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## <sup>1</sup> Scale-by-scale turbulent energy budget in the intermediate wake of two-dimensional <sup>2</sup> 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

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

Although the conditions for self-preservation, or sometimes self-similarity, have been in-20 <sup>21</sup> vestigated by considering the mean momentum and turbulent energy equations at one point <sup>22</sup> in space, it is of course more powerful and general to consider the self-preservation of the two-<sup>23</sup> point velocity correlation tensor. The precursor of this approach was outlined by Ref. 7 who <sup>24</sup> obtained a similarity solution of the two-point velocity correlation for homogeneous isotropic <sup>25</sup> turbulence when the Reynolds number is very large. George<sup>8</sup> proposed an equilibrium simi-<sup>26</sup> 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 <sup>28</sup> et al.<sup>9</sup>, in the context of the transport equation for the turbulent energy structure function  ${}_{^{29}}\overline{(\Delta q)^2} = \overline{(\Delta u)^2} + \overline{(\Delta v)^2} + \overline{(\Delta w)^2} , \text{ where } \Delta u, \, \Delta v, \, \Delta w \text{ are the velocity increments for } u, \, v, \, w$ 30 between two points separated by a distance r (e.g.  $\Delta u = u(x+r) - u(x)$ ), and the overbar <sup>31</sup> denotes time-averaging. A plausible interpretation of the transport equation of  $(\Delta q)^2$  is <sup>32</sup> that it represents an energy budget at each scale, or a scale-by-scale (hereafter s-b-s) en-<sup>33</sup> ergy budget. There have been several attempts at investigating s-b-s budgets, using either <sup>34</sup> 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 <sup>36</sup> 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 <sup>38</sup> flows, different physical phenomena are at play, for example grid turbulence is dominated by <sup>39</sup> the advection of energy whilst diffusion is prominent near the centreline of a channel flow. <sup>40</sup> On the centreline of an axisymmetric jet, the major source of (large scale) inhomogeneity <sup>41</sup> is associated with the streamwise decay of turbulent energy with only a small contribution <sup>42</sup> from the production of energy arising from the difference in the Reynolds normal stresses.

In this paper, we examine the s-b-s energy budget in the intermediate wake of different 43 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 <sup>48</sup> coherent structures, the latter depending clearly on the type of initial condition, *viz.* the <sup>49</sup> geometry of the wake-generating obstacle. Detailed velocity and vorticity measurements <sup>50</sup> 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 s2 same Reynolds number based on the Taylor microscale  $R_{\lambda} \approx 200$ . Unfortunately, there was <sup>53</sup> no attempt in Ref. 1 at measuring the one point energy budget or establishing if similarity <sup>54</sup> applied in this region of the flow. This would have enabled the construction of the the s-b-s <sup>55</sup> energy budget. To remedy this shortcoming, we present new measurements made at several <sup>56</sup> 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 <sup>58</sup> equilibrium similarity is tested (section IV) and exploited in section V to examine the effect <sup>59</sup> of initial conditions on the s-b-s energy budget. Finally, conclusions are drawn in section 60 VI.

