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Temperature fluctuations and scales in grid-generated turbulence

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To study the mixing of a passive scalar in nearly isotropic turbulence, experiments have been made in isotropically mixed thermal fields with thermal mesh size M_θ (*a*) equal to the momentum mesh size M, (*b*) larger than M (obtained by heating only alternate rods of the turbulence generating grid), and (*c*) smaller than M. This last condition was achieved by inserting a fine heating screen with $M_\theta < M$, at locations downstream of the turbulence grid. The heating screen was designed to produce negligible statistical change in the velocity field a short distance downstream. In all the heated grid experiments, for a given initial configuration of the thermal field, the intensity of temperature fluctuations θ normalized by the mean temperature rise ΔT, and the decay rate of $\overline{\theta^2}$ were both independent of the temperature of the grid. The principal effect of having $M_\theta > M$ was an increase in the relative intensity of temperature fluctuations compared with the $M_\theta = M$ case, and a marginal increase in their decay rate; contrary to expectation, the ratio R of temperature to velocity integral scales in the region of approximate homogeneity did not differ from that corresponding to $M_\theta = M$. In heated screen experiments, the relative decay rate was independent of M_θ/M and ΔT. For the three locations of the heating screen used in these experiments, the decay rate was also independent of the relative distance x_s of the heating screen from the turbulence generating grid; however, larger x_s was associated with larger relative intensity of fluctuations. To a first approximation, the ratio R approached unity according to the empirical relation $R = 1 - A \exp[-\alpha x_\theta/(UT_0)]$, where x_θ is downstream distance measured from the heating screen, and T_0 is a characteristic turbulence decay time scale at $x_\theta = 0$. It was also verified that the skewness of the streamwise temperature derivative is approximately zero sufficiently downstream of the heating screen. Where the present study overlaps with previous measurements, an extensive comparison reveals several points of agreement as well as departure.

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1. Introduction

A principal effect of turbulence is the mixing of scalar contaminants, most simply hastening the approach to molecular scale uniformity of an admixture (or contaminant) which (while being uniform in the mean) has an initially non-uniform concentration. In studies devoted to this aspect of turbulence, heat is most often the convenient paradigm, with temperature being its 'concentration'. Even in the simplest case of an isotropic scalar field embedded in a statistically prescribed isotropic turbulence field, the phenomenon is not fully understood. In particular, no accurate theoretical predictions are available.

Experiments approximating this simple mixing problem have been made by Kistler, O'Brien & Corrsin (1954, 1956), Mills *et al.* (1958), Gibson & Schwarz (1963), Yeh & Van Atta (1973), Sepri (1976), Newman, Warhaft & Lumley (1977) and Warhaft & Lumley (1978*a*, *b*). Gibson & Schwarz (1963) measured salinity fluctuations behind a regular grid discharging salt water into fresh-water flow; all other experiments were thermal, using a heated grid in a cooler airstream. In all these cases except some experiments of Newman *et al.* and Warhaft & Lumley, the turbulence generating grid was used also for creating the scalar field fluctuations, perhaps thus implying an implicit coupling of the initial conditions for the scalar and velocity fields.†

One of the analytical attempts at practical application of the existing information on turbulent mixing (Corrsin 1964) gave an explicit dependence of the relative rates of decay of concentration and velocity fluctuations on the ratio of the integral scales or microscales of the two random fields. It provided the motivation for the present study. A modified version of Corrsin's earlier theory (1951*b*) was given by Newman *et al.* (1977). Since most of the experiments cited above used grid configurations with scalar field mesh size equal to the momentum field mesh size, the theoretical predictions of the dependence of relative rates of decay on the integral scale ratio cannot be tested over a range of scale ratios.

One of the objectives of this work is to study this dependence, by creating thermal fields of mesh size larger as well as smaller than the momentum mesh size. First, the temperature field was produced in the conventional way, by heating uniformly all the turbulence generating rods of a square-mesh, bi-plane, round rod grid spanning a uniform air stream; this can serve as a reference for the other cases. Second, a temperature field with thermal mesh size twice that of the momentum field was generated by heating only the alternate rods of the turbulence generating grid. This procedure retains a possible shortcoming of the case with all rods heated (although perhaps to a lesser extent), namely a possible undesired initial statistical coupling of velocity and temperature fields.

In order to introduce a temperature fluctuation field with initial integral scale smaller than that of the turbulence, a square array of thin wires (having negligible effect on the velocity field) with mesh smaller than turbulence grid mesh, was inserted downstream of the turbulence-generating grid. In fact, two cases, with the heated wire mesh nearly equal and nearly half the mesh size of the turbulence generating grid,

† The configuration used by Kistler *et al.* had heated wires located just upstream of a wooden grid, but as the sheets of temperatures coming from the heated wires would be swept directly into the boundary layers of the grid rods, the initial statistical coupling of the two fields is effectively the same as in the conventional heated grid configuration.

were studied. Furthermore, it turned out that far enough downstream, the temperature fluctuation field, though initially pathologically pulse-like, quickly 'filled up' the fluid and decoupled statistically from the velocity field; for example, the correlation coefficient $\rho_{u\theta}$ was considerably smaller in magnitude than that behind the heated rod grids.

While the present study was under way, three papers reporting closely related experiments appeared (Newman *et al.* 1977, Warhaft & Lumley 1978*a*, *b*). In these papers, the authors specifically discuss anomalies which they found when the same grid was used to produce both velocity and temperature fluctuation fields. They found that the decay rate of temperature fluctuations varied over a wide range and depended on the heating power applied to the grid; the decay rate increased with the increase in heating. Increased decay rate could imply decreasing temperature length scale, but measurements show the opposite result, that the length scale increases with increasing heating power. Warhaft & Lumley suggest that heating the turbulence-generating grid somehow affects the 'geometry' of the temperature field, including its length scale, but no obvious mechanism can be thought of. By evaluating a buoyancy parameter far downstream of the grid, Warhaft & Lumley (1978*b*) suggest that heat was indeed a passive scalar in their experiments, but it is not clear whether the relevant buoyancy parameter should be evaluated in the shear layers immediately downstream of the heating rods. In fact, it is not clear what the most appropriate buoyancy parameter should be.

It should be remarked that their component turbulent energies (their figure 3) show relative histories appreciably different from those found by most other investigators. In particular, their transverse fluctuations $(\overline{v^2})$ decayed more slowly than the streamwise fluctuations $(\overline{u^2})$, actually becoming larger beyond $x/M = 50$. This is rather different from the usual situation (Comte-Bellot & Corrsin 1966), but the authors make no comment about it.

A third anomalous behaviour found by Warhaft & Lumley (1978*a*, *b*) was the appearance of a sizeable region with $-\frac{5}{3}$ slope in the spectrum of the temperature fluctuations, even at relatively low Reynolds and Péclet numbers. They also found a relatively high coherence between temperature and longitudinal velocity fluctuations, as well as a persisting non-negligible correlation between them (the correlation coefficient being about -0.3), and suggest that the coupled initial conditions might persist throughout the life-time of the fluctuations. To remedy this, Newman *et al.* and Warhaft & Lumley placed an array of parallel, uniformly heated wires† downstream of the unheated turbulence-generating grid, and studied the subsequent decay of the resulting temperature fluctuations. They found that the anomalous $-\frac{5}{3}$ region, as well as the strong correlation and coherence between the temperature and the velocity fields, disappeared. In these experiments, however, the relation between the

† We may here note that, at their tunnel speed, the wire Reynolds number based on the kinematic viscosity at the ambient temperature was about 140. Clearly, because of the increased kinematic viscosity due to heating, the effective Reynolds number will be substantially smaller, but the precise value of the temperature to be used here is not clear. Our experience is that the use of kinematic viscosity at the average of the wire surface temperature and the ambient temperature results in too low a value for the effective Reynolds number. Considering the possible operating temperature of the Chromel-A wires used by Warhaft & Lumley, it seems likely that there was some vortex-shedding behind these wires.

temperature decay rate and the amount of heating cannot be determined with any certainty because more than one parameter was changed between one experiment and another (see also § 4).

Our interest here is in isotropic turbulent mixing of a passive scalar contaminant. At small enough temperature differences, heat should be passive. A concern will be to examine this aspect of the problem as well as to pursue the central goal of studying mixing.

2. Experimental conditions and instrumentation

The experiments were carried out in an open return wind tunnel with a $30{\cdot}5 \times 45{\cdot}7$ cm rectangular test section. Turbulence was produced by a bi-planar grid consisting of a square mesh of round heating rods (Chromalox TSSM 14XX) with diameter 0·635 cm, positioned 2·54 cm between centres. The momentum mesh M was therefore 2·54 cm, and the grid solidity was 0·44. The mean speed U of the tunnel was 4·4 m s^{-1}, so the grid Reynolds number was

$$R_M \equiv UM/\nu \simeq 7420.$$

During the first part of the study, all or alternate rods were heated electrically to produce temperature fluctuation fields generated by thermal mesh M_θ equal to or twice M. In some cases, only horizontal or alternate horizontal rods were heated.

For the second part of the study, a screen with square mesh of 1·11 cm was constructed using Nichrome V wires of diameter 0·127 mm. The screen solidity was therefore 0·023, while the wire Reynolds number, using the kinematic viscosity at the ambient temperature, was about 38, just lower than the critical value of 40 for the incidence of vortex shedding (e.g. Kovasznay 1949). The latter was verified by placing hot wires directly behind the heating wires; regular sinusoidal oscillations which appeared at higher tunnel speeds did not occur at the test speed of 4·4 m s^{-1}.

