A review of research on the effects of gravity on combustion processes is presented, with an emphasis on a discussion of the ways in which reduced-gravity experiments and modeling has led to new understanding. Comparison of time scales shows that the removal of buoyancy-induced convection leads to manifestations of other transport mechanisms, notably radiative heat transfer and diffusional processes such as Lewis number effects. Examples from premixed-gas combustion, non-premixed gas-jet flames, droplet combustion, flame spread over solid and liquid fuels, and other fields are presented. Promising directions for new research are outlined, the most important of which is suggested to be radiative reabsorption effects in weakly burning flames.

Comparison of Time Scales for Premixed-Gas Combustion

To determine the conditions in which gravity can affect flames, the estimated time scales for chemical reaction \( t_{\text{chem}} \), inviscid buoyant convection \( t_{\text{invis}} \), viscous buoyant convection \( t_{\text{vis}} \), conductive heat loss to walls \( t_{\text{cond}} \), and radiant heat loss \( t_{\text{rad}} \) are compared. Premixed laminar flames are considered first because of their simplicity. Subsequent sections introduce time scales for other flames.

The chemical time scale is (see Nomenclature) \( t_{\text{chem}} = \delta/a \), where \( \delta = a/S_l \), thus \( t_{\text{chem}} = \alpha S_l^2 \). The convective transport time scale is \( d/U \), where \( d \) is a characteristic flow length scale, \( U \approx (gd/\rho)^{1/2} \) is the buoyant convection velocity, and \( \rho \) the density change across the flame. Because \( \Delta \rho \approx 1 \) for flames, \( t_{\text{invis}} = d/(gd)^{1/2} \approx (d/g)^{1/2} \). For inviscid flow, \( d \) is determined by the apparatus dimensions, for example the burner or tube diameter. For viscous flow, \( d \) cannot be specified independently; instead \( d \approx \sqrt{U} \), thus \( t_{\text{invis}} = d/(gd)^{1/2} = (d/g)^{1/2} \). For both \( d \approx \sqrt{U} \) and \( t_{\text{invis}} \approx d/(gd)^{1/2} \), the conduction time scale \( t_{\text{cond}} \) is the flame front temperature \( T_f \) divided by the rate of temperature decrease due to conductive loss; thus \( t_{\text{cond}} \approx \sqrt{\rho C_p \Delta T} \). Similarly, for optically thin radiation, \( t_{\text{rad}} \approx \sqrt{\rho C_p a \Delta T} \).

Two sets of time scales are shown in Table 1, one for near-stoichiometric hydrocarbon-air flames and one for near-limit flames, both at \( g = 1 \) atm. For near-stoichiometric flames \( S_l \approx 0.40 \text{ m/s}, T_f \approx 2200 \text{ K}, \alpha \approx 1.5 \times 10^{-4} \text{ m}^2/\text{s} \), and \( \Delta p \approx 0.56 \text{ m}^3/\text{m}^2 \). For near-limit flames, \( S_l \approx 0.02 \text{ m/s}, T_f \approx 1500 \text{ K}, \alpha \approx 1.0 \times 10^{-4} \text{ m}^2/\text{s} \), and \( \Delta p \approx 0.83 \text{ m}^3/\text{m}^2 \). For both cases \( g \approx 9.8 \text{ m/s}^2 \), \( \gamma \approx 13.5 \), \( T \approx 300 \text{ K} \), and \( d \approx 0.05 \text{ m} \) (a typical apparatus dimension).

Several observations can be made on the basis of these simple estimates:

1. Buoyant convection is unimportant for near-stoichiometric flames because \( t_{\text{vis}} \gg t_{\text{chem}} \) and \( t_{\text{invis}} \gg t_{\text{chem}} \).
2. Buoyant convection strongly influences near-limit flames at \( g \) because \( t_{\text{vis}} \ll t_{\text{chem}} \) and \( t_{\text{invis}} \ll t_{\text{chem}} \).
3. Radiation effects are unimportant at \( g \) compared to buoyant convection because \( t_{\text{rad}} \ll t_{\text{vis}} \) and \( t_{\text{chem}} \ll t_{\text{rad}} \).
4. Radiation effects dominate near-limit flames because \( t_{\text{rad}} \approx t_{\text{chem}} \), but these effects are only observable at reduced gravity because of observation 3.