#### 61 II. EXPERIMENTAL DETAILS

<sup>62</sup> X-wire measurements were made in the wake of a circular cylinder in the CORIA wind <sup>63</sup> tunnel, which is of recirculating type. The residual turbulence level is smaller than 0.2%. <sup>64</sup> The square test section is  $(0.4 \times 0.4 \text{m}^2)$  has a length of 2.5m. The mean pressure gradient <sup>65</sup> was set to zero. A circular cylinder of diameter d = 10mm, placed horizontally downstream <sup>66</sup> the contraction, spanned the width of the test section. The upstream velocity  $U_0$  was <sup>67</sup>  $6.5 \text{ms}^{-1}$ , corresponding to a Reynolds number  $Re_d$ , based on the cylinder diameter, of 4333. <sup>68</sup> Measurements of velocity fluctuations u and v were made at several streamwise locations <sup>69</sup> in the range  $40 \leq x/d \leq 110$ . The X-wire probe (Dantec 55P51) was calibrated using a <sup>70</sup> look-up table technique (e.g. Refs. 19 and 20) with velocity and angle increment of 1ms<sup>-1</sup> <sup>71</sup> and 5° respectively. The hot wires (Pt-W, 5µm diameter, 1.25mm length) were operated <sup>72</sup> with a Dantec constant temperature bridge at an overheat ratio of 1.6. Voltage signals were <sup>73</sup> passed though buck and amplifiers (SRS SIM983) and low-pass filtered (SRS SIM965) at a <sup>74</sup> frequency  $f_c$  close to the Kolmogorov frequency  $f_K \equiv \overline{U}/2\pi\eta$  ( $\overline{U}$  is the local mean velocity <sup>75</sup> and the Kolmogorov length-scale  $\eta = (\nu^3/\overline{\epsilon})^{1/4}$ ). The sampling frequency  $f_s = 20kHz$  and <sup>76</sup> the sampling period was about 4min. The air temperature in the wind tunnel was kept <sup>77</sup> constant during the calibration and measurements using a regulating cross-flow water-air <sup>78</sup> cooler system which allowed the mean temperature to be controlled to  $\pm 0.5^{\circ}C$ .

<sup>79</sup> Measurements made at the University of Newcastle in the wakes downstream of 5 different <sup>80</sup> two dimensional obstacles, each with the same characteristic dimension d = 25.4mm have <sup>81</sup> 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).

The data were obtained in an open circuit wind tunnel with a 2.4m long square test section  $(0.35 \times 0.35 \text{m}^2)$ . The five generators (a circular cylinder CC, a square cylinder SqC, a screen cylinder ScC, a solid plate P normal to the flow and a screen strip SS, Fig. 1) were so located 20cm downstream of the end of the contraction and spanned the width of the test section. All measurements were made on the centreline at x = 70d at several values of  $Re_d$ . A one-component vorticity probe (Fig. 1), which included an X-wire straddled by 2 parallel hot wires, was used to determine either  $\omega_z$  when the X-wire was in the plane (x,y) or  $\omega_y$ when the X-wire was in the (x,z) plane. All the hot wires were etched from Pt-10%Rh to a diameter of 2.5 $\mu$ m and an active length of approximatively 0.5mm (length to diameter ratio  $\approx 200$ ). They were operated at an overheat ratio of 1.5 with in-house constant temperature anemometers. The anemometer output signals were passed through buck and gain circuits, prior to low-pass filtering at a frequency  $f_c \approx f_K$ . The filtered signals were then digitized at a frequency  $f_s = 2f_c$  into a PC using a 12bit A-D converter. The sampling period was in the range 45 to 100s. The mean velocity  $U_0$ , upstream of the wake generator, was varied so that the Taylor microscale Reynolds number  $R_{\lambda_u} (\equiv \overline{u^{21/2}}\lambda_u/\nu$ , where  $\lambda_u = 15\nu\sqrt{\overline{u^2}/\overline{\epsilon}}$  is or the Taylor microscale) was in the range 150  $\leq R_{\lambda_u} \leq 300$  at x/d = 70.

#### 98 III. ONE-POINT ENERGY BUDGET

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

$$\frac{1}{2}\overline{U}_{i}\frac{\partial\overline{q^{2}}}{\partial x_{i}} + \overline{u_{i}u_{j}}\frac{\partial\overline{U_{i}}}{\partial x_{j}} + \frac{1}{2}\frac{\partial}{\partial x_{j}}\left(2\overline{u_{j}p} + \overline{u_{j}q^{2}}\right) + \overline{\epsilon} = 0.$$
(1)