The heating wires were divided in two alternating groups connected to the electric power separately; this way initial conditions with thermal meshes M_θ equal to either 1·11 cm or 2·22 cm could be produced. The screen was annealed for several hours, and then the wires were retightened to minimize possible sagging or distortion. The screen could be positioned at distances x_s/M equal to 20, 34, and 54 downstream of the turbulence generating grid.

For some measurements, a heating wire of diameter 0·16 mm was used instead of the 0·127 mm diameter wire. The thicker wire was also operating below the critical Reynolds number when heated to an estimated temperature of about 100 °C.

The streamwise velocity fluctuation was measured with a tungsten hot wire of diameter 5 μm and length 1·2 mm, powered by a DISA 55D01 constant temperature anemometer circuit, while the fluctuating temperature was in most cases measured with a platinum cold-wire of diameter 1 μm and length 0·4 mm, supplied with a constant current of 0·30 mA by a 'home-made' electronic circuit (Tavoularis 1978*a*); a home-made cold wire (platinum–rhodium) of 0·6 μm diameter and 0·8 mm length was also used for some measurements. The mean temperature of the flow T as well as the upstream 'ambient' temperature T_a were measured with two Fenwal Electronics thermistor probes. On-line data processing was performed with a PDP 11/40 digital mini-computer. All measured quantities were corrected for noise contamination.

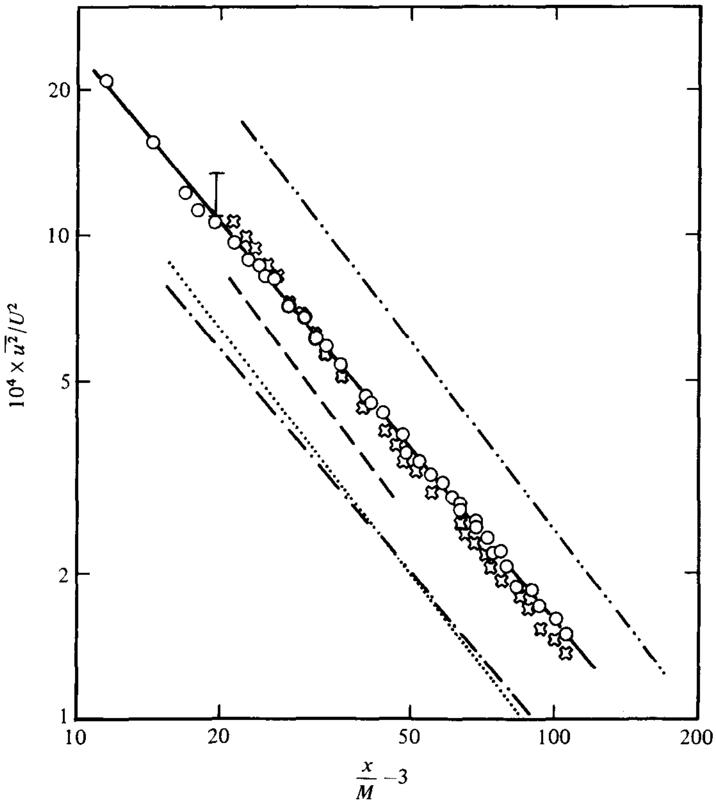

FIGURE 1. Downstream decay of the mean-square velocity fluctuation $\overline{u^2}$. —·—, Mills *et al.*; ·····, Kistler *et al.*; ——, Yeh & Van Atta and Sepri; —··—, Warhaft & Lumley. Present measurements: ○, without the heating screen; ☆, with the heating screen located about 20 meshes downstream of the turbulence grid. —— indicates $0{\cdot}04\,(x/M-3)^{-1{\cdot}20}$. Vertical bar indicates variability in $\overline{u^2}/U^2$ a mesh-length downstream of the heating screen, depending on the position of the hot-wire. $M_\theta/M = 1$ unless otherwise indicated.

3. The measurements

3.1. *The decay of velocity fluctuations*

The downstream decay of the mean-squared streamwise turbulent velocity behind the unheated grid is shown in figure 1. A relation of the type

$$\frac{\overline{u^2}}{U^2} = \alpha_1 \left(\frac{x}{M} - \frac{x_0}{M}\right)^{-n} \tag{1}$$

was fitted to the data, with α_1, x_0 and n chosen to give the maximum possible straight line fit to the data in logarithmic co-ordinates. With $\alpha_1 = 0{\cdot}04$, $x_0/M = 3$, and $n = 1{\cdot}2$, (1) approximates the data well over the entire range of measurements. Equation (1) with these empirical values agrees well with the decay law derived by Saffman (1967), as well as the data reported by other workers, e.g. Comte-Bellot & Corrsin (1966). The corresponding kinetic energy decay data of Kistler *et al.*, Mills *et al.*, Yeh & Van Atta and Warhaft & Lumley are also plotted in figure 1. The decay data of Sepri (1976) are indistinguishable from those of Yeh & Van Atta.

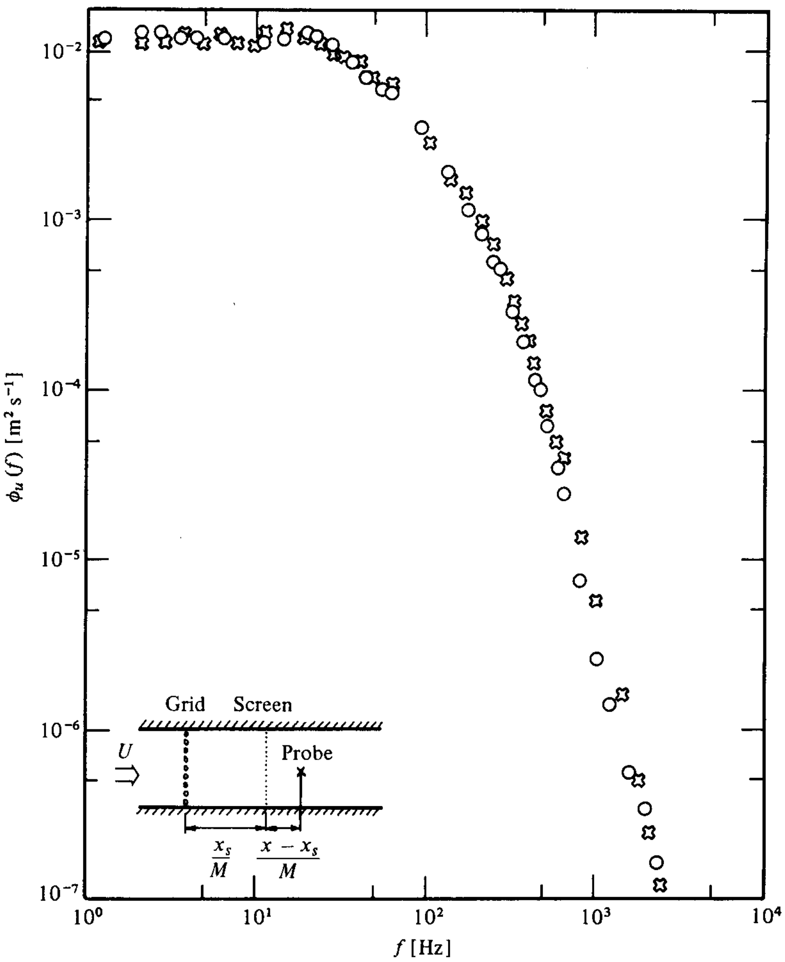

FIGURE 2. Comparison of the power spectral density of the velocity fluctuation u at $x/M \simeq 27$, without the heating screen (○), and with the heating screen (✠) at $x_s/M = 20$.

It was verified experimentally that the density changes caused by heating the grid to give mean airflow overheat $\Delta T \leqslant 5$ °C were too small to cause a perceptible change in the decay data. On the other hand, the effect of the fine-wire screen on the velocity field was not completely negligible and was studied in some detail. The downstream decay of $\overline{u^2}$ behind the unheated wire screen positioned at $x_s/M = 20$ is shown in figure 1. For $(x-x_s)/M \leqslant 10$, $\overline{u^2}$ is somewhat larger than when the wire screen is absent, although it was ensured that no 'von Kármán–Bénard' vortices were shed behind the wires. This increased intensity is due to the straining action on the pre-existing turbulence by the mean velocity field set-up by the array of wires forming the screen. Sufficiently far downstream of the screen (in practice for $(x-x_s)/M \geqslant 20$ in this case), the net damping effect of the screen (see, for example, Dryden & Schubauer 1947; Townsend 1951; Batchelor 1953) may predominate resulting in a reduced turbulence intensity. Kellogg (1965) and Kellogg & Corrsin (1979) have reported in detail the action of a low-solidity, parallel-wire screen on pre-existing grid turbulence.

Typical one-dimensional velocity spectra with and without the wire screen are shown in figure 2. Except for a slight difference in the total energy, no appreciable influence of the screen on the velocity spectrum can be observed. The Kellogg Corrsin array of parallel wires which was of 5 % solidity produced, on the other hand, a substantial, deliberate distortion of the one-dimensional spectral data.

