale ( hereafter s-b-s) en-  
33 ergy budget. There have been several attempts at investigating s-b-s budgets, using either  
34 experimental or numerical data, in other types of flows, *e.g.* in homogeneous shear flow<sup>10–12</sup>,  
35 in the central region of a channel flow<sup>13</sup>, in a circular jet<sup>14–16</sup>, in wall-bounded flows<sup>17,18</sup>. In  
36 general, closure of the budget has been satisfactory validated by the data, allowing for the  
37 assumptions that are made when deriving the transport equations. In these previously cited  
38 flows, different physical phenomena are at play, for example grid turbulence is dominated by  
39 the advection of energy whilst diffusion is prominent near the centreline of a channel flow.

40 On the centreline of an axisymmetric jet, the major source of (large scale) inhomogeneity  
41 is associated with the streamwise decay of turbulent energy with only a small contribution  
42 from the production of energy arising from the difference in the Reynolds normal stresses.

43 In this paper, we examine the s-b-s energy budget in the intermediate wake of different  
44 obstacles. We focus on this region primarily because we expect the energy transfer in this  
45 region of the flow to be dominated, as in the case of grid turbulence and along the axis of  
46 a circular jet, by the streamwise decay of turbulent energy. However, unlike the latter two  
47 flows, one also expects that the transfer may feel the influence of the organised motion or  
48 coherent structures, the latter depending clearly on the type of initial condition, *viz.* the  
49 geometry of the wake-generating obstacle. Detailed velocity and vorticity measurements  
50 were made by Antonia, Zhou & Romano<sup>1</sup>, hereafter AZR, at  $x/d = 70$ , in wakes generated  
51 by 5 different bluff-bodies, each with the same characteristic dimension  $d$  at nominally the  
52 same Reynolds number based on the Taylor microscale  $R_\lambda \approx 200$ . Unfortunately, there was  
53 no attempt in Ref. 1 at measuring the one point energy budget or establishing if similarity  
54 applied in this region of the flow. This would have enabled the construction of the the s-b-s  
55 energy budget. To remedy this shortcoming, we present new measurements made at several  
56 streamwise locations in the intermediate wake of a circular cylinder. Experimental details  
57 are given in section II. The one-point energy budget (section III) is then obtained and the  
58 equilibrium similarity is tested (section IV) and exploited in section V to examine the effect  
59 of initial conditions on the s-b-s energy budget. Finally, conclusions are drawn in section  
60 VI.

## 61 II. EXPERIMENTAL DETAILS

62 X-wire measurements were made in the wake of a circular cylinder in the CORIA wind  
63 tunnel, which is of recirculating type. The residual turbulence level is smaller than 0.2%.  
64 The square test section is  $(0.4 \times 0.4\text{m}^2)$  has a length of 2.5m. The mean pressure gradient  
65 was set to zero. A circular cylinder of diameter  $d = 10\text{mm}$ , placed horizontally downstream  
66 the contraction, spanned the width of the test section. The upstream velocity  $U_0$  was  
67  $6.5\text{ms}^{-1}$ , corresponding to a Reynolds number  $Re_d$ , based on the cylinder diameter, of 4333.  
68 Measurements of velocity fluctuations  $u$  and  $v$  were made at several streamwise locations  
69 in the range  $40 \leq x/d \leq 110$ . The X-wire probe (Dantec 55P51) was calibrated using a

70 look-up table technique (*e.g.* Refs. 19 and 20) with velocity and angle increment of  $1\text{ms}^{-1}$   
71 and  $5^\circ$  respectively. The hot wires (Pt-W,  $5\mu\text{m}$  diameter,  $1.25\text{mm}$  length) were operated  
72 with a Dantec constant temperature bridge at an overheat ratio of 1.6. Voltage signals were  
73 passed through buck and amplifiers (SRS SIM983) and low-pass filtered (SRS SIM965) at a  
74 frequency  $f_c$  close to the Kolmogorov frequency  $f_K \equiv \bar{U}/2\pi\eta$  ( $\bar{U}$  is the local mean velocity  
75 and the Kolmogorov length-scale  $\eta = (\nu^3/\bar{\epsilon})^{1/4}$ ). The sampling frequency  $f_s = 20\text{kHz}$  and  
76 the sampling period was about  $4\text{min}$ . The air temperature in the wind tunnel was kept  
77 constant during the calibration and measurements using a regulating cross-flow water-air  
78 cooler system which allowed the mean temperature to be controlled to  $\pm 0.5^\circ\text{C}$ .

79 Measurements made at the University of Newcastle in the wakes downstream of 5 different  
80 two dimensional obstacles, each with the same characteristic dimension  $d = 25.4\text{mm}$  have  
81 been detailed in Refs. 1 and 21. Only a brief description is given below (see Fig. 1).

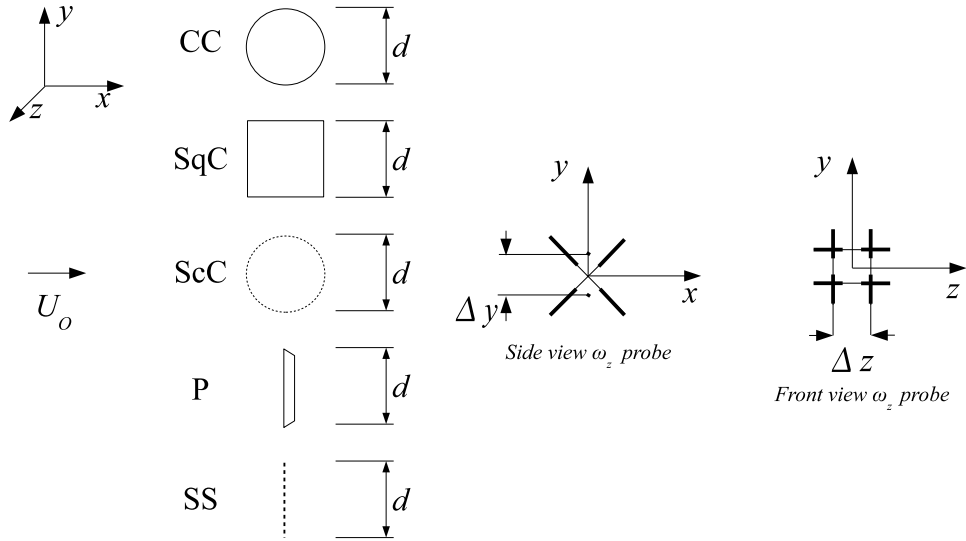


FIG. 1. Cross section of the five different generators used for the experiments carried out in Newcastle by Ref. 1. From top to bottom: CC, SqC, ScC, P, SS. Also sketched is a one-component vorticity probe in the  $\omega_z$  configuration ( $\omega_z$  is the vorticity component in the  $z$  direction).