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

Profiles along the y axis of  $\overline{vq^2}$  approximated by  $2\overline{vu^2} + \overline{v^3}$  indicated that the diffusion <sup>109</sup> of  $\overline{q^2}$  by v is negligible along the axis since  $\overline{vq^2} \approx$  constant on either side of the axis. A <sup>110</sup> similar conclusion was arrived at by examining the data supplied to us by Zhou<sup>22</sup>. The <sup>111</sup> latter measurements were made with a 3-component vorticity probe at x/d = 10, 20 and <sup>112</sup> 40. These data, with simultaneous information for u, v and w, indicated that  $\overline{vq^2}$  was more <sup>113</sup> 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 <sup>114</sup> superior to that with respect to x. The latter axisymmetry along the y axis for large-scale <sup>115</sup> statistics contrasts with Ref. 23 who observed a local axisymmetry with respect to the x axis <sup>116</sup> at the level of the smallest scales. However, the present result appears reasonable since the <sup>117</sup> influence of the Bénard-von Kármán street (a large-scale effect), which induces flow along <sup>118</sup> 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  $U_0^3/d$ .  $\Box -\frac{1}{2}\overline{U}\frac{\partial q^2}{\partial x}$ ,  $\circ \overline{\epsilon}$ ,  $\diamondsuit \left(\overline{u^2} - \overline{v^2}\right)\frac{\partial \overline{U}}{\partial x}$ .

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

$$\frac{1}{2}\overline{U}\frac{\partial\overline{q^2}}{\partial x} + \left(\overline{u^2} - \overline{v^2}\right)\frac{\partial\overline{U}}{\partial x} + \overline{\epsilon} = 0$$
<sup>(2)</sup>

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

$$\frac{1}{2}\overline{U}\frac{\partial\overline{q^2}}{\partial x} + \overline{\epsilon} = 0 \tag{3}$$

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

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

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

135 may be written as

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

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

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

138 where

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

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

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

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

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

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

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

<sup>143</sup> should be equal to 0.5. The expected linearity of  $\lambda^2$  with respect to x was used to estimate <sup>144</sup>  $x_0 ~(\approx 8d)$ . 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 <sup>145</sup> 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, <sup>146</sup> 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 ;  $\circ log_{10}(\overline{q^2}/U_0^2)$ ,  $\lhd log_{10}(\overline{u^2}/U_0^2)$ ,  $\Box log_{10}(\overline{\epsilon}d/U_0^3)$ ,  $\diamondsuit log_{10}(10^{-1}\lambda/d)$ ,  $\bigtriangleup log_{10}(U_d^2/U_0^2)$ , solid lines are linear fits to the data.

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

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

$$G_q + D_q + I_q = \frac{4}{3} \tag{11}$$

<sup>150</sup> 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 <sup>151</sup> 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.

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

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

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

We first consider the distributions of  $f_q$  (Fig. 5) for the five different wakes. The main 171 point which emerges from this figure is that the effect of initial conditions cannot be dis-172 counted, at least when  $\tilde{r}$  is sufficiently large. The collapse at small  $\tilde{r}$  is almost as good as 173 <sup>174</sup> 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 functions of  $r^*$  (the asterisk stands for the normalization by the Kolmogorov scales). However, although Kolmogorov scaling may be superior in the dissipative range, scaling on  $\lambda$ 176 and  $\overline{q^2}$  extends the collapse to a wider range of scales. In general, distributions for different 177 wakes depart from each other at  $\tilde{r} \approx 2$ . Note also the significant overshoot in  $f_q = \overline{(\Delta q)^2}/\overline{q^2}$ 178 <sup>179</sup> for the screen strip (SS) with a maximum at  $\tilde{r} = 12$ . Included in Fig. 5 are the distributions <sup>180</sup> (thin and bold curves) obtained from the models of Kurien & Sreenivasan<sup>26</sup> (hereafter KS) <sup>181</sup> 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,  $\Box$  SqC,  $\diamond$  ScC,  $\triangleright$  P, + SS.