Finally, it was verified that the streamwise velocity integral scale L_f and the Taylor microscale λ_g did not show any change beyond the experimental scatter, with or without the screen. Data will be shown later, in § 3.4. Further, the Taylor microscale λ_g, evaluated as

$$\lambda_g = [\overline{u^2}/\overline{(\partial u/\partial x_1)^2}]^{\frac{1}{2}}, \tag{2}$$

agreed well with that implied by the decay law, equation (1), namely,

$$\lambda_g = [-10\nu\overline{u^2}/(d\overline{u^2}/dt)]^{\frac{1}{2}}. \tag{3}$$

Taylor's approximation was used to replace streamwise derivative with temporal derivative in (2), and vice versa in (3).

3.2. *Isotropy and homogeneity of velocity and temperature fields*

The mean velocity U and the mean-squared turbulent velocity $\overline{u^2}$ profiles become uniform (to within ± 2 %) for $x/M \gtrsim 40$. At about the same distance, the mean temperature rise ΔT also becomes nearly uniform (to within 2 % with all rods heated and 5 % with alternate rods heated). Comparable uniformity was observed for the mean-square temperature fluctuation $\overline{\theta^2}$.

For the case of the heated downstream screen, the results of Corrsin & Karweit (1972) on the angular dispersion of fluid line elements in isotropic turbulence can be used to obtain a rough estimate of the distance required for the development of isotropy of the temperature field. This suggests that between 25 and 40 velocity mesh lengths are required, the lower value corresponding to the case when the heated screen is located at $x_s/M = 20$ and the higher value to $x_s/M = 54$. In the present experiment, clearly translational dispersion is also relevant. The actual measurements suggest that a somewhat smaller distance of about 15 to 25 velocity mesh lengths is sufficient to attain reasonable homogeneity and isotropy.

The near-isotropy of the temperature field was verified in all cases by a comparison of the measured correlation functions $R_\theta(r_i)$, $i = 1, 2, 3$, along the three co-ordinate axes. Within the measurement accuracy, the three curves were found to be identical, as required by isotropy.

3.3. *The decay of temperature fluctuations*

The downstream decay of $\overline{\theta^2}$ behind the heated grid is shown in figure 3 for several heating modes and several heating powers. The various heating modes correspond to the cases with:

(*a*) all rods heated;
(*b*) all horizontal rods heated;
(*c*) all alternate rods heated;
(*d*) only alternate horizontal rods heated.

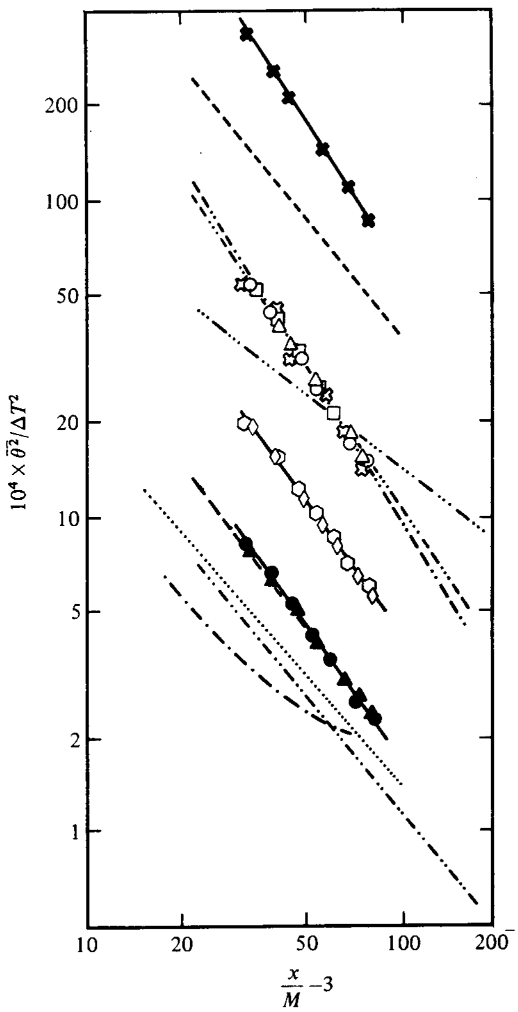

FIGURE 3. Downstream decay of the mean-square temperature fluctuation in the heated grid experiments. The symbols are explained in table 1.

In each case, relations of the type

$$\frac{\overline{\theta^2}}{(\Delta T)^2} = \beta\left(\frac{x}{M} - \frac{x_0'}{M}\right)^{-m} \tag{4}$$

were tried with several choices of x_0'/M with the aim of covering the largest span of experimental data by the relation (4). Also shown in the figure are the temperature decay data from all other previous experiments reported in the literature for heated grids.

Concentrating on the present data, we note first that in all cases, the temperature decay power law has effectively the same origin as the velocity decay law (1), i.e. $x_0' \simeq x_0 \simeq 3M$. Secondly, the constants m and β are essentially the same for a given heating mode and are independent of ΔT (for $\Delta T \leqslant 5$ °C), as should be expected if the temperature fluctuations are indeed passive, and the rod boundary layer heating

Authors	Symbol	M_θ/M	ΔT (°C)	β	m	m/n
Mills *et al.*	—·—	1	5·0	not applicable		
Kistler *et al.*	·····	1	5·0	0·030	1·17	0·92
Yeh & Van Atta; Sepri	———	1	10·0	0·096	1·39	1·06
Warhaft & Lumley	—···—	1	1·66	0·057	0·81	0·63
	—··—	1	9·09	0·036	1·26	0·98
	-----	2	0·50	1·210	1·26	0·98
	---··--	2	2·00	1·150	1·52	1·18
	--·--	2	3·30	1·970	1·66	1·29
Present experiments	●	1	2·1	0·124†	1·44†	1·20†
	▲	1	4·4			
	◇	1‡	1·4	0·263†	1·40†	1·17†
	O	1‡	2·1			
	△	2	0·75	1·191†	1·54†	1·28†
	⌘	2	1·75			
	○	2	2·00			
	□	2	2·70			
	✖	2‡	0·50	7·70	1·55	1·29

† These values are the average for the set.
‡ In these cases, only horizontal rods were heated. M_θ here refers to the distance between parallel rods.

TABLE 1. Details of constants of temperature decay downstream of heated grids, relevant to figure 3.

is not dynamically important. Thirdly, m (to some extent) and β differ for different heating modes. When all the horizontal rods only are heated, i.e. case (b), the fluctuation level (i.e. β) is higher (by a factor of 2·5) than when all the rods of the grid are heated, i.e. case (a); the index m however is not significantly different in the two cases. When all the alternate rods are heated, i.e. case (c) [$M_\theta = 2M$], the fluctuation level is even higher (by a factor of 6 to 7 compared with case (a)); the index m is also slightly higher (by about 7 %). In case (d), when only alternate horizontal rods are heated, β is much higher, but m is about the same as in case (c). These details are set out in table 1, and can be summarized as follows:

(i) The constant x_0' is approximately the same as x_0.

(ii) The constants m and β are effectively independent of ΔT for a given heating configuration.

(iii) For a given thermal mesh size M_θ (whether a square mesh or formed by an array of horizontal heating rods), m is essentially the same. However, it increases slightly, but definitely, with doubling the spacing of the heated rods relative to the turbulent mesh size M.

(iv) The value of the constant β increases with increasing spacing of the heated rods. This of course is to be expected because, for a given ΔT, smaller M_θ corresponds to smaller initial mean temperature gradients, hence smaller production rate of $\overline{\theta^2}$ close to the grid.

For consistency of comparison, decay data from other authors are plotted in figure 3 also with an assumed value of $x_0/M = 3$. In some cases, this improves the fit of equation (4) to the data, and in some cases slightly increases the scatter, but in general

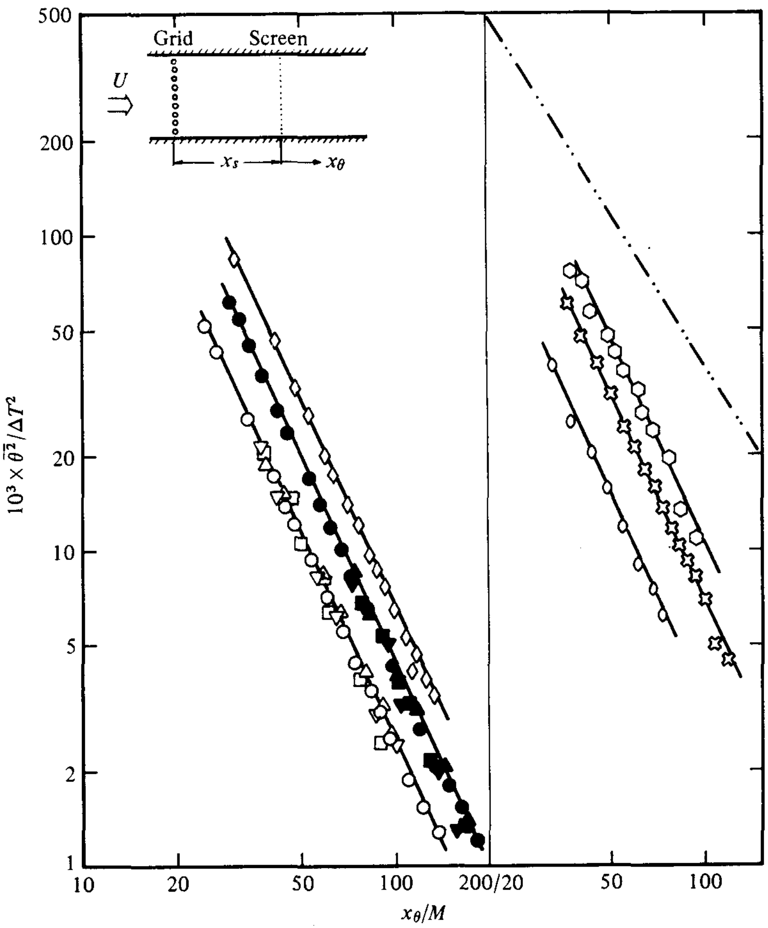

FIGURE 4. Downstream decay of the mean-square temperature fluctuation in the heated screen experiments. The symbols are explained in table 2.