5. The apparatus size \((d)\) must be larger than about 0.03 m to observe radiation-induced extinction; otherwise, conduction losses exceed radiative losses \((t_{\text{cond}} < t_{\text{rad}})\).

6. Many radiative loss effects can be studied in drop-towers (test duration 2-10 s), because these times are typically larger than \( t_{\text{rad}} \).

7. Because \( t_{\text{rad}} \approx \frac{g}{d} \) and \( t_{\text{visc}} \approx \frac{g}{d^2} \), aircraft-based \( \mu g \) experiments at \( g \approx 10^{-2} \) \( g_0 \) may not provide sufficiently reduced buoyancy to observe radiative effects.

8. Because \( t_{\text{visc}} \approx \frac{g}{d^2} \) and \( t_{\text{rad}} \approx \frac{g}{d} \), but at higher \( P \), buoyancy effects interfere more strongly with radiative effects.

9. A Reynolds number \( Re_d = Ud/d \) for buoyant flow is estimated as \( (gd)^{1/2}d^2 = Gr_d^2 \), where \( Gr_d = \frac{gV^2}{d} \) is a Grashof number. For \( Re_d \approx 10^3 \), thus \( Gr_d \approx 10^6 \), buoyant flow at \( 1 g \) is generally turbulent, thus, it is difficult to obtain steady laminar flames in large systems at \( 1 g \).

The implications of these observations are discussed in the following sections. Essentially, any combustion process where \( t_{\text{chem}} \) or \( t_{\text{rad}} \) exceeds \( t_{\text{visc}} \) or \( t_{\text{visc}} \) may be affected by gravity and is worthy of \( \mu g \) investigation.

### Premixed Gas Flames

**Flammability Limits**

The previous section showed that gravity effects are significant only for mixtures with low \( S_L \), implying mixtures highly diluted with excess fuel, oxidant, or inert gas, but sufficient dilution causes flammability limits. Thus, gravity effects have significant influences on near-limit behavior, which is expected since limits are different for upward, downward, and horizontal propagation [5].

Practically all flammability limit studies show that burning velocity at the flammability limit \( S_{L,\lim} \) is nonzero. Giovangi and Snooile [6] have shown that there is no purely chemical flammability limit criterion for planar unstretched flames: without losses, \( S_L \) decreases asymptotically to zero as dilution increases. Consequently, loss mechanisms such as those discussed below are needed to explain limit mechanisms. The resulting predictions of \( S_{L,\lim} \) indicate that \( S_{L,\lim} \) usually depends only weakly on chemical reaction rate parameters. Thus, limit mechanisms may be inferred by comparing predicted and measured \( S_{L,\lim} \) without detailed chemical knowledge. The mixture composition at the limit affects \( S_{L,\lim} \) only weakly through \( T_f \); thus, comparing predicted and measured limit compositions is not especially enlightening; comparisons of \( S_{L,\lim} \) values is much more useful. Consequently, this discussion emphasizes comparisons of predicted and measured values of \( S_{L,\lim} \).