82 The data were obtained in an open circuit wind tunnel with a  $2.4\text{m}$  long square test  
83 section ( $0.35 \times 0.35\text{m}^2$ ). The five generators (a circular cylinder CC, a square cylinder SqC,  
84 a screen cylinder ScC, a solid plate P normal to the flow and a screen strip SS, Fig. 1) were  
85 located  $20\text{cm}$  downstream of the end of the contraction and spanned the width of the test

86 section. All measurements were made on the centreline at  $x = 70d$  at several values of  $Re_d$ .  
87 A one-component vorticity probe (Fig. 1), which included an X-wire straddled by 2 parallel  
88 hot wires, was used to determine either  $\omega_z$  when the X-wire was in the plane  $(x,y)$  or  $\omega_y$   
89 when the X-wire was in the  $(x,z)$  plane. All the hot wires were etched from Pt-10%Rh to a  
90 diameter of  $2.5\mu\text{m}$  and an active length of approximately  $0.5\text{mm}$  (length to diameter ratio  
91  $\approx 200$ ). They were operated at an overheat ratio of 1.5 with in-house constant temperature  
92 anemometers. The anemometer output signals were passed through buck and gain circuits,  
93 prior to low-pass filtering at a frequency  $f_c \approx f_K$ . The filtered signals were then digitized  
94 at a frequency  $f_s = 2f_c$  into a PC using a 12bit A-D converter. The sampling period was  
95 in the range 45 to 100s. The mean velocity  $U_0$ , upstream of the wake generator, was varied  
96 so that the Taylor microscale Reynolds number  $R_{\lambda_u}$  ( $\equiv \overline{u^2}^{1/2} \lambda_u / \nu$ , where  $\lambda_u = 15\nu \sqrt{\overline{u^2}/\bar{\epsilon}}$  is  
97 the Taylor microscale) was in the range  $150 \lesssim R_{\lambda_u} \lesssim 300$  at  $x/d = 70$ .

### 98 III. ONE-POINT ENERGY BUDGET

99 It is important that the one-point energy budget is known relatively accurately since it  
100 represents the limiting state, as the separation  $r$  approaches the integral scale, of the s-b-s  
101 energy budgets. Here, we focus on the region along the axis of the intermediate wake. The  
102 general formulation of the one-point energy budget is

$$\frac{1}{2} \overline{U}_i \frac{\partial \overline{q^2}}{\partial x_i} + \overline{u_i u_j} \frac{\partial \overline{U}_i}{\partial x_j} + \frac{1}{2} \frac{\partial}{\partial x_j} \left( 2\overline{u_j p} + \overline{u_j q^2} \right) + \bar{\epsilon} = 0. \quad (1)$$

103 (the summation convention applies). The first term on left hand side of Eq. (1) represents the  
104 advection of turbulent kinetic energy by the mean velocity. The second term is interpreted  
105 as the production of kinetic energy through the mean velocity gradient. The third term  
106 corresponds to the diffusion of turbulent energy while the last term is the mean energy  
107 dissipation rate.

108 Profiles along the  $y$  axis of  $\overline{vq^2}$  approximated by  $2\overline{vu^2} + \overline{v^3}$  indicated that the diffusion  
109 of  $\overline{q^2}$  by  $v$  is negligible along the axis since  $\overline{vq^2} \approx \text{constant}$  on either side of the axis. A  
110 similar conclusion was arrived at by examining the data supplied to us by Zhou<sup>22</sup>. The  
111 latter measurements were made with a 3-component vorticity probe at  $x/d = 10, 20$  and  
112  $40$ . These data, with simultaneous information for  $u, v$  and  $w$ , indicated that  $\overline{vq^2}$  was more  
113 closely approximated by  $2\overline{vu^2} + \overline{v^3}$  than by  $\overline{vu^2} + 2\overline{v^3}$ , *i.e.* axisymmetry with respect to  $y$  is

114 superior to that with respect to  $x$ . The latter axisymmetry along the  $y$  axis for large-scale  
 115 statistics contrasts with Ref. 23 who observed a local axisymmetry with respect to the  $x$  axis  
 116 at the level of the smallest scales. However, the present result appears reasonable since the  
 117 influence of the Bénard-von Kármán street (a large-scale effect), which induces flow along  
 118 the  $y$  direction cannot be neglected in this region of the wake.

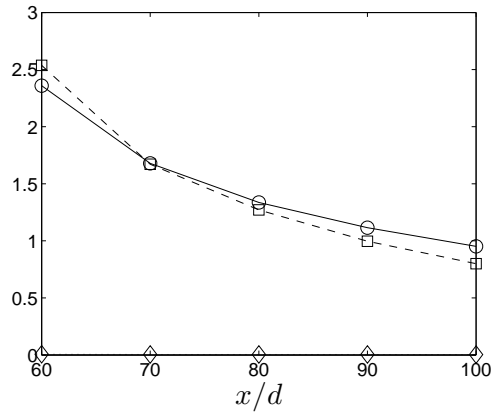


FIG. 2. Terms of the one-point kinetic energy budget (Eq. (2)) on the center-line of the wake of a circular cylinder normalized by  $U_0^3/d$ .  $\square$   $-\frac{1}{2}\bar{U}\frac{\partial \bar{q}^2}{\partial x}$ ,  $\circ$   $\bar{\epsilon}$ ,  $\diamond$   $(\overline{u^2} - \overline{v^2}) \frac{\partial \bar{U}}{\partial x}$ .