equation (4) holds for virtually all the data. The one unexplained exception is the data of Mills *et al.* (1958), for which a simple power law of the form of equation (4) cannot be fitted. Further, it seems reasonable to normalize the root-mean-square temperature by the mean temperature rise ΔT (instead of the absolute temperature T used by Warhaft & Lumley). With this normalization, all the decay data for a given M_θ/M plot in a certain vicinity of each other, except one set of the Warhaft & Lumley data with $\Delta T \simeq 1{\cdot}66$ °C. This set also has a power law exponent ($m = 0{\cdot}81$) that is vastly different from all other data. Leaving aside this set of data for a moment, and considering only cases with $M_\theta/M = 1$, the average value of m is 1·31, with $\pm 10\,\%$ variation. This variation is quite comparable to that of the index n for the velocity decay data. Furthermore, the collected data do not suggest any definite relation between m and ΔT (or the heating power). (If, instead, we take into account the Warhaft & Lumley set as well, the mean value of m will be about 1·25, with variations of the order $+15\,\%$ and $-35\,\%$ about the mean).

Authors	Symbol	M_θ/M	ΔT (°C)	x_s/M	β_s	m_s	m_s/n
Warhaft & Lumley†	—··—	1	0·14	20	275	2·10	1·63
Present experiments	○	0·44	1·00	20	68	2·23	1·86
	□		0·75				
	△		0·61				
	▽		0·50				
	●		1·00	34	120	2·28	1·86
	■		0·75				
	▲		0·61				
	▼		0·50				
	◇		0·95	54	184	2·24	1·86
	⬯	0·87	0·50	20	89	2·20	1·86
	✕		0·50	34	180	2·23	1·86
	⬡		0·50	54	282	2·20	1·86

† In these experiments, the heated screen had only an array of horizontal wires. M_θ refers to the distance between a pair of wires.

TABLE 2. Details of constants of temperature decay downstream of heated screen, relevant to figure 4

Figure 4 shows the new decay data for mean-squared temperature fluctuation downstream of the thermal screen for $x_s/M = 20$, 34, and 54, and $M_\theta/M = 0{\cdot}44$ and 0·87. Also shown is a set of Warhaft & Lumley's '*mandoline*' data corresponding to roughly equal M_θ/M (1·0 compared with the present 0·87), comparable x_s/M. The large difference in $\overline{\theta^2}/(\Delta T)^2$ (we have estimated their ΔT from heating current and voltage) probably reflects the difference in heating screen geometries. Consistent with the contrast between our horizontal-rods-only-heated case and our all-rods-heated case (figure 3), the parallel wire screen gives larger $\overline{\theta^2}/(\Delta T)^2$ than does the square-mesh screen. The other three cases of Warhaft & Lumley's data (their figure 14) correspond to $M_\theta/M = 2$, hence cannot be compared here. To a good approximation, a relation of the type

$$\frac{\overline{\theta^2}}{(\Delta T)^2} = \beta_s \left(\frac{x - x_s}{M_\theta}\right)^{-m_s} \tag{5}$$

fits the data well, where m_s is a constant ($\simeq 2{\cdot}2$) and is essentially independent of x_s/M and M_θ/M. To a good approximation, the location of the heating screen was also found to be the effective origin. As might be expected, the coefficient β_s depends on x_s/M and M_θ/M, as shown in table 2. As before, larger M_θ/M is associated with larger β_s. A qualitative explanation for the increasing value of β_s with increasing x_s/M is that at larger x_s, the turbulent fluctuations responsible for the mixing process are of a larger scale and lower intensity.

An important feature of these measurements is the universality of the decay coefficient m_s in equation (5) for different conditions. A possible reason for this empirically observed universality is that, at the point of introduction of temperature fluctuations, the ratio L_f/M_θ varies by a factor of only about 3 among the various experiments. Figure 4 and table 2 show that m_s is again practically independent of ΔT.

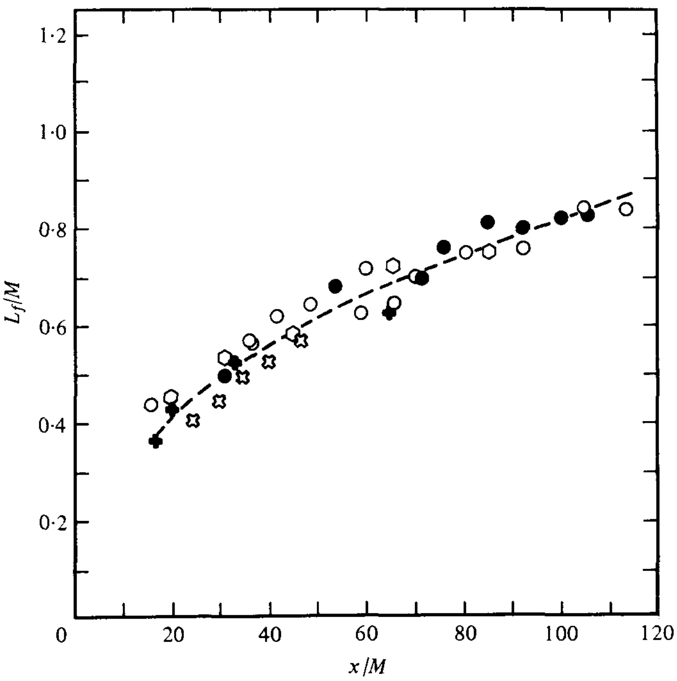

FIGURE 5. Downstream development of the longitudinal integral scale L_f of the velocity fluctuation u. ○, Kistler *et al.*; ✚, Mills *et al.*; ✿, Yeh & Van Atta and Sepri. Present measurements: ○, without the heating screen; ●, with the heating screen at $x_s/M = 20$. ——, $0{\cdot}13\,(x/M-3)^{0{\cdot}4}$.

3.4. *The downstream development of the integral length scales*

The streamwise integral length scale L_f of the streamwise velocity fluctuation is shown in figure 5. Here and elsewhere in this paper, the integral length scale of a signal is obtained by evaluating the integral of its auto-correlation function up to its first zero and then using Taylor's 'frozen-field' approximation. The rationale of this former procedure has been discussed by Comte-Bellot & Corrsin (1971). In view of the considerable scatter in the length scale data of figure 5, it was decided for later use to fit the data with a relation consistent with the self-similarity hypothesis (von Kármán & Howarth 1938), and the decay law (1). The resulting best expression is

$$L_f/M = 0{\cdot}13\,(x/M-3)^{0{\cdot}4}, \tag{6}$$

shown as a dashed line in figure 5. This expression provides a reasonably good fit to our data as well as the integral scale data of Kistler *et al.*, Mills *et al.* and Yeh & Van Atta, also included in figure 5. (Even though the decay data of Mills *et al.* seem unusual, we see no obvious problems with their integral scale data.) The figure further shows that there is very little difference in integral scale values whether or not the heating screen is present.

Figure 6 shows the integral scale ratio $R \equiv L_\theta/L_f$ for all the heated grid experiments. Here, L_θ is the integral scale of temperature fluctuation, and L_f is approximated by (6). In all these experiments, this ratio is a constant of about 0·8–0·85, essentially independent of x. Probably the most interesting result is the fact that the integral scale ratio L_θ/L_f is hardly affected by doubling the ratio M_θ/M, although in principle

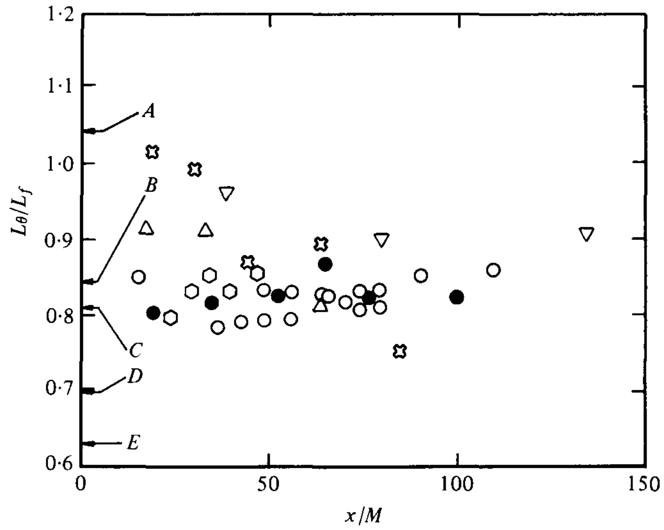

FIGURE 6. The integral scale ratio $R = L_\theta/L_f$ in the heated grid experiments, with L_f approximated by equation (6). ⬡, Kistler *et al.*; △, Mills *et al.*; ✡, Yeh & Van Atta; ▽, Warhaft & Lumley, $M_\theta/M = 1$, estimated from their one-dimensional temperature spectra. Present measurements: ●, $M_\theta/M = 1$; ○, $M_\theta/M = 2$. Arrows refer to R computed according to Corrsin's (1964) theory from the ratio of the measured decay rates of temperature and velocity fluctuations; A, for Kistler *et al.*'s data; B, for Yeh & Van Atta's data; C, for Warhaft & Lumley's data; D, for present data with $M_\theta/M = 1$; E, for present data with $M_\theta/M = 2$. Equation (9) of Newman *et al.* (1977) gives almost exactly twice the value for R in each case.

there is no reason to expect this result if $M_\theta \neq M$, except as an asymptotic result when the 'mixing action' has persisted for several integral time scales L_f/u'. It seems that the velocity field imposes its own scale on the temperature fluctuation field, irrespective of the thermal mesh size, when turbulence and temperature fluctuations originate at the same grid.