For upward propagation, Levy [7] showed that the flame rise speed at the limit \((S_{L,\lim})\) is identical to that of an inviscid hot gas bubble. This relation was later verified for a wide range of tube diameters and mixtures [8]. Buckmaster and Mikhail [9] showed how this minimum rise speed causes hydrodynamic strain at the flame tip, which causes extinguishment for sufficiently low \( S_L \). The predicted burning velocity of the limit mixture \((S_{L,\lim})\) is, after temperature-averaging transport properties:

\[
S_{L,\lim} = 2.8 \exp\left[\frac{P}{4} \left(1 - \frac{1}{L_e}\right)\left(1 - \frac{T_f}{T_{\text{ch}}}ight)\left(\frac{g}{d}\right)^{1/4}\right]
\]  

\( (1) \)

The form of equation 1, \( S_{L,\lim} \approx (g^2d/\rho)^{1/4} \), can be obtained by setting \( t_{\text{visc}} = t_{\text{chem}} \).

For downward propagation in tubes, centrifuge experiments [10] indicate that \( S_{L,\lim} \approx g^{1/3} \), independent of \( L_e \), which reasonable since downward-propagating near-limit flames are nearly flat and unstrained. Experiments [11] and numerical simulations [12] suggest extinction results from sinking regions of cooling burned gas near the walls over-taking the flame and blocking it from fresh reactants, although the \( g^{1/3} \) scaling was not tested in Refs. [11,12]. The \( g^{1/3} \) scaling can be obtained by setting \( t_{\text{chem}} = t_{\text{visc}} \) thus

\[
S_{L,\lim} \approx (g^2d)^{1/4}
\]  

\( (2) \)

Experiments [8] employing varying diluent gases and pressures confirm the \( g^{1/3} \) scaling and lack of dependence on tube size, which, together with the \( g^{2/3} \) scaling mentioned above, supports the proposed mechanism.

Both upward- and downward-limit mechanisms indicate that as \( g \to 0 \), \( S_{L,\lim} \to 0 \), implying arbitrarily
weak mixtures could burn very slowly. However, conductive or radiative losses prevent arbitrarily weak mixtures from burning even at $g = 0$. Theories that relate flammability limit to heat losses \[13-15\] predict a minimum $T_f$ below which flame propagation cannot occur because chemical reaction rates are much stronger functions of temperature than heat loss rates (exponential vs. algebraic). Consequently, because dilution decreases $T_f$, dilution increases the impact of losses, leading to flammability limits. For conductive losses, setting $t_{\text{chem}} = t_{\text{cond}}$ leads to

$$P_{\text{h,lim}} = S_{\text{h,lim}} d/\alpha = \text{constant} \quad (3)$$

with experiments \[8,16\] and computations \[17\] indicating $P_{\text{h,lim}} \approx 40$. For radiative losses, setting $t_{\text{chem}} = t_{\text{rad}}$ leads to \[18,19\]

$$S_{\text{r,lim}} = \frac{1}{\rho \cdot C_p} \sqrt{\frac{1.2 \beta A L_f}{T_f}} \quad (4)$$

For lean-limit CH$_4$-air mixtures at 1 atm, equation 4 yields $S_{\text{r,lim}} \approx 0.023$ m/s, similar to detailed numerical model predictions \[20,21\]. Such small $S_{\text{r,lim}}$ are not observed at 1g because of buoyant convection ($t_{\text{conv}} < t_{\text{rad}}$ and $t_{\text{rad}} < t_{\text{chem}}$); equations 1 and 2 yield $S_{\text{r,lim}} \approx 0.033$ and 0.078 m/s for upward and downward propagation, respectively. At $\mu g$, however, predictions of equation 4 compare favorably to experiments in large combustion vessels \[19,22\] using varying pressures, fuels, and inert gases. Also, similar results were obtained for CH$_4$-air mixtures at 1 atm in tubes with $d = 0.056$ m \[23\], suggesting these limits are apparatus-independent. Thus, radiative losses may cause flammability limits when extrinsic losses (conduction, buoyant convection, etc.) are eliminated. In this instance, $\mu g$ experiments enabled observation of phenomena not observable at 1g.