182 second-order structure functions

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

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

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

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

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

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

195 numerator, *i.e.* 

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

<sup>196</sup> As c < 1,  $\overline{(\Delta q)^{*2}}|_{K} < \overline{(\Delta q)^{*2}}|_{A}$  for every separation r as observed in Fig. 5. Thus,  $I_{q}|_{K} < I_{q}|_{A}$ <sup>197</sup> and so  $G_{q}|_{K} > G_{q}|_{A}$ . Nevertheless, in the limit of large scales

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

198 and for small separations

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

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

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

 $G_q$  is much more affected by whether the wake generator is porous or solid than the other terms. For the porous wakes, the magnitude of  $G_q$  remains larger than that predicted by the wo models when  $\tilde{r} \geq 2$ . This reflects the relatively strong organization in these two wakes and hence the stronger influence of the initial conditions. Note that the A<sup>27</sup> model fails to adequately represent  $G_q$  in these two cases (for  $\tilde{r} \geq 2$ ) whereas it represents adequately 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 close to the expected value of 4/3 for the solid generators. There is however a significant these two wakes, the local anisotropy is greater than for the other 3 wakes<sup>21</sup>, reflecting the larger degree of organisation in the porous generator wakes than in the impervious generator 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.  $\circ, -\cdot -: G_q; \triangleright, \cdots : D_q; \diamondsuit, - - -: I_q; \Box, ---: (G_q + D_q + I_q)$ . The thick horizontal line corresponds to 4/3.

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

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

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

$$\overline{\epsilon} = C_{\epsilon c} \frac{\left(\overline{q^2}/3\right)^{3/2}}{L_c}.$$
(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  $A^{27}$  (thick lines) or  $KS^{26}$  (thin lines). (b) Evolution of the crossover length scale  $L_c$  normalized by  $\lambda$  with the Taylor microscale Reynolds number  $R_{\lambda}$ . Models; —  $A^{27}$  (thick lines), —  $KS^{26}$  (thin lines), —  $\cdot$  —  $L_s/\lambda_q = C_{\epsilon_c}R_{\lambda q}/15$  using  $C_{\epsilon c} = 0.34$ . Experiments  $\circ$  CC,  $\Box$  SqC,  $\diamond$  ScC,  $\triangleright$  P, + SS.

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

$$\overline{\epsilon} = C_{\epsilon} \frac{\left(\overline{u^2}\right)^{3/2}}{L_u},\tag{20}$$

where  $L_u$  is the longitudinal integral length-scale. Fig. 7(a) shows the evolution of the normalized dissipation rate parameter, defined with either Eq. (20) or (19). As was emphasized by Ref. 1,  $C_{\epsilon}$  reaches very different values depending on the type of obstacle. Indeed,  $C_{\epsilon}$ is equal to about 2 for CC wake and has a quite different value ( $\approx 0.6$ ) in the SS wake. On the other hand, when the normalization of  $\overline{\epsilon}$  is based on  $L_c$  and  $\overline{q^2}$ ,  $C_{\epsilon c}$  appears to be remarkably universal, *i.e.* it is independent of both the type of obstacle and the Reynolds number. The constancy of  $C_{\epsilon c}$  is also reasonably well reproduced by the two models. The measured value of  $C_{\epsilon c}$  is equal to 0.34 whereas the values predicted by the KS and A models are equal to 0.28 and 0.4 respectively.

It is evident that  $C_{\epsilon c}$  has a much stronger claim to universality than  $C_{\epsilon}$ . This is due in part to the fact that  $L_c$  is much more precisely determined than  $L_u$  but more importantly because  $L_c$  is an intermediate scale between that at which energy is injected into the flow and that that  $L_{40}$   $(r \sim \lambda)$  at which the transfer of energy is maximum. In other words, the cross-over between  $L_{41}$   $G_q$  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.