Corrsin (1964) obtained a theoretical estimate of the relative rates of decay of scalar and velocity fluctuations, under assumptions appropriate to high Reynolds and Péclet numbers. If θ' and u' are the root-mean-square scalar and velocity fluctuation intensities,

$$\frac{d\theta'/\theta'}{du'/u'} \simeq \tfrac{1}{2}(3-\sigma^2)\left(\tfrac{2}{3}L_f/L_\theta\right)^{\frac{2}{3}}, \tag{7}$$

where σ is the fluid Prandtl number. A possible criticism of (7) is that it postulates that a large scale property such as the length scale ratio depends explicitly on the molecular Prandtl number. Using simple modifications of an earlier analysis due to Corrsin (1951*b*), Newman *et al.* (1977) obtained an alternative expression

$$\frac{d\theta'/\theta'}{du'/u'} = \frac{2}{B}\left(\frac{1}{k_0 L_\theta}\right)^{\frac{2}{3}}, \tag{8}$$

where k_0 is the wavenumber corresponding to the peak of the velocity spectrum (= to the cut-off wavenumber in Corrsin's high-Reynolds-number approximation), and B

is the constant in the three-dimensional temperature spectrum in the inertial-convective subrange (Oboukhov 1949; Corrsin 1951*b*)

$$E_\theta(k) = B\epsilon^{-\frac{1}{3}}\epsilon_\theta k^{-\frac{5}{3}}.$$

Here, ϵ and ϵ_θ are the rates of velocity and temperature dissipation functions, respectively. Using Corrsin's (1964) estimate that

$$k_0 = \frac{3\pi}{10} L_f^{-1} \approx L_f^{-1},$$

equation (8) can be rewritten as

$$\frac{d\theta'/\theta'}{du'/u'} = \frac{2}{B}\left(\frac{L_f}{L_\theta}\right)^{\frac{2}{3}}. \tag{9}$$

This equation, unlike Corrsin's expression (7), contains the constant B explicitly. Newman *et al.* (1977) use a value of 1·32 for B, but we may note the wide variability (between 1·20 and 3·8, see for example Monin & Yaglom 1975, p. 504) among the various estimates of this constant. Equations (7) and (9) are identical if $B \simeq 2{\cdot}1$.

We can now evaluate the left-hand side of equations (7) and (9) from the decay data of velocity and temperature fluctuations, and can treat these equations as relations predicting the ratio L_θ/L_f. In the heated grid experiments, since both $\overline{u^2}$ and $\overline{\theta^2}$ decay according to power laws with effectively the same origin, the left hand side of equations (7) and (9) is constant for a given experiment, and so is the ratio L_θ/L_f. The values of L_θ/L_f evaluated from the observed decay rates according to equation (7) for each of the experiments are shown at the ordinate of figure 6. The agreement appears to be in order of magnitude only. It is clear that the length scale ratio given by (9) depends on the precise value assigned to B. Choosing $B = 1{\cdot}32$ (as preferred by Newman *et al.* 1977) gives almost exactly twice the corresponding values given by (7).

Perhaps the considerable discrepancy between experimental and theoretical values of the integral scale ratio R is due in part to the fact that the experimental $R_\lambda \equiv u'\lambda_g/\nu$ is fairly small, $R_\lambda \simeq 30$. In the other limit of extremely small Reynolds and Péclet numbers, Corrsin (1951*a*) showed that

$$\frac{L_\theta}{L_f} = \frac{1}{\sqrt{\sigma}} \simeq 1{\cdot}18 \quad \text{(for air)}.$$

This is higher than the experimental values of L_θ/L_f (see figure 6*a*). Furthermore, this expression predicts no dependence of L_θ/L_f on the relative rates of decay of u and θ; experiments, on the other hand, do show some dependence.

In the heated screen experiments, measurements show that R varies with x and is substantially smaller than unity for both $M_\theta/M = 0{\cdot}44$ and $0{\cdot}87$. Because x_s/M appears as an additional parameter in the heated screen experiments, it is not convenient to plot the data against x_θ/M (where $x_\theta = x - x_s$), but against the 'time of flight' x_θ/U relative to the heating screen, normalized by an integral time scale $T_0 \equiv L_f/u'$ at $x_\theta = 0$. Figure 7 shows that the ratio R exhibits approximate similarity in this co-ordinate system; the data cover three values of x_s/M and two values of M_θ/M. Because of turbulence mixing, although in principle we expect R to approach approximately unity eventually, data presented here show that the persistence of

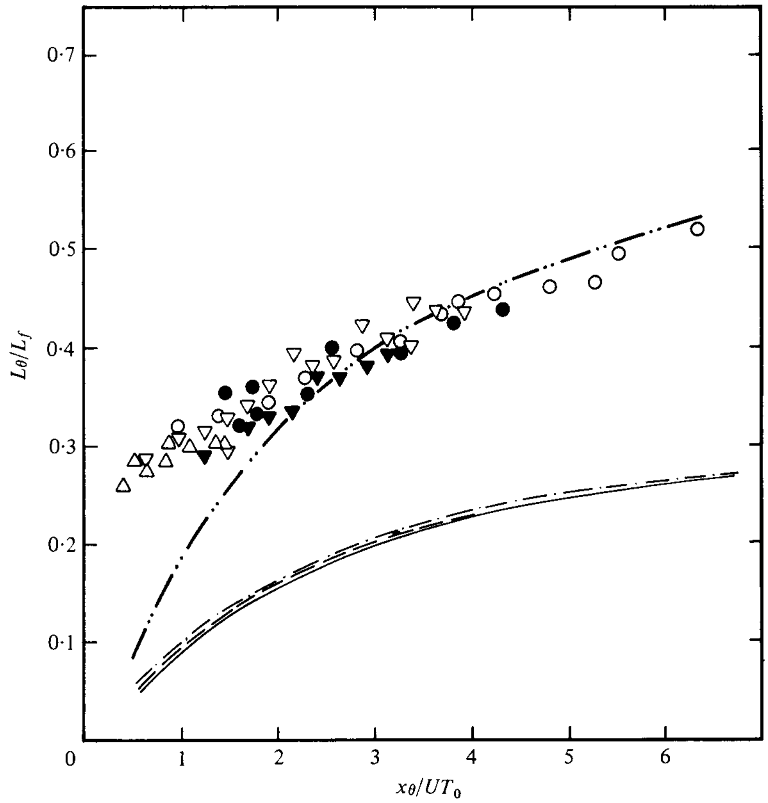

FIGURE 7. Downstream development of the integral scale ratio in the heated screen experiments, with L_f approximated by equation (6), expressed in the similarity co-ordinate x_θ/UT_0. Open symbols correspond to $M_\theta/M = 0{\cdot}44$, and filled symbols to $M_\theta/M = 0{\cdot}87$. Location of heating screen: ○, $x_s/M = 20$; ▽, $x_s/M = 34$; △, $x_s/M = 54$. Bottom curves correspond to Corrsin's (1964) theory: ——, $x_s/M = 20$; — —, $x_s/M = 34$; —·—, $x_s/M = 54$. —··—, from the theory of Newman *et al.* (1977), equation (9), $x_s/M = 20$. The curves for $x_s/M = 34$ and 54 are nearly identical with that for $x_s/M = 20$.

turbulent mixing for 5 to 6 integral time scales is not enough to attain this asymptotic result. In ordinary wind-tunnel practice therefore, R is unlikely to become unity even if sufficient tunnel length is available; the reason is that the approximate self-preservation state is unlikely to hold far downstream, where conditions closer to the final period of decay are likely to set in.