119 With the assumption that diffusion by the pressure  $p$  is also negligible on the axis, the  
 120 one-point energy budget Eq. (1) reduces to

$$\frac{1}{2}\bar{U}\frac{\partial \bar{q}^2}{\partial x} + (\overline{u^2} - \overline{v^2}) \frac{\partial \bar{U}}{\partial x} + \bar{\epsilon} = 0 \quad (2)$$

121 The distributions (Fig. 2) indicate that the production term associated with the difference  
 122 in the Reynolds normal stresses is negligible. Consequently, Eq. (2) further simplifies to

$$\frac{1}{2}\bar{U}\frac{\partial \bar{q}^2}{\partial x} + \bar{\epsilon} = 0 \quad (3)$$

123 Although Fig. 2 shows that the advection and energy dissipation rate data crossover  
 124 at  $x/d = 70$ , the magnitude of the two distributions are sufficiently close to each other to  
 125 claim that the energy budget over this range of  $x/d$  is quite similar to that for decaying  
 126 homogeneous isotropic turbulence (or HIT), as approximated by a turbulent grid flow. This  
 127 is a considerable simplification in the context of estimating the s-b-s energy budget since  
 128 only the decay term needs to be taken into account when  $r$  is sufficiently large.

## 129 IV. SCALE-BY-SCALE ENERGY BUDGETS

130 In view of the analogy with decaying turbulence HIT, the s-b-s budget is in essence  
 131 provided by the Kármán-Howarth equation<sup>7</sup> as detailed in Ref. 9. The latter authors  
 132 assumed equilibrium similarity, as proposed by Ref. 8. The similarity scales are  $\lambda$  and  $\overline{q^2}$ .  $\overline{q^2}$   
 133 decays according to a power-law, whose exponent depends on the initial conditions. Second-  
 134 and third-order structure functions, as they appear in the following transport equation<sup>24,25</sup>

$$-\overline{\Delta u(\Delta q)^2} + 2\nu \frac{\partial}{\partial r} \overline{(\Delta q)^2} - \frac{\overline{U}}{r^2} \int_0^r s^2 \frac{\partial}{\partial x} \overline{(\Delta q)^2} ds = \frac{4}{3} \epsilon r, \quad (4)$$

135 may be written as

$$f_q(\tilde{r}) = \frac{\overline{(\Delta q)^2}}{\overline{q^2}}, \quad g_q(\tilde{r}) = -\overline{\Delta u(\Delta q)^2} \frac{3^{1/2} R_\lambda}{(\overline{q^2})^{3/2}} \quad (5)$$

136 where  $\tilde{r} = r/\lambda$  and  $s$  in Eq. (4) is a dummy integration variable. Note that  $f_q$  and  $g_q$  depend  
 137 only on  $\tilde{r}$ . Using Eq.(5), Eq. (4) can be rewritten as

$$g_q + 2f'_q - \left( \frac{5\Gamma_{q1}}{m_q} - 10\Gamma_{q2} \right) \tilde{r}^{-2} = \frac{20}{3} \tilde{r} \quad (6)$$

138 where

$$\Gamma_{q1} = \int_0^{\tilde{r}} \tilde{s}^3 f'_q d\tilde{s}, \quad \Gamma_{q2} = \int_0^{\tilde{r}} \tilde{s}^2 f_q d\tilde{s}. \quad (7)$$

139 The prime denotes differentiation with respect to  $\tilde{r}$ . The parameter  $m_q$  is the power-law  
 140 decay exponent of  $\overline{q^2}$ , *i.e.*

$$\overline{q^2} \propto (x - x_0)^{m_q} \quad (8)$$

141 where  $x_0$  is a virtual origin. Eq. (3) implies that  $\overline{\epsilon}$  will also decay in a power-law fashion

$$\overline{\epsilon} \propto (x - x_0)^{m_\epsilon} \quad (9)$$

142 with  $m_\epsilon = m_q - 1$ . Since  $\lambda^2 = 5\nu\overline{q^2}/\overline{\epsilon}$ , it follows that the exponent  $m_\lambda$ , *viz.*

$$\lambda \propto (x - x_0)^{m_\lambda} \quad (10)$$

143 should be equal to 0.5. The expected linearity of  $\lambda^2$  with respect to  $x$  was used to estimate  
 144  $x_0$  ( $\approx 8d$ ). The distributions, along the wake centreline of  $\overline{u^2}$ ,  $\overline{q^2}$ ,  $\overline{\epsilon}$ ,  $\lambda$  and  $\overline{U_d^2}$  ( $\overline{U_d}$  is the  
 145 maximum velocity defect) are shown in Fig. 3. Note that  $m_{U_d}$  ( $\overline{U_d^2} \propto (x - x_0)^{m_{U_d}}$ ) is -0.68,  
 146 which is substantially different from the value of -1, expected in the far wake. Also  $R_\lambda$



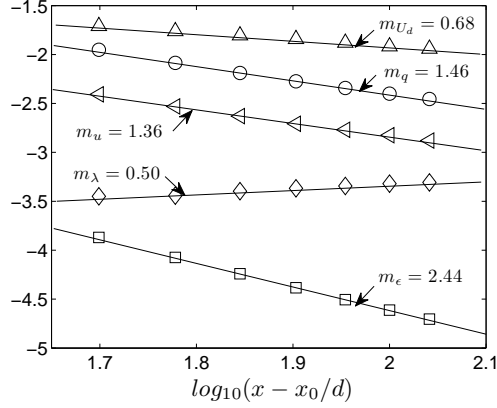


FIG. 3. Decay exponents on the wake center-line of the circular cylinder wake ;  $\circ \log_{10}(\overline{q^2}/U_0^2)$ ,  $\triangleleft \log_{10}(\overline{u^2}/U_0^2)$ ,  $\square \log_{10}(\overline{\epsilon d}/U_0^3)$ ,  $\diamond \log_{10}(10^{-1}\lambda/d)$ ,  $\triangle \log_{10}(U_d^2/U_0^2)$ , solid lines are linear fits to the data.

147 must decay with  $x$  through the intermediate wake ( $R_\lambda \propto (x - x_0)^{-0.23}$ ) whereas it should be  
 148 independent of  $x$  in the far wake.

149 For convenience, Eqs.(6) can be rewritten as

$$G_q + D_q + I_q = \frac{4}{3} \quad (11)$$

150 where  $G_q = g_q \tilde{r}^{-1}/5$ ,  $D_q = 2f'_q \tilde{r}^{-1}/5$  and  $I_q = -(\Gamma_{q1}/m_q - 2\Gamma_{q2}) \tilde{r}^{-3}$ . Each term can be  
 151 estimated from measurements made at one spatial location provided  $m_q$  is available.