These radiative effects apply only for optically thin gases (i.e., no reabsorption of emitted radiation), which are inappropriate for large systems, high pressures, or mixtures with strongly absorbing material. With this motivation, $\mu g$ experiments \[24\] were conducted using lean CH$_4$-air mixtures seeded with SiC particles. Because solids emit/absorb as black- or gray-bodies, whereas gases radiate in narrow spectral bands, particle-seeded gases emit/absorb more radiation than particle-free gases. Measurements of propagation rates, pressures and postflame thermal decays showed that, consistent with theoretical predictions \[25\], at low particle loadings the particles increase radiative loss (optically thin conditions), whereas at higher loadings reabsorption of emitted radiation becomes significant, which decreases net radiative loss and augments conductive heat transport.

Even for gases, computations \[26\] using detailed statistical narrow-band radiation models show that flammability limits are extended considerably with reabsorption (Fig. 1). With gases, however, two mechanisms lead to flammability limits even with reabsorption. One is the difference in composition between reactants and products; if H$_2$O or other radiatively active combustion products are absent from the reactants, radiation from these species that is emitted upstream cannot be reabsorbed by the reactants. The second mechanism is that emission spectra are broader at $\mu g$ than $g$, thus, some radiation emitted near the flame front cannot be absorbed by the reactants. Upstream loss occurs via both mechanisms, leading to extinction of weak mixtures. These results suggest that fundamental (domain- and gravity-independent) flammability limits due to radiative losses may exist at $\mu g$, but these limits are strongly dependent on emission/absorption spectra of reactant and product gases and their temperature dependence and cannot be predicted using gray-gas or optically-thin model parameters.

**Stretched Flames**

Premixed gas flames generally are not flat and steady nor do they propagate into quiescent flows. Consequently, flames are subject to “flame stretch,” $\Sigma = (1/A)(dA/dt)$ \[18\], which affects $S_t$ and extinction conditions \[19,27\]. At 1g, buoyancy imposes flame stretch comparable to $t_{\text{conv}}$ or $t_{\text{rad}}$. At $\mu g$, weak flame stretch effects that are insignificant at 1g may dominate. One example is expanding spherical flames for which

$$\Sigma = \frac{1}{A} \frac{dA}{dt} = \frac{1}{4\pi r^2} \frac{d}{dt} \left(4\pi r^2\right) = \frac{2}{r} \frac{dr}{dt} \quad (5)$$

For $Le < 1$, positive stretch increases $T_f$ because the increased chemical enthalpy diffusion to the flame front in the form of scarce reactant exceeds the increased thermal enthalpy loss. Because heat release
The curvature effect \( (2S/R) \) opposes heat loss \( (Q) \), allowing mixtures that are nonflammable as plane flames \( (Q > 1/e) \) to exhibit expanding spherical flames until \( r_i \) grows too large and thus the curvature benefit too small. For mixtures just outside the limit, the extinction radius may be very large. Such behavior, termed self-extinguishing flames (SEFs), is observed experimentally \([19,22]\) (Fig. 2) at \( \mu g \) in near-limit mixtures with \( Le \) slightly less than unity. (Mixtures with lower \( Le \) exhibit diffusive-thermal instabilities or flame balls discussed below.) Equation 6 also predicts, consistent with experimental observations \([29]\), that SEFs cannot occur for \( Le > 1 \) (thus, \( R < 0 \)) because both curvature and heat loss weaken the flame.

Two experimental observations not predicted by equation 6 are that narrow mixture ranges exhibit both SEFs and normal flames and that the energy release before extinguishment can be orders of magnitude greater than the ignition energy. Such behavior is predicted by computations \([30]\) not subject to scaling limitations of activation energy asymptotics used to derive equation 6. These calculations also show that for small initial \( r_i \), all mixtures exhibit extinguishment, corresponding to nonignition behavior \([31]\). Thus, in mixtures exhibiting SEFs, flames extinguish at large curvature (small \( r_i \)) due to large \( \Sigma \) and at small curvature due to radiative losses. This dual-limit behavior is also exhibited by many other types of flames described later.