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

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

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

The magnitude of  $L_c$ , identified with the values of  $\tilde{r}$  for which  $G_q = I_q$  is shown in Fig. 253 <sup>254</sup> 7(b). Also included in this figure are the values of  $\tilde{L_c}$  which correspond to the two models.  $\tilde{L_c}$ 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 255 between the two models; although the range of  $R_{\lambda}$  is small, all the data are quite consistent 256 with a linear variation of  $L_c$  with  $R_{\lambda}$  and are very close to the line  $L_c/\lambda_q = C_{\epsilon_c}R_{\lambda_q}/15$ . 257 The interpretation for  $L_c$  provided here is similar to that for the shear length-scale  $L_s$ , 258 <sup>259</sup> first introduced by Corrsin<sup>31</sup> in the context of shear-dominated flows, and further physically interpreted in Refs. 11, 17, 18, and 32. Indeed, the authors emphasized that  $L_S \equiv \sqrt{S^3/\overline{\epsilon}^{31}}$ 260 corresponds to the crossover between contributions from the production term due to the 261 mean shear S and the nonlinear transfer term in the spherically averaged scale-by-scale 262 kinetic energy budget. Note also that the present authors have recently generalized the def-263 inition of  $L_S$  in order to account for the additional effect of the coherent strain using similar 264 arguments as in Ref. 11, *i.e.* based on scale-by-scale budgets (see Ref. 21). This suggests 265 that there exists a large variety of different relevant injection length-scales depending on 266  $_{267}$  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 <sup>269</sup> crossover between the production (due to both the mean shear and the coherent shear) and <sup>270</sup> the transfer term.

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

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

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



FIG. 8. Evolution of A<sub>q</sub> with R<sub>λ</sub> and the effect of m<sub>q</sub> on A<sub>q</sub>. Thick and thin curves correspond to the models of A<sup>27</sup> and KS<sup>26</sup> respectively. —: m<sub>q</sub> = 1.25, - - -: m<sub>q</sub> = 1.5, ·····: m<sub>q</sub> = 1.75.
O CC, □ SqC, ◊ ScC, ▷ P, + SS.

We now turn our attention to the maximum value of the energy transfer  $A_q \equiv \max(\overline{\Delta u(\Delta q)^2}/\overline{\epsilon}r)$ depicted in Fig. 8. Vertical and horizontal error bars are also given in order to appraise uncertainties in assessing  $A_q$ . Vertical error bars correspond to the standard deviation between different (isotropic) estimations of the dissipation rate. From the one-component vorticity probe oriented in either the x - y plane or x - z plane, one can calculate  $\overline{\epsilon}$  from <sup>285</sup> the following six expressions

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

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

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

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

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

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

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

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

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

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

#### 305 VI. CONCLUSIONS

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

<sup>306</sup> (i) Firstly, the X-wire measurements at various streamwise locations along the axis of <sup>309</sup> the wake indicate that equilibrium similarity, based on the Taylor microscale  $\lambda$  and the <sup>310</sup> turbulent energy  $\overline{q^2}$ , is verified satisfactorily. In particular, the collapse of the energy transfer <sup>311</sup> term  $\overline{\Delta \tilde{u} (\Delta \tilde{q})^2} R_{\lambda} q^{2^{-3/2}} \tilde{r}^{-1}$  appears to be of better quality than that previously obtained for <sup>312</sup> decaying turbulence downstream a grid<sup>9</sup>. In the intermediate wake and for large separations, <sup>313</sup> the one-point energy budget is closely approximated by a balance between advection and <sup>314</sup> energy dissipation rate. This is substantially different to what is generally observed in the <sup>315</sup> far field<sup>34</sup> in which the diffusion term has also to be taken into account. This difference <sup>316</sup> between the intermediate and far wake might be related to a reorganization or rescaling <sup>317</sup> of the large-scale structures<sup>4,35,36</sup>, sometimes referred to as the secondary vortex street. <sup>318</sup> Further investigations of velocity statistics and corresponding length-scales obtained over a <sup>319</sup> sufficiently large distance between the near and far fields are needed to confirm this.

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

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

(*iii*) Finally, the maximum value of the energy transfer term is systematically larger than that predicted by models for decaying HIT. This reflects the presence of the coherent <sup>336</sup> motion. The latter observation is important in the context of providing some insight into the <sup>337</sup> way the coherent motion interacts with small scales. Some progress has been made on this <sup>338</sup> interaction<sup>33</sup>, by quantifying the contribution of the coherent motion to the energy transfer <sup>339</sup> after writing the s-b-s budgets for both the coherent and random motions.

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