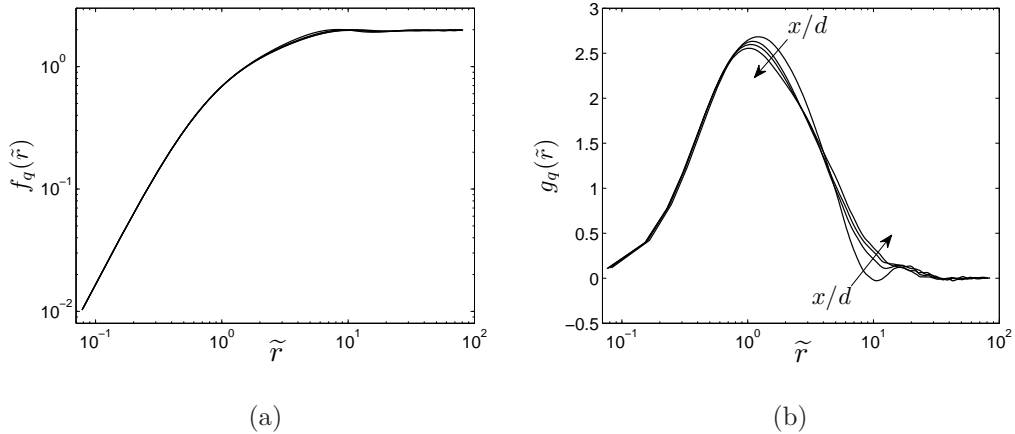


FIG. 4. (a)  $f_q$  and (b)  $g_q/\tilde{r}$  as a function of  $\tilde{r}$ , for  $x$  between  $50d$  and  $110d$  of the circular cylinder wake.

152 Measured distributions of  $f_q$  in Fig. 4(a) confirm that the collapse, based on  $\lambda$  and  
 153  $\overline{q^2}$ , is quite satisfactory. Note that  $f_q$  approaches its limiting value ( $= 2$ ) at large  $\tilde{r}$  after  
 154 overshooting it, due mainly to the contribution from  $\overline{(\Delta v)^2}$ , which is affected by the coherent  
 155 motion<sup>21</sup>. The quality of the collapse for  $g_q$  (Fig. 4(b)) is not as good as that in Fig. 4(a); it  
 156 is however at least as good, if not better, than that reported by Ref. 9 for grid turbulence.  
 157 Note also that the curve of  $g_q$  at  $x = 50d$  differs quite significantly from those beyond  
 158  $x = 70d$  for which the collapse is much more satisfactory irrespectively of the downstream  
 159 distance. Particularly, at a separation  $\tilde{r} \approx 10$ ,  $g_q$  reaches negative values which is most likely  
 160 due to the presence of the organized motion. This may highlight that a transition in the  
 161 behaviour of the coherent motion occurs between  $x = 50d$  and  $70d$ .

## 162 V. RESULTS FOR SCALE-BY-SCALE ENERGY BUDGETS IN THE 5 163 WAKES

164 Since all measurements in the five different wakes were made<sup>1</sup> at only one streamwise  
 165 location ( $x/d = 70$ ), equilibrium similarity, which was verified in the previous section for  
 166 the wake of the circular cylinder, will be used, together with the value of  $m_q$  reported in  
 167 Fig. 3, to construct the s-b-s budgets for these 5 wakes. Although one would expect some  
 168 dependence of  $m_q$  on both  $Re_d$  and the initial conditions, we have assumed that it remains  
 169 constant for the present purpose, since its influence on the s-b-s energy budget, especially  
 170 for  $R_\lambda \approx 200$  is likely to be small.

171 We first consider the distributions of  $f_q$  (Fig. 5) for the five different wakes. The main  
 172 point which emerges from this figure is that the effect of initial conditions cannot be dis-  
 173 counted, at least when  $\tilde{r}$  is sufficiently large. The collapse at small  $\tilde{r}$  is almost as good as  
 174 that reported at small  $r^*$  ( $\equiv r/\eta$ ) in Fig. 7 of<sup>1</sup>, where  $\overline{(\Delta u^*)^2}$  and  $\overline{(\Delta v^*)^2}$  were plotted as  
 175 functions of  $r^*$  (the asterisk stands for the normalization by the Kolmogorov scales). How-  
 176 ever, although Kolmogorov scaling may be superior in the dissipative range, scaling on  $\lambda$   
 177 and  $\overline{q^2}$  extends the collapse to a wider range of scales. In general, distributions for different  
 178 wakes depart from each other at  $\tilde{r} \approx 2$ . Note also the significant overshoot in  $f_q = \overline{(\Delta q)^2}/\overline{q^2}$   
 179 for the screen strip (SS) with a maximum at  $\tilde{r} = 12$ . Included in Fig. 5 are the distributions  
 180 (thin and bold curves) obtained from the models of Kurien & Sreenivasan<sup>26</sup> (hereafter KS)  
 181 and Aivalis *et al.*<sup>27</sup> (henceforth A). The latter models rely on a parametric equation for

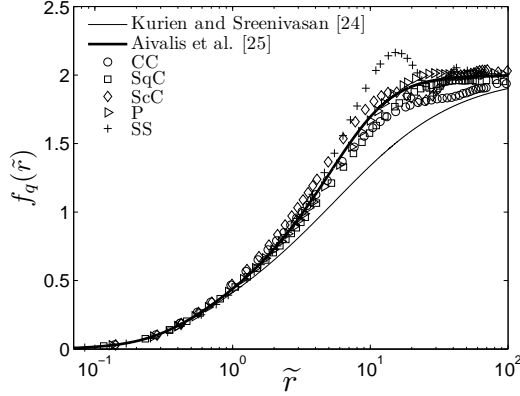


FIG. 5. Measured distributions of  $f_q$  as a function of  $\tilde{r}$  for the five different wakes. Also shown is  $f_q$  inferred from KS<sup>26</sup> (thin line) and A<sup>27</sup> (bold line). The Taylor microscale Reynolds number is approximatively the same in each case, *viz.*  $R_{\lambda_u} \approx 195$ .  $\circ$  CC,  $\square$  SqC,  $\diamond$  ScC,  $\triangleright$  P,  $+$  SS.