Flames in hydrodynamic strain induced by counterflowing round jets are frequently employed to model turbulence-induced flame stretch effects. At steady state, the flame resides at the axial location \( y \) where the axial velocity \( (U_t) \) equals \( S_t \) for the given \( \Sigma = dU/dy \). As \( \Sigma \) increases, \( U_t \) increases, thus the flame moves toward the stagnation plane (smaller \( y \)) and the burned gas volume (thus radiative loss) decreases. As with curvature-induced stretch, for \( Le \) less than/greater than unity, moderate hydrodynamic strain increases/decreases \( S_t \), but for all \( Le \), large strain extinguishes the flame \([32]\). Consequently, \( \mu g \) experiments \([33]\) in low-\( Le \) mixtures (Fig. 3) reveal extinguishment behavior analogous to spherical flames. For large \( \Sigma \), the short residence time \( (\sim \Sigma^{-1}) \) causes extinguishment \( (\Sigma^{-1} \sim t_{bem}) \) along the “normal flame” branch, analogous to nonignition behavior of spherical flames. In contrast, for low \( \Sigma \), the residence time and burned gas volume are large; thus, radiative loss is significant \( (t_{rel} \sim t_{bem}) \), so radiative loss extinguishes the flame along the “weak flame” branch, analogous to SEFs. The optimal \( \Sigma \) \( (\sim 13 \text{ s}^{-1}) \) producing the minimum flammable fuel concentration corresponds to \( \Sigma^{-1} = 0.08 \text{ s} \), which is less than \( t_{rel} \) or \( t_{bem} \); thus, the C-shaped response and the entire weak-flame branch cannot be observed at \( 1g \). The optimal \( \Sigma \) is nearly the same for model and experiment, indicating that loss rates are modeled well, but the computed limit

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**Fig. 2.** Characteristics of self-extinguishing flames in CH₄-air mixtures at 1 atm for various mole percent CH₄ and spark ignition energies \([112]\). The *x* symbols denote extinction.

**Fig. 3.** Measured and predicted extinction strain rates for strained premixed CH₄-air flames at \( \mu g \) \([33]\) showing dual-limit behavior, that is, residence-time limited extinction at high strain rates (upper branch, “strong flames”) and radiative loss extinction at low strain rates (lower branch, “weak flames”).

Reactions have high activation energies, small \( T_f \) causes large changes in reaction rate and thus \( S \). An evolution equation for nonadiabatic expanding spherical flames is given by \([28]\):

\[
\frac{dS}{dR} + S^2 \ln S^2 = \frac{2S}{R} - Q \tag{6}
\]

where \( S = (dr_i/dt)(S_i(\rho_i/\rho_f)) \), \( R = r_i/(\beta \beta_i(Le, e)) \), \( \beta(Le, e) \) is a scaling function \( (I > 0 \text{ for } Le < 1 \text{ and } I < 0 \text{ for } Le > 1) \) and \( Q = [\beta A(T_f^{-1})]^{-1}(T_f^{-1} - T_2) \). The terms in equation 6 represent unsteadiness, heat release, curvature-induced stretch, and heat loss, respectively. For steady planar flames, equation 6 becomes \( S^2 \ln S^2 = -Q \), which exhibits a maximum \( Q = 1/e = 0.3678 \ldots \) at \( S = e^{-1/2} \), which corresponds to \( S_{lim} \) from equation 4. For \( Le < 1 \),
I several new extinction branches depending on plus the reduced radiative loss at larger nism to reduce radiative loss by flame expanding flames, no flammability extension occurs [34]. (For spherically ciently high decreases at moderate δ, the flammability limit extension also occurs for 2 > 1, though, for sufficiently high 2, no C-shaped response or flammability limit extension occurs [34]. (For spherically expanding flames, no flammability extension occurs for 2 > 1 because in this case there is no mechanism to reduce radiative loss by flame stretch.)