The observation that the length scale ratio quickly attains its asymptotic value when $M_\theta > M$, but not when $M_\theta < M$, can be given a *qualitative* explanation that is independent of the manner in which the temperature field is produced. When $M_\theta \gg M$, the mixing effect is mainly one of wrinkling of the temperature phase-lines by a quasi-diffusion effect due to the velocity field, and has a characteristic time scale τ_1 given by (Comte-Bellot & Corrsin 1971)

$$\tau_1 \sim \frac{L_\theta^2}{\varepsilon} \sim \frac{L_\theta^2}{L_f u'} = \left(\frac{L_\theta}{L_f}\right)^2 \frac{L_f}{u'},$$

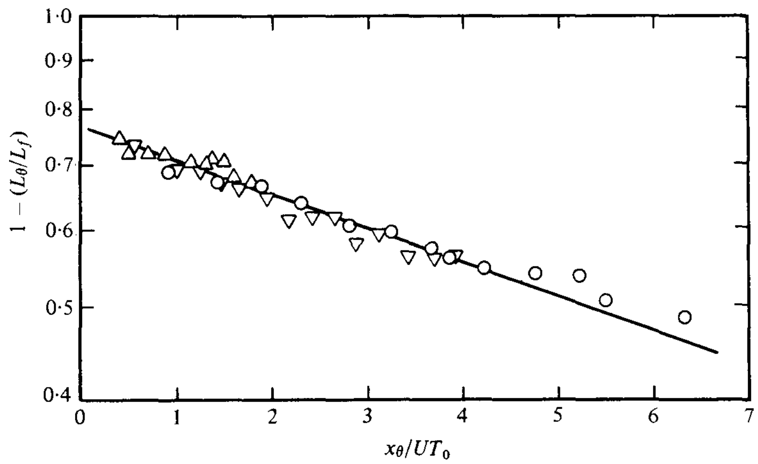

FIGURE 8. A replot of the data of figure 7 for $M_\theta/M = 0{\cdot}44$. ——, $0{\cdot}76\exp[-0{\cdot}81x_\theta/UT_0]$.

where $\varepsilon \sim L_f u'$ is the eddy-diffusivity due to the velocity field. On the other hand, when $M_\theta \ll M$, the mixing is essentially a straining effect of the velocity field, and the characteristic time scale τ_2 is given by

$$\tau_2 \sim \frac{L_f}{u'}.$$

For comparable L_f/u', the ratio

$$\tau_2/\tau_1 \sim (L_f/L_\theta)^2 > 1$$

in view of the experimental result that $L_f > L_\theta$ (even when $M_\theta/M = 2$).

Also shown in figure 7 is the ratio R as given by Corrsin's (1964) estimate, by expressing, for each x_s/M, the left-hand side of equation (7) in terms of x_θ (with the help of equations (1) and (4)), and normalizing it by the appropriate measured value of UT_0. Although the theoretical estimate is significantly smaller than the measured L_θ/L_f, it is consistent with the similarity suggested by the experimental data. The corresponding values from (9) (with $B = 1{\cdot}32$) are also plotted in figure 7, and are seen to be in much better agreement with the measurement, especially for $x_\theta U/T_0 \gtrsim 2$. So, in general, we conclude that the Newman *et al.*'s expression with $B = 1{\cdot}32$ performs better in relation to the screen data, but rather poorly with the grid data. (As already remarked, with $B = 2{\cdot}1$, its performance is indistinguishable from that of the equation (7).)

A replot in figure 8 of the integral scale data shows that the approach to unity of the ratio R is essentially exponential, and can be written as

$$R = 1 - A\exp[-\alpha(x_\theta/UT_0)],$$

where the constants A and α are, respectively, 0·76 and 0·081. Again, the small value of α suggests that it takes many integral time scales for R to become approximately unity. This relation can be expected to breakdown for large values of x_θ/UT_0, as indeed the data seem to suggest.

Figure 9 shows that the downstream development of the ratio λ_θ/λ_f in the co-ordinate x_θ/UT_0 also exhibits a similarity for the heated screen experiments. Here,

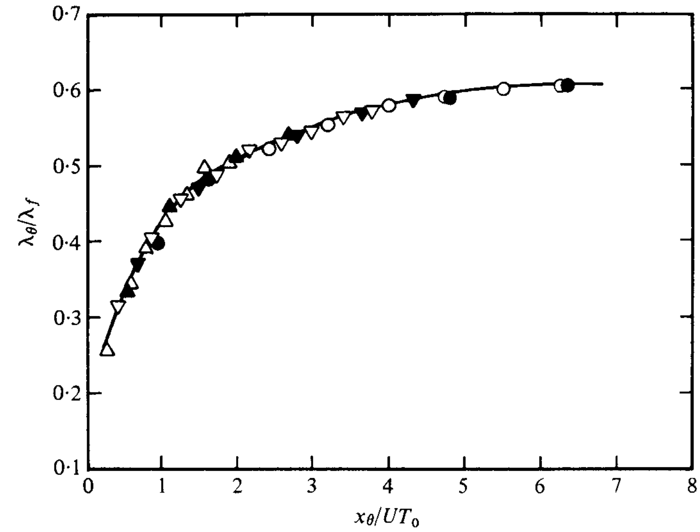

FIGURE 9. Downstream development of the ratio of the temperature microscale to the longitudinal microscale of the velocity fluctuation, u, in the heated screen experiments. Symbols as in figure 7.

λ_f was obtained from the decay law (1) using the relation (3), and the isotropic result $\lambda_f = \sqrt{2}\lambda_g$; and λ_θ was obtained from the decay law (4) using the relation (Corrsin 1951*b*)

$$\lambda_\theta = (-12\gamma\overline{\theta^2}/(d\overline{\theta^2}/dt))^{\frac{1}{2}},$$

where γ is the thermal diffusivity in air. Corrsin (1951*a*) also showed that

$$\frac{\lambda_\theta}{\lambda_f} = \frac{1}{\sigma^{\frac{1}{2}}} \simeq 1{\cdot}18 \quad \text{(for air)},$$

when either (*a*) the Reynolds and Péclet numbers are so small that the convective effects are negligible for both heat and momentum, or (*b*) the Reynolds and Péclet numbers are both large, but the ratio of the rates of growth of the integral scales is prescribed in terms of the ratio of the integral scales at an initial instant. Mills *et al.* (1958) obtained this same form as an asymptotic result for very large Reynolds and Péclet numbers, with the spectral shapes of velocity and temperature fluctuations represented completely by their respective inertial subranges. Conditions corresponding to either (*a*) or (*b*) or the assumption of Mills *et al.* do not obtain in the present case, but with the identity of small and large Reynolds and Péclet number estimates, we might expect the intermediate cases to be at least roughly the same. Excluding the data near the heating screen, we see that the theoretical estimate is high by 50 %.

3.5. *Higher moments of temperature fluctuations*

The temperature signal immediately downstream of the heating screen is essentially pulse-like, but this character largely disappears for $x_\theta/M_\theta \gtrsim 15$ or 20. However, different statistical properties of the temperature signal are sensitive to different degrees to the remnants of this pulse-like character. Here, we examine the skewness S_θ and flatness factor F_θ (figures 10*a*, *b*). These values were obtained by ensemble

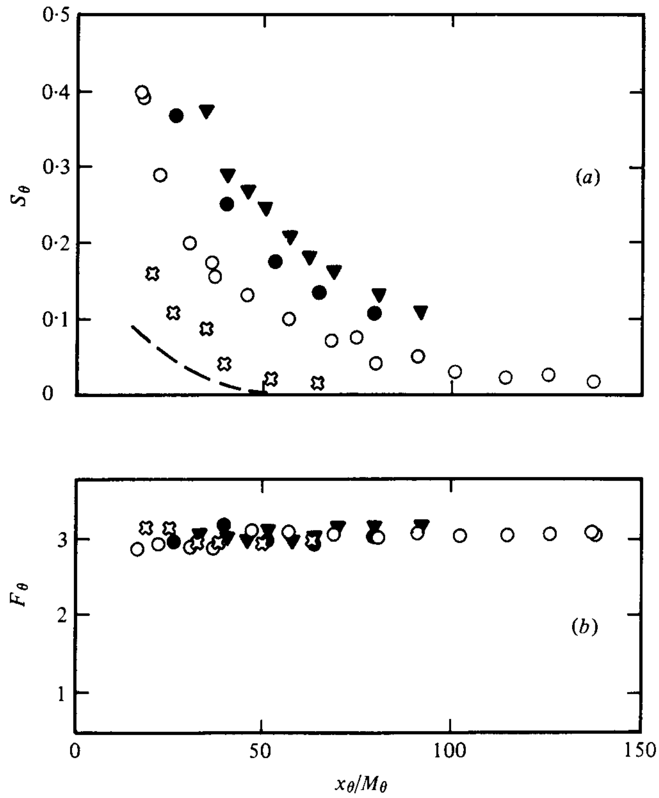

FIGURE 10. Downstream variation of the (a) skewness and (b) flatness factor of temperature fluctuation. ✧, heated grid, $M_\theta/M = 1\,(x_s = 0)$. ——, data of Mills *et al.*, $M_\theta/M = 1\,(x_s = 0)$. Other symbols (for the heated screen experiments) as in figure 7.

averaging about 50 records, each of about $1\frac{1}{2}$ s duration; the uncertainty due to the finite length of record is thus small. The skewness data of Mills *et al.* are also shown in figure 10(*a*). While $F_\theta \simeq 3$ for all cases for $x_\theta/M_\theta \gtrsim 25$, the skewness S_θ behaves differently for different experimental conditions; this is true whether the data are plotted against x_θ/M_θ or against x_θ/UT_0. This indicates that the skewness is a more sensitive parameter for detecting the remnants of the pulse-like character of the temperature fluctuation. In the heated grid experiments, including that of Mills *et al.*, S_θ decreases to nearly zero by about 40 thermal mesh sizes downstream of the grid, while in the heated screen experiments, it takes distances of the order of $100M_\theta$ or more to attain near-zero values; thus we have a case of skewed temperature fluctuations even in nearly isotropic temperature field. This persistence of skewness is however not inconsistent with the isotropy of the temperature field, because there is no mathematical reason why isotropic temperature fluctuations should be statistically symmetric except as a result of the mixing action of isotropic (hence non-skew) velocity fluctuations. But as we have already seen in the case of length scales, this may take a relatively long time.