182 second-order structure functions

$$\overline{(\Delta u)^{*2}}|_K = \frac{r^{*2} (1 + \beta r^{*})^{(2c-2)}}{15 (1 + \alpha r^{*2})^c} \quad (12)$$

$$\overline{(\Delta u)^{*2}}|_A = \frac{r^{*2} (1 + (\beta r^{*})^2)^{(c-1)}}{15 (1 + \alpha r^{*2})^c}, \quad (13)$$

183 where the subscript  $K$  and  $A$  refer to the models of KS<sup>26</sup> and A<sup>27</sup> respectively. In Eqs. (12)  
 184 and (13),  $\beta = L^{*-1}$  where  $L^* = C_\epsilon R_\lambda^{3/2} 15^{-3/4}$  ( $C_\epsilon = 1$  as in Refs. 9 and 28) is a measure of  
 185 the integral length-scale.  $\alpha = (15C_u)^{-3/2}$  ( $C_u = 2$  was chosen<sup>9,28</sup>) designates the cross-over  
 186 between the dissipative and inertial range. Finally,  $c = 1 - \zeta_u/2$  where  $\zeta_u$  is the inertial  
 187 range exponent of the longitudinal second order structure function. In contrast to Danaila  
 188 *et al.*<sup>29</sup> or Thiesset *et al.*<sup>30</sup>, for which  $\zeta_u$  is function of the turbulent Reynolds number  $R_\lambda$ ,  
 189 we assume here  $\zeta_u = 2/3$  as in Refs. 9 and 28. Then,  $\overline{(\Delta q)^2}$  is calculated using the isotropic  
 190 relation

$$\overline{(\Delta q)^{*2}}|_{K,A} = (3 + r^* \frac{\partial}{\partial r^*}) \overline{(\Delta u)^{*2}}|_{K,A}. \quad (14)$$

191 After inserting Eq. (14) in Eq. (11), one can calculate the third-order structure function by  
 192 difference. These models are thus only parametrized by the Reynolds number  $R_\lambda$  and the  
 193 decay exponent  $m_q$ .

194 A comparison of Eqs.(13) and (12) indicates that the two models differ only through their

195 numerator, *i.e.*

$$\frac{\overline{(\Delta u)^{*2}}|_K}{\overline{(\Delta u)^{*2}}|_A} = \left[ \frac{1 + 2\beta r^* + (\beta r^*)^2}{1 + (\beta r^*)^2} \right]^{(c-1)} \quad (15)$$

196 As  $c < 1$ ,  $\overline{(\Delta q)^{*2}}|_K < \overline{(\Delta q)^{*2}}|_A$  for every separation  $r$  as observed in Fig. 5. Thus,  $I_q|_K < I_q|_A$   
 197 and so  $G_q|_K > G_q|_A$ . Nevertheless, in the limit of large scales

$$\lim_{r \rightarrow \infty} \overline{(\Delta q)^{*2}}|_K = \lim_{r \rightarrow \infty} \overline{(\Delta q)^{*2}}|_A = \frac{6R_\lambda}{15^{1/2}} \quad (16)$$

198 and for small separations

$$\lim_{r \rightarrow 0} \overline{(\Delta q)^{*2}}|_K = \lim_{r \rightarrow 0} \overline{(\Delta q)^{*2}}|_A = \frac{r^{*2}}{3} \quad (17)$$

199 Therefore, the two models lead to a same asymptotic distributions of energy in the limit  
 200 of small and large separations. The main difference between the two models occurs at  
 201 intermediate scales.

202 We now turn our attention to the different terms in the s-b-s energy budgets displayed in  
 203 Fig. 6. All the terms in Eq. (11) are shown in this figure. Differences in the budgets between  
 204 the 5 wakes tend to occur when  $\tilde{r} \geq 2$ . This separation corresponds approximatively to the  
 205 cross-over  $L_c$  (identified by vertical arrow in Fig. 6) between the transfer term and the large  
 206 scale decay or inhomogeneous term. It follows that the energy distribution becomes more  
 207 sensitive to the initial conditions the larger  $\tilde{r}$  is by comparison to  $L_c$ . This distribution is  
 208 more influenced by the contribution of the relatively organized large scales which are more  
 209 likely to remember the initial conditions.

210  $G_q$  is much more affected by whether the wake generator is porous or solid than the other  
 211 terms. For the porous wakes, the magnitude of  $G_q$  remains larger than that predicted by the  
 212 two models when  $\tilde{r} \geq 2$ . This reflects the relatively strong organization in these two wakes  
 213 and hence the stronger influence of the initial conditions. Note that the A<sup>27</sup> model fails  
 214 to adequately represent  $G_q$  in these two cases (for  $\tilde{r} \geq 2$ ) whereas it represents adequately  
 215 the other terms ( $I_q$  and  $D_q$ ) and is superior to the KS model<sup>26</sup>. The sum  $G_q + I_q + D_q$  is  
 216 close to the expected value of 4/3 for the solid generators. There is however a significant  
 217 departure, typically in the range  $2 \leq \tilde{r} \leq 10$ , for the porous generators and especially SS. For  
 218 these two wakes, the local anisotropy is greater than for the other 3 wakes<sup>21</sup>, reflecting the  
 219 larger degree of organisation in the porous generator wakes than in the impervious generator  
 220 wakes<sup>21</sup>.