The combination of nonmonotonic response to δ plus the reduced radiative loss at larger δ causes several new extinction branches depending on δext, δint, δ−1, and 2 [34,35]. It is uncertain whether these branches are physically observable because they have not been identified experimentally and stability analyses have not been performed.

### Flame Balls

Over 50 years ago, Zeldovich [36] showed that the steady mass, energy, and species conservation equations admit solutions corresponding to stationary spherical flames, characterized by a flame radius (rf). Fuel and oxygen diffuse from the ambient mixture inward to the reaction zone while heat and combustion products diffuse outward (Fig. 4). Mass conservation requires that the fluid velocity be zero everywhere. The temperature and species mass fraction profiles have the form c1 + c2r, where c1 and c2 are constants. Corresponding solutions in planar and cylindrical geometry cannot exist because the solution forms c1 + c2r and c1 + c2r, respectively, are unbounded as r → ∞. Zeldovich [36] and others [37,38] also showed that flame-ball solutions are unstable and thus probably not physically observable, although these solutions are related to flame ignition [37].

Forty years after Zeldovich [36], apparently stable flame balls were accidentally discovered in drop-tower experiments using H2-air mixtures [39] and aircraft-based μg experiments using various low-2 mixtures [40]. The μg environment facilitated spherical symmetry and prevented buoyancy-induced extinction. For mixtures sufficiently far from flammability limits, expanding spherical fronts composed of many individual cells were observed, whereas for more dilute mixtures, cells that formed initially did not split and instead closed up on themselves to form flame balls. (For still more dilute mixtures all flames eventually extinguished.) It was inferred that flame balls can occur in all near-limit low-2 mixtures; however, the short duration of drop-tower experiments and substantial g fluctuations in the aircraft-based μg experiments precluded definite conclusions. Recent space shuttle experiments [41] confirmed that flame balls can exist for at least 500 seconds (the entire experiment duration.) Zeldovich [36] noted that radiative losses might stabilize flame balls; consequently, after their experimental observation, radiative loss effects on flame balls were analyzed [42]. For moderate loss, two solutions branches are predicted (Fig. 5), a strongly nonadiabatic large-radius branch and a nearly-adiabatic small-radius branch. For sufficiently strong losses, no solutions exist, indicating extinction limits. Stability analyses [42] predict that all small flames are unstable to radial disturbances, and large flames with weak loss (far from flammability limits) are unstable to three-dimensional disturbances. Close to extinction limits, the large-radius branch is stable to both disturbances. These predictions are consistent.
with the observed splitting cellular flames away from limits and stable balls close to limits. For $Le$ close to or larger than unity, all flame balls are unstable for any loss magnitude [43], explaining why they are never observed in (for example) CH$_4$-air mixtures ($Le \approx 0.9$) or C$_3$H$_8$-air mixtures ($Le \approx 1.7$).

Numerical predictions of non-adiabatic flame balls employing detailed chemistry, diffusion, and radiation models [44,45] are qualitatively consistent with these experimental and theoretical results. Still, quantitative agreement has been elusive (Fig. 6) for at least two reasons. First, flame-ball properties are very sensitive to the three-body recombination step $H + O_2 + H_2O \rightarrow HO_2 + H_2O$ [45] whose rate varies widely between different published $H_2-O_2$ reaction mechanisms. The second reason is reabsorption of emitted radiation in mixtures diluted with radiatively active CO$_2$ or SF$_6$. An upper bound on self-absorption of diluent radiation ($a_T \rightarrow \infty$) is assessed by neglecting diluent radiation entirely because as $a_T \rightarrow \infty$, radiative loss from the diluent vanishes and no additional heat transport occurs due to radiation. Agreement between predicted and measured flame radii is much better in this case (Fig. 7), strongly suggesting radiation modeling including reabsorption is needed for accurate predictions in these cases.