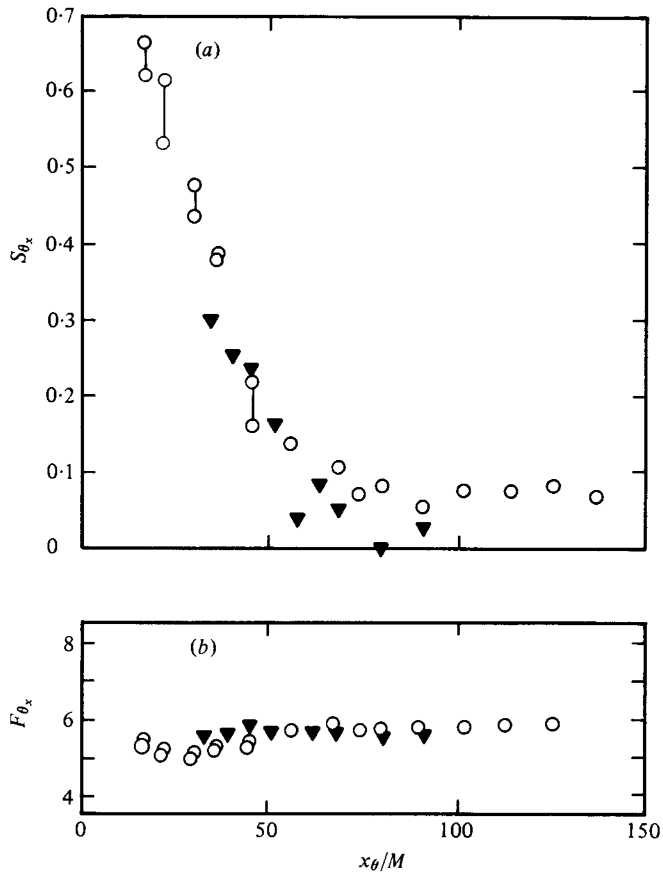

FIGURE 11. Downstream variation in the heated screen experiments of the (*a*) skewness and (*b*) flatness factor of the streamwise derivative of the temperature fluctuation. Symbols as in figure 7.

3.6. *Statistical properties of temperature derivative*

If turbulent fluctuations are isotropic, reflexional symmetry requires that the skewness of the temperature derivative be zero. More generally, this result must hold when the Reynolds and Péclet numbers are sufficiently large for $\overline{(\partial\theta/\partial x)^3}$ to be determined entirely by locally isotropic wavenumbers. In contrast, it is now well known that in flows with mean shear (both moderately large Reynolds number laboratory flows, and high Reynolds number atmospheric flows), the skewness is a non-zero constant of magnitude about 0·8; see, for example, the summary sketch of Sreenivasan & Antonia (1978), Tavoularis (1978*b*), and Sreenivasan & Tavoularis (1980). Sreenivasan & Antonia showed that the asymmetry of the large-scale ramp-like signals occurring in temperature traces is responsible for this anomalous behaviour. Since these ramps cannot be observed in the temperature fluctuations of isotropic grid turbulence, measurement of the skewness of temperature derivative in the present flow is of some interest.

Figures 11 (*a*, *b*) show the variation of the skewness S_{θ_x} and the flatness factor F_{θ_x} of the temperature derivative $\partial\theta/\partial x \simeq (1/U)\,\partial\theta/\partial t$ in the heated screen experiments. For $x_\theta/M_\theta \lesssim 60$, S_{θ_x} is a decreasing function of x_θ although, close to the heating screen, its magnitude depends on the precise location of the probe (since homogeneity has not yet been attained here). For $x_\theta/M_\theta \gtrsim 60$, $S_{\theta_x} \simeq 0{\cdot}08$. Considering possible measurement errors and the fairly low Reynolds number of the flow, the result is consistent with local isotropy.

Flatness factor, on the other hand, is not very sensitive to the initial departures from homogeneity and isotropy, and is about 5·5. This value compares well with that of Antonia *et al.* (1978) measured far away from the grid.

4. Further discussion of results

From (1) and (3), it follows that the Reynolds number R_λ is a slowly decreasing function of x; at $x/M \simeq 40$, $R_\lambda \simeq 34$. The ratio $(\lambda_f/L_f)\, R_\lambda$ is about 18. In the region of approximate homogeneity of the temperature field, the Péclet number P_{λ_θ} ($\equiv u'\lambda_\theta/\nu$) is essentially a constant for each experiment. Among various experiments mentioned here, it varies between 14 and 20; the ratio $(\lambda_\theta/L_\theta)P_{\lambda_\theta}$ is about 13. As in previous experiments, the Reynolds number of the present experiments is low. It is thus difficult to make a realistic assessment of the accuracy and practical applicability of the theoretical results of Corrsin (1964) and Newman *et al.* (1977) relating the length scale ratio to the relative rates of decay, because the theoretical estimates are based on high Reynolds and Péclet number approximations.

Newman *et al.* and Warhaft & Lumley studied essentially the same problem, using heated grid configurations in which every rod or every alternate rod was heated; they also introduced a fine-wire screen (which was made, however, of only an array of thin parallel wires rather than a square-mesh) at different locations downstream of the turbulence-generating grid. One of their experiments, with the wire-distance equal to the grid mesh size and the screen located at $x_s/M = 20$ is directly relevant to the present configuration.

Newman *et al.* and Warhaft & Lumley found that different experiments on the decay rates downstream of the heating system (grid in all cases except the Lin & Lin (1973) experiments; see later discussion) give vastly different values of the decay exponent m (varying between 0·87 and 3·1, according to figure 1 of Newman *et al.* and Warhaft & Lumley (1978*a*, *b*)) unlike the case of velocity data. Warhaft & Lumley suggested that this was due largely to the differences in heating powers, the speculation being that different degrees of heating affected the effective geometry of the thermal field differently, even though the solid boundary geometry remained unaltered. They found that m increased with increased heating,† and determined a correlation between them based on previous data as well as their own. We have seen that the present data contradict this finding; the decay exponent m, as well as the intensity of fluctuations normalized by ΔT, remained constant (to within measurable accuracy) for a given physical configuration of the heating grid, independent of changes in the moderate heating level. (Changing the configuration of the heating grid changes both m and β in equation (4), but we shall come back to this question later.) It is thus useful to re-examine some of the data on which Warhaft & Lumley based their conclusion.

First, among the different experiments they considered, the two sets of data with the lowest values of m are those of Mills *et al.* and one of their own. We recall that

† However, the one set of Lin & Lin's data which Warhaft & Lumley ignore because of large scatter in the decay measurements differs from the bulk of the Warhaft & Lumley correlation.

the former data did not obey a power law of the type (4), in contrast to all the other data, and thus must be treated with some reserve in this regard. Note that in figure 3 all the decay data with $M_\theta/M = 1$ fall in a certain vicinity of each other, except the one set of Warhaft & Lumley with this unusually small m; this at least suggests the presence of some unexplained feature with regard to this set of data. Furthermore, Warhaft & Lumley's suggestion that their lower grid heating 'essentially simulates the data of Mills *et al.*' is not persuasive, since the two experiments had appreciably different ΔT, equal to 1·66 °C and 5·0 °C, respectively. The present grid data with $\Delta T \simeq 2{\cdot}1$ °C and 4·4 °C, plus other data with 5 °C $\leqslant \Delta T \leqslant$ 10 °C support the conclusion that within the $\pm$ 10 % scatter mentioned earlier, the temperature decay rate was practically independent of ΔT for the range 2·1 °C $\leqslant \Delta T \leqslant$ 10 °C. Secondly, consider the data points with the largest m, which are obtained from the experiments of Lin & Lin (1973). The configuration used by Lin & Lin was in fact not a grid in the conventional sense, but a complex structure of heating elements repeatedly folded into compact flow channels, which themselves were arranged in a square-mesh fashion. The effective solidity of this configuration being high, an unstable flow might have resulted downstream (Corrsin 1963). In addition, the high overheats (of up to 80 °C) used by Lin & Lin could lead to further uncertainties. In any case, if as pointed out by Warhaft & Lumley the changes in decay rate are due to the implicit changes in the initial conditions caused by heating, physically different initial conditions ought to be at least as relevant, and so Lin & Lin's data should be considered an exceptional case which, though valuable in other contexts, are irrelevant to the present considerations. Excluding thus the doubtful possibilities which happen to be on either extreme in the range of values of m, we see that among heated grid experiments with $M_\theta/M = 1$, m varies by no more than $\pm$ 10 % about a mean value of 1·31.

Heating only alternate rods (i.e. making $M_\theta/M = 2$) increases β by a factor of 6 to 7, and also increases m to a value of about 1·55. Again, both m and β are independent of ΔT in the present data. The only other source of similar data are three of Warhaft & Lumley experiments with ΔT's of about 3·3 °C, 2·0 °C, and 0·5 °C, having m equal to 1·66, 1·52, and 1·26, respectively. (These values of m are slightly lower than those given by Warhaft & Lumley due to a different effective origin in our plot.) Interestingly, the first two sets of data having m values effectively the same as the present value also plot very close to the present data; that with a somewhat different m (= 1·26) plots considerably farther away, as in the experiments with $M_\theta/M = 1$.