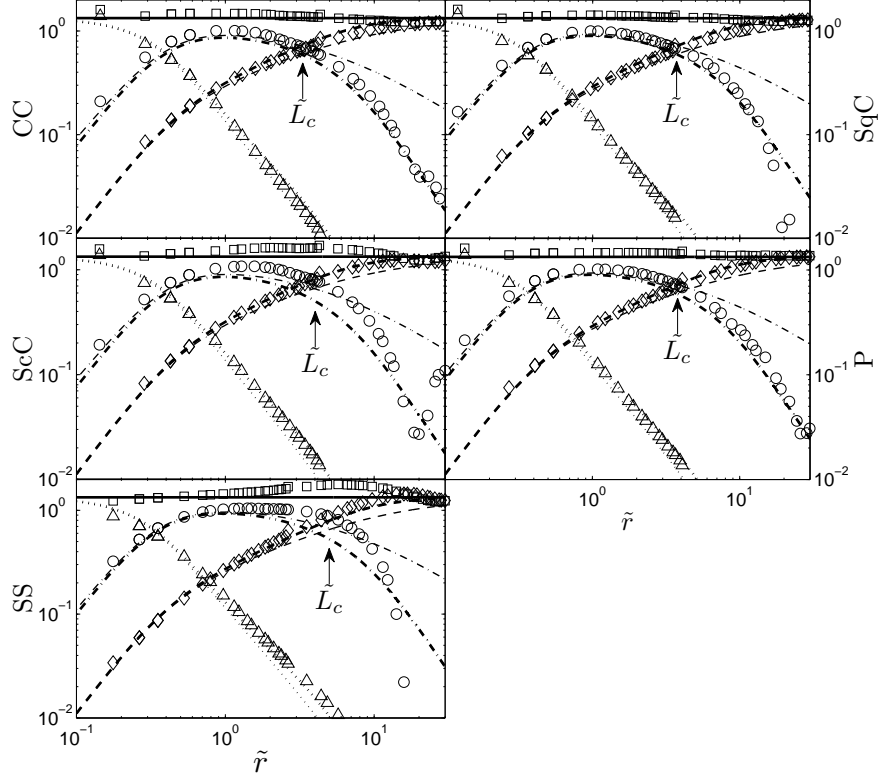


FIG. 6. Scale-by-scale energy budgets Eq. (11) in the five different wakes. Thin curves correspond to the model of KS<sup>26</sup> and bold curves to the model of A<sup>27</sup>. Symbols denote the experimental data.  $\circ$ ,  $-\cdot-$ :  $G_q$ ;  $\triangleright$ ,  $\cdots$ :  $D_q$ ;  $\diamond$ ,  $- - -$ :  $I_q$ ;  $\square$ ,  $—$ :  $(G_q + D_q + I_q)$ . The thick horizontal line corresponds to  $4/3$ .

221 The crossover length scale  $L_c$  has been identified with the separation at which  $G_q = I_q$ .  
 222 From Eq. (6), we can thus write

$$\tilde{L}_c = \left[ \frac{1}{g_q(\tilde{L}_c)} \left( \frac{5\Gamma_{q1}(\tilde{L}_c)}{m_q} - 10\Gamma_{q2}(\tilde{L}_c) \right) \right]^{1/2}. \quad (18)$$

223 One expects  $\tilde{L}_c$  to depend on initial conditions since  $g_q$ ,  $\Gamma_{q1}$ ,  $\Gamma_{q2}$  and  $m_q$  should all be affected  
 224 by the initial conditions. A plausible interpretation of  $L_c$  is that it represents a length scale  
 225 at which the energy injected in the system starts being transferred towards small scales. It  
 226 thus seems appropriate to normalize the mean energy dissipation rate  $\bar{\epsilon}$  by  $L_c$  and  $\overline{q^2}$ , *i.e.*

$$\bar{\epsilon} = C_{ec} \frac{\left( \overline{q^2}/3 \right)^{3/2}}{L_c}. \quad (19)$$

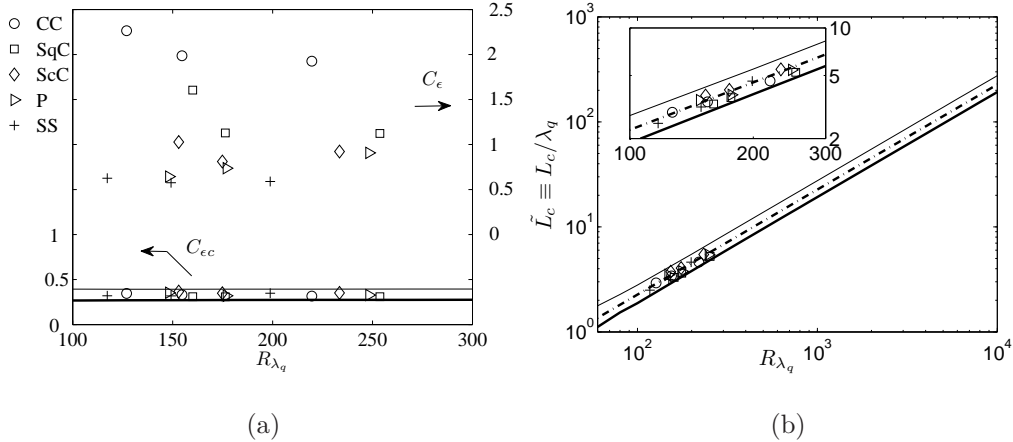


FIG. 7. (a) Evolution of the normalized dissipation rate parameters  $C_\epsilon$  and  $C_{\epsilon c}$  with  $R_\lambda$  for the five wakes. Also presented is the expectation for  $C_{\epsilon c}$  using either the model of A<sup>27</sup> (thick lines) or KS<sup>26</sup> (thin lines). (b) Evolution of the crossover length scale  $L_c$  normalized by  $\lambda$  with the Taylor microscale Reynolds number  $R_\lambda$ . Models; — A<sup>27</sup> (thick lines), — KS<sup>26</sup> (thin lines), — · —  $L_s/\lambda_q = C_{\epsilon c}R_{\lambda q}/15$  using  $C_{\epsilon c} = 0.34$ . Experiments  $\circ$  CC,  $\square$  SqC,  $\diamond$  ScC,  $\triangleright$  P, + SS.

227 This differs from the classical definition of the normalized dissipation rate parameter  $C_\epsilon$

$$\bar{\epsilon} = C_\epsilon \frac{(\overline{u^2})^{3/2}}{L_u}, \quad (20)$$

228 where  $L_u$  is the longitudinal integral length-scale. Fig. 7(a) shows the evolution of the nor-  
 229 malized dissipation rate parameter, defined with either Eq. (20) or (19). As was emphasized  
 230 by Ref. 1,  $C_\epsilon$  reaches very different values depending on the type of obstacle. Indeed,  $C_\epsilon$   
 231 is equal to about 2 for CC wake and has a quite different value ( $\approx 0.6$ ) in the SS wake.  
 232 On the other hand, when the normalization of  $\bar{\epsilon}$  is based on  $L_c$  and  $\overline{q^2}$ ,  $C_{\epsilon c}$  appears to be  
 233 remarkably universal, *i.e.* it is independent of both the type of obstacle and the Reynolds  
 234 number. The constancy of  $C_{\epsilon c}$  is also reasonably well reproduced by the two models. The  
 235 measured value of  $C_{\epsilon c}$  is equal to 0.34 whereas the values predicted by the KS and A models  
 236 are equal to 0.28 and 0.4 respectively.

237 It is evident that  $C_{\epsilon c}$  has a much stronger claim to universality than  $C_\epsilon$ . This is due in part  
 238 to the fact that  $L_c$  is much more precisely determined than  $L_u$  but more importantly because  
 239  $L_c$  is an intermediate scale between that at which energy is injected into the flow and that  
 240 ( $r \sim \lambda$ ) at which the transfer of energy is maximum. In other words, the cross-over between  
 241  $G_q$  and  $I_q$  defines the scale at which the energy associated with all the inhomogeneous large

242 scales, arising from the way energy is injected into the flow, is exactly balanced by the energy  
 243 transferred to the dissipative scales.