Changes in m (though slight in the present experiments) brought about by changes in the heating grid configuration suggests a possible connection between the rate of decay of $\overline{\theta^2}$ and the initial conditions corresponding to the way in which the temperature field is created. Warhaft & Lumley suggested that this may come about by the persistence of a significant correlation between the two fields. Figure 12 shows that the correlation coefficient $\rho_{u\theta}$ between u and θ is about $-0{\cdot}20$ (instead of the isotropic value zero) when $M_\theta/M = 1$. The correlation between u and θ is somewhat stronger in the experiments of Warhaft & Lumley (1978), while it is much weaker in the experiments of Yeh & Van Atta (1973) (figure 12). Yeh & Van Atta suggested that the persistent negative correlation $\rho_{u\theta}$ indicates that the large eddies in the shear layers only slowly forget the method of their initial generation. The negative sign for $\rho_{u\theta}$ appears reasonable because in the heated wakes of the individual grid rods the velocity is low and the temperature is high relative to the mean values downstream. The

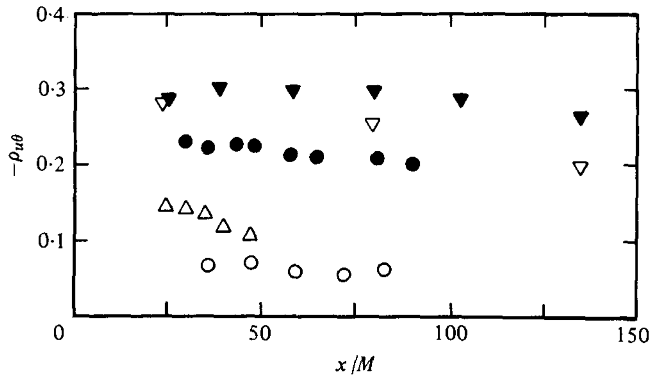

FIGURE 12. Cross-correlation between u and θ for the heated grid experiments. ●, $M_\theta/M = 1$; ○, $M_\theta/M = 2$. ▼, ▽, Warhaft & Lumley, $M_\theta/M = 1$; △, Yeh & Van Atta, $M_\theta/M = 1$.

differences in the magnitude of $\rho_{u\theta}$ among the various experiments suggest that it may depend in some complicated way on the grid solidity, the temperature level in the shear layers separating from the grid rods, etc. Yeh & Van Atta's suggestion also implies that, for a given grid, the correlation coefficient $\rho_{u\theta}$ must be smaller when only alternate rods are heated. Figure 12 shows that $\rho_{u\theta}$ is indeed nearly zero when $M_\theta/M = 2$.

Although Warhaft & Lumley did not measure integral scales, they inferred from the peaks of their three-dimensional spectra of temperature fluctuations that the length scale of the temperature field increased with increase in the grid heating. Excluding their data, table 2 shows that different experiments have been performed at different ΔT, while figure 6 shows no obvious trend of L_θ with ΔT. It appears reasonable to expect a certain range of small ΔT over which the characteristics of the thermal field should be independent of ΔT. If a critical ΔT does indeed exist, presumably this is a function of other parameters of the flow, and would thus be different for different flows.

Figure 13 shows the variation of the length scale L_θ with ΔT measured at two downstream locations $x/M = 50$ and 73. In this figure, each point is an average of several length scale measurements. It appears that, *for the present experimental configurations*, L_θ remains constant up to a ΔT of 2–3 °C, but increases beyond that.

It is not clear why L_θ should increase with ΔT, particularly because the increase in L_θ is not reflected in the decay data of temperature fluctuations (since m is independent of ΔT). It is possible that this anomalous behaviour becomes observable in the decay rates at a much higher critical value of ΔT than that corresponding to L_θ. A similar effect of ΔT on L_f is not as pronounced (see figure 13) as with L_θ.

Three points of difference of the heated screen experiments from the heated grid experiments are that, in the former case,

(i) m_s is independent of M_θ/M (as well as small ΔT), this indicates that the temperature field is independent of initial conditions under which it is created;

(ii) L_θ/L_f does not attain a near-unity constant that one expects, even after the mixing action has persisted for 5–6 integral time scales of turbulence at $x_\theta = 0$;

(iii) the skewness of the temperature field does not become zero as fast with x_θ/M_θ

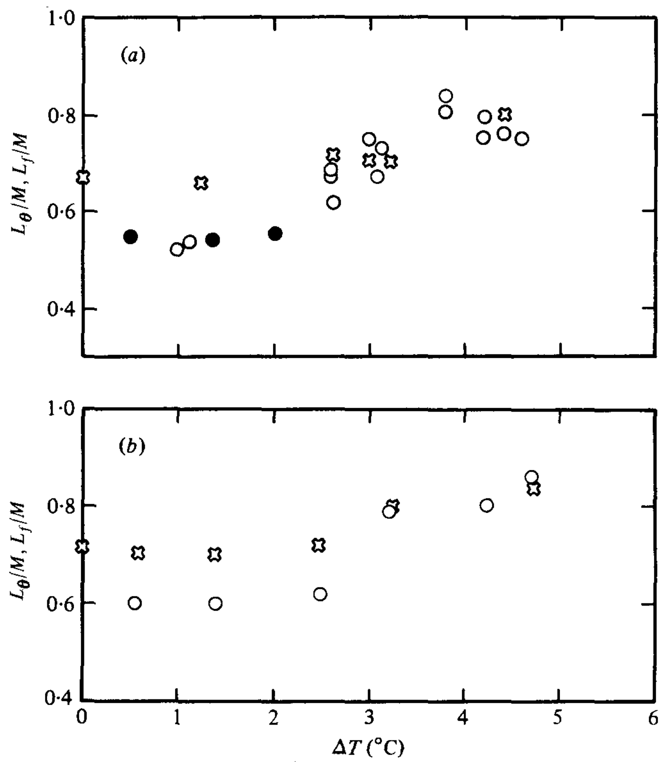

FIGURE 13. Variation of the temperature integral scale L_θ/M in the heated grid experiments and the velocity scale L_f/M with ΔT: (a) at $x/M = 50$, (b) at $x/M = 73$. ○, L_θ/M for $M_\theta/M = 1$; ●, L_θ/M for $M_\theta/M = 2$; ✡, L_f/M (with all rods heated).

as in the heated grid experiments, which reminds us that an isotropic scalar field variable can be skew.

Again, no dependence of m_s on ΔT was observed. This result, although not contradicted, cannot be confirmed by the data of Warhaft & Lumley, since their three experiments with different ΔT at a fixed M_θ (table 3) also correspond to different dimensionless distances between the screen and the grid†. Consistent with the observations of Warhaft & Lumley, $|\rho_{u\theta}|$ was found to be numerically less than 0·1.

5. Conclusions

In these experiments it has been found that, for a given configuration of the heating grid, the temperature decay data can be represented by a relation of the type (4) with both m and β substantially independent of ΔT. For the case with $M_\theta/M = 1$, the average value of m is about 1·31, with a $\pm 10\,\%$ variation. The ratio R ($\equiv L_\theta/L_f$) is about 0·8 and is independent of x. For ΔT beyond a particular critical value, it appears that L_θ may increase with ΔT; this unexplained feature is possibly a function of the particular grid and its solidity. Because the effects of ΔT on L_f are less certain, one cannot attribute this effect to a modified velocity field.

† On the other hand, in a recent manuscript Warhaft has confirmed the observation that $\overline{\theta^2}$ decay rate behind the *mandoline* is independent of overheat for small overheat.

For the case with $M_\theta/M = 2$, β increases six- to seven-fold, while m increases (by a magnitude of the order of 10 %) to about 1·55. The temperature length scale L_θ does not increase perceptibly. With only alternate rods heated, the correlation between u and θ (which was found to be about $-0{\cdot}2$ when $M_\theta/M = 1$) is nearly zero. The effect of ΔT on L_θ is not known in this case because, in the present experiments, ΔT could not be increased beyond about 2·7 °C.

Heating all the horizontal rods only increases β by a factor of $2\frac{1}{2}$, but does not increase m significantly when compared with the case $M_\theta/M = 1$. Similarly, heating only the alternate horizontal rods increases β six-fold, but does not affect m, when compared with the case $M_\theta/M = 2$.

In the heated screen experiments, there exists no appreciable initial coupling of the velocity and temperature fields. The value of m is thus independent not only of ΔT, but also of M_θ/M and of the downstream distance x_s of the location of the heating screen with respect to the grid (at least for the range covered here); β increases with M_θ/M as in the heated-grid case, and also with the screen distance x_s. However, for a given M_θ/M and x_s, β is also independent of ΔT. The ratio R of the integral scales L_θ/L_f slowly approaches unity approximately exponentially when the 'time of flight' is expressed in terms of the characteristic turbulence decay time at $x_\theta = 0$.

Corrsin's (1964) high-Reynolds-number estimate of the relative rates of decay of velocity and temperature intensities in terms of the ratio of the integral scales is generally unsuccessful, although in the heated grid case, it agrees roughly with experiment. The modification suggested by Newman *et al.* (1977) is generally more successful in the heated screen experiments, but gives poorer agreement for the heated grid experiments. The general failure of either theory to predict the length scale ratio correctly in both cases may be due in part to the fact that the experiments are at comparatively low Reynolds and Péclet numbers.

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