244 The use of  $L_c$  and  $\overline{q^2}$  as the normalization scales for the normalized dissipation rate  
 245 parameter has two main practical advantages. First, replacing  $\overline{u^2}$  by  $\overline{q^2}$  allows the large-  
 246 scale anisotropy to be accounted for. Second, the estimation of  $L_c$  is much less arbitrary  
 247 than that of  $L_u$ , the latter being usually inferred by integrating the longitudinal correlation  
 248 function to its first zero crossing.

249 It seems more appropriate to consider the ratio  $L_c/\lambda$  since  $\lambda$  is the scale at which the  
 250 energy transfer is maximum. This ratio can be thought of as a measure of the separation  
 251 between large scale and inertial range effects. By using Eq. (19) and since  $\lambda_q = 5\nu\overline{q^2}/\overline{\epsilon}$ , we  
 252 obtain

$$L_c/\lambda_q = \frac{1}{15}C_{cc}R_{\lambda_q}. \quad (21)$$

253 The magnitude of  $\tilde{L}_c$ , identified with the values of  $\tilde{r}$  for which  $G_q = I_q$  is shown in Fig.  
 254 7(b). Also included in this figure are the values of  $\tilde{L}_c$  which correspond to the two models.  $\tilde{L}_c$   
 255 is always smaller for the A model<sup>27</sup> than for the KS model<sup>26</sup>. The measured values of  $\tilde{L}_c$  lie  
 256 between the two models; although the range of  $R_\lambda$  is small, all the data are quite consistent  
 257 with a linear variation of  $\tilde{L}_c$  with  $R_\lambda$  and are very close to the line  $L_c/\lambda_q = C_{cc}R_{\lambda_q}/15$ .

258 The interpretation for  $L_c$  provided here is similar to that for the shear length-scale  $L_S$ ,  
 259 first introduced by Corrsin<sup>31</sup> in the context of shear-dominated flows, and further physically  
 260 interpreted in Refs. 11, 17, 18, and 32. Indeed, the authors emphasized that  $L_S \equiv \sqrt{S^3/\overline{\epsilon}}$ <sup>31</sup>  
 261 corresponds to the crossover between contributions from the production term due to the  
 262 mean shear  $S$  and the nonlinear transfer term in the spherically averaged scale-by-scale  
 263 kinetic energy budget. Note also that the present authors have recently generalized the def-  
 264 inition of  $L_S$  in order to account for the additional effect of the coherent strain using similar  
 265 arguments as in Ref. 11, *i.e.* based on scale-by-scale budgets (see Ref. 21). This suggests  
 266 that there exists a large variety of different relevant injection length-scales depending on  
 267 the flow under consideration. For example, as far as wake flows are concerned,  $L_c$  and  $L_S$   
 268 coexist,  $L_c$  being the crossover between the decay term and the transfer term and  $L_S$  the  
 269 crossover between the production (due to both the mean shear and the coherent shear) and  
 270 the transfer term.

271 In addition, even though we observe that  $L_c$  restores universality for the dissipation rate,

272 its applicability for modeling the dissipation rate has the shortcoming that this applies only  
 273 to decaying flows. In shear-dominated turbulence,  $L_S$  is likely to be the relevant injection  
 274 length-scale for normalizing the dissipation rate. This point is worth investigating in the  
 275 future.

276 Finally,  $L_c$  represents the largest length-scale that is in essence affected by the shape of  
 277 the obstacle. Therefore, as shown in Ref. 32 for shear flows, for a proper implementation of  
 278 a Large Eddy Simulation (LES),  $L_c$  represents the threshold scale that has to be simulated,  
 279 scales smaller than  $L_c$  having the best prospect of being modeled by a universal function.

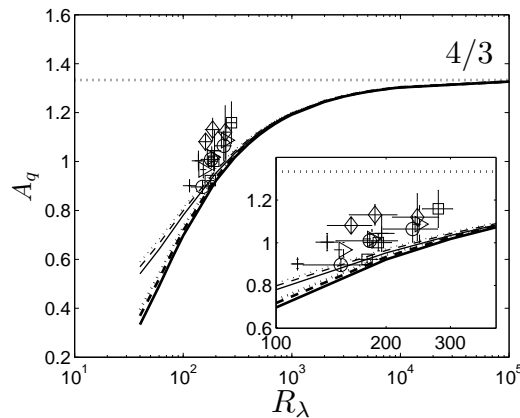


FIG. 8. Evolution of  $A_q$  with  $R_\lambda$  and the effect of  $m_q$  on  $A_q$ . Thick and thin curves correspond to the models of A<sup>27</sup> and KS<sup>26</sup> respectively. —:  $m_q = 1.25$ , - - -:  $m_q = 1.5$ , ·····:  $m_q = 1.75$ .  
 ○ CC, □ SqC, ◇ ScC, ▷ P, + SS.

280 We now turn our attention to the maximum value of the energy transfer  $A_q \equiv \max(\overline{(\Delta u)(\Delta q)^2}/\bar{\epsilon}r)$   
 281 depicted in Fig. 8. Vertical and horizontal error bars are also given in order to appraise  
 282 uncertainties in assessing  $A_q$ . Vertical error bars correspond to the standard deviation  
 283 between different (isotropic) estimations of the dissipation rate. From the one-component  
 284 vorticity probe oriented in either the  $x - y$  plane or  $x - z$  plane, one can calculate  $\bar{\epsilon}$  from



285 the following six expressions

$$\bar{\epsilon}_{iso} = 15\nu \overline{\left(\frac{\partial u}{\partial x}\right)^2} \quad (22)$$

$$\bar{\epsilon}_q = 3\nu \left[ \overline{\left(\frac{\partial u}{\partial x}\right)^2} + \overline{\left(\frac{\partial v}{\partial x}\right)^2} + \overline{\left(\frac{\partial w}{\partial x}\right)^2} \right] \quad (23)$$

$$\bar{\epsilon}_{\omega_z} = \nu \left( 2\overline{\omega_z^2} + \overline{\omega_y^2} \right) \quad (24)$$

$$\bar{\epsilon}_{\omega_y} = \nu \left( \overline{\omega_z^2} + 2\overline{\omega_y^2} \right) \quad (25)$$

$$\bar{\epsilon}_{xy} = \nu \left[ 6\overline{\left(\frac{\partial u}{\partial x}\right)^2} + 3\overline{\left(\frac{\partial u}{\partial y}\right)^2} + 2\overline{\left(\frac{\partial v}{\partial x}\right)^2} + \overline{\left(\frac{\partial u}{\partial y} \frac{\partial v}{\partial x}\right)} \right] \quad (26)$$

$$\bar{\epsilon}_{xz} = \nu \left[ 6\overline{\left(\frac{\partial u}{\partial x}\right)^2} + 3\overline{\left(\frac{\partial u}{\partial z}\right)^2} + 2\overline{\left(\frac{\partial w}{\partial x}\right)^2} + \overline{\left(\frac{\partial u}{\partial z} \frac{\partial w}{\partial x}\right)} \right]. \quad (27)$$

286 Vertical error bars can thus be interpreted as a measure of anisotropy at the level of  
 287 the smallest scales.  $A_q$  is calculated using  $\bar{\epsilon}_{iso}$  which appears to be close to the mean value  
 288 between these 6 different surrogates. Horizontal error bars are calculated from the difference  
 289 between  $R_{\lambda_u} = \sqrt{u^2}\lambda_u/\nu$  and  $R_{\lambda_q} = \sqrt{q^2}\lambda_q/\sqrt{3}\nu$ , where  $\lambda_u^2 = 15\nu\overline{u^2}/\bar{\epsilon}_{iso}$  and  $\lambda_q^2 = 5\nu\overline{q^2}/\bar{\epsilon}_{iso}$ .  
 290 The ratio between  $R_{\lambda_u}$  and  $R_{\lambda_q}$  is thus

$$\frac{R_{\lambda_u}}{R_{\lambda_q}} = \sqrt{\frac{3\overline{u^2}}{\overline{q^2}}} \left( \frac{\lambda_u}{\lambda_q} \right) = \frac{3\overline{u^2}}{\overline{q^2}} \quad (28)$$

291 which gives a measure of the anisotropy at the largest scales. Measured values of  $A_q$  are also  
 292 compared together with those inferred from the A and KS models, for three different values  
 293 of  $m_q$  (1.25, 1.5 and 1.75). Note that variations in the magnitude of  $m_q$  has virtually no  
 294 effect on the way  $A_q$  approaches the asymptote of 4/3. The influence of the decay exponent  
 295 is more discernible when  $R_\lambda$  is small, as the contribution from the large scale term becomes  
 296 more pronounced. For  $R_\lambda$  larger than 100, the differences are not perceptible and the choice  
 297 of a single value of  $m_q$  for the 5 wakes seems therefore reasonable.

298 Measured values of  $A_q$  are systematically larger than those predicted from the models  
 299 of decaying isotropic turbulence by on average 15%. The two porous obstacles, for which  
 300 CM is more prominent, appear to produce a higher maximum value of the energy transfer,  
 301 independently of the Reynolds number. This observation indicates that the energy transfer  
 302 is affected by the presence of the coherent motion, whose effect is to locally enhance the  
 303 cascade mechanism. A similar observation was made by Ref. 33, for the wake of a circular  
 304 cylinder at  $x = 40d$  and a smaller Reynolds number ( $R_\lambda \approx 70$ ).

## 305 VI. CONCLUSIONS

306 There are three major conclusions which arise from this work, all pertaining to the inter-  
 307 mediate region of a turbulent wake.

308 (i) Firstly, the X-wire measurements at various streamwise locations along the axis of  
 309 the wake indicate that equilibrium similarity, based on the Taylor microscale  $\lambda$  and the  
 310 turbulent energy  $\overline{q^2}$ , is verified satisfactorily. In particular, the collapse of the energy transfer  
 311 term  $\overline{\Delta \tilde{u}(\Delta \tilde{q})^2} R_\lambda \overline{q^2}^{-3/2} \tilde{r}^{-1}$  appears to be of better quality than that previously obtained for  
 312 decaying turbulence downstream a grid<sup>9</sup>. In the intermediate wake and for large separations,  
 313 the one-point energy budget is closely approximated by a balance between advection and  
 314 energy dissipation rate. This is substantially different to what is generally observed in the  
 315 far field<sup>34</sup> in which the diffusion term has also to be taken into account. This difference  
 316 between the intermediate and far wake might be related to a reorganization or rescaling  
 317 of the large-scale structures<sup>4,35,36</sup>, sometimes referred to as the secondary vortex street.  
 318 Further investigations of velocity statistics and corresponding length-scales obtained over a  
 319 sufficiently large distance between the near and far fields are needed to confirm this.

320 (ii) Secondly, the equilibrium similarity is used to estimate the s-b-s budgets, also along  
 321 the centreline, for wakes generated by five two-dimensional bodies of different shapes. The  
 322 closure of the budgets is less satisfactory for the two porous bodies (screen strip and screen  
 323 cylinder) than for the solid bodies (cylinder, plate and square), reflecting the persistence  
 324 of the coherent motion arising from the Kelvin-Helmholtz instability in the shear layers  
 325 associated with the porous bodies.

326 Initial conditions are first felt at a separation  $L_c$ , which corresponds to the intersection  
 327 point between the energy transfer and the large scale terms. For  $r < L_c$ , the shape of the  
 328 generator is not important and the statistics are similar to those for decaying homogeneous  
 329 isotropic turbulence. The parameter  $C_{ec}$  which involves normalization by  $L_c$  and  $\overline{q^2}$ , is  
 330 remarkably independent of  $R_\lambda$  and almost unaffected by the initial conditions. Clearly,  
 331 the ramification of this universality and its possible extension to other flows deserves to be  
 332 investigated further.  $L_c/\lambda$  varies linearly with  $R_\lambda$ , a result in accord with existing models  
 333 of  $\overline{(\Delta q)^2}$  as well as our measurements.

334 (iii) Finally, the maximum value of the energy transfer term is systematically larger  
 335 than that predicted by models for decaying HIT. This reflects the presence of the coherent

336 motion. The latter observation is important in the context of providing some insight into the  
337 way the coherent motion interacts with small scales. Some progress has been made on this  
338 interaction<sup>33</sup>, by quantifying the contribution of the coherent motion to the energy transfer  
339 after writing the s-b-s budgets for both the coherent and random motions.

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