A key difference between propagating flames and flame balls is that propagating flames have convective-diffusive zones where temperature and concentration approach their ambient values in proportion to $e^{-a'T}$, whereas flame balls have purely diffusive zones where the approach is proportional to $1/r$. Plane flames respond on short time scales $t_{flow} = \delta^2/a$, whereas the gradual $1/r$ flame-ball profiles produce properties dominated by the far-field length scale $\beta r$, and thus diffusion time scales $(\beta r)^2/a$ [42], typically 100 s. These scales are relevant to stability and extinction limits because they affect the times for radiant combustion products to diffuse to the far-field and indicate large volumes of gas ($\sim \beta^2 r^2$) where radiative loss affects flame balls. Such large scales are confirmed by space experiments [41] and numerical simulations [44,45]. Droplet and candle flames (discussed later), which have quasi-spherical, diffusion-dominated far-fields, exhibit analogous behavior.

**Gaseous Non-Premixed Flames**

**Stretched Flames**

Non-premixed flames, where fuel and oxidant are separated before combustion, are affected by stretch differently than are premixed flames. The most significant difference is that the flame position is determined by the location of stoichiometric mixture fraction, dictated by mixing considerations, rather than being determined by balances between $S$ and $U$ as in premixed flames. Consequently, non-premixed flames have considerably less freedom of movement. Also, premixed flames have characteristic thicknesses $\delta \sim a/S$ unrelated to flow environment, whereas non-premixed flames have only the diffusion length scale $\delta \sim (a/S)^{1/2}$. With fixed flame location and $a$ increasing monotonically with decreasing $S$, non-premixed flames with radiative loss exhibit only simple C-shaped responses to strain (Fig. 8) [46], with a short residence time extinction branch ($t_{flow} > \Sigma^{-1}$) and a radiative loss extinction
I
nered by the model, probably axial conductive heat
behavior was seen for premixed flames (Fig. 3). This
model predicts flames at arbitrarily low
flames exist below some value of \( Z' \), whereas the
and the large-radius branch of flame balls (Fig. 5).
lower branch of strained premixed flames (Fig. 3)
cussed later. It is also somewhat analogous to the
tinction of other types of non-premixed flames dis-
time (\( t_{\text{rad}} \)) for non-premixed round-jet and slot-jet flames
branch (\( t_{\text{rad}} < \Sigma^{-1} \)) rather than the complica-
ted responses found for premixed flames [34]. The
only significant difference in flame structures near
the two limits is \( \delta \) [46], a situation quite unlike pre-
mixed flames.

For the radiative extinction branch \( t_{\text{chem}} \) is still a
factor because \( t_{\text{rad}} = \Sigma^{-1} \) results in order unity de-
creases in flame temperature, thus causing exponen-
tially large decreases in \( t_{\text{chem}} \). Even conditions far
from extinction at \( l_0 \), therefore, may exhibit radiat-
ive extinction at \( \mu \) due to much larger residence
times. This mechanism also applies to radiative ex-
tinction of other types of non-premixed flames dis-
cussed later. It is also somewhat analogous to the
lower branch of strained premixed flames (Fig. 3)
and the large-radius branch of flame balls (Fig. 5).

The experiments shown in Fig. 8 suggest that no
flames exist below some value of \( \Sigma \), whereas the
model predicts flames at arbitrarily low \( \Sigma \). Similar
behavior was seen for premixed flames (Fig. 3). This
suggests an additional loss mechanism not consid-
ered by the model, probably axial conductive heat
losses to the jets or radial conductive loss to inert
gases surrounding the reactant streams. This would
induce \( t_{\text{chem}} = 2.9 \text{ s}^{-1} \) if \( d \) is the jet spacing (25 \text{ mm}),
or \( 7.1 \text{ s}^{-1} \) if \( d \) is the jet diameter (16 \text{ mm}). Either of
these are roughly consistent with the minimum \( \Sigma \) in
Fig. 8. Thus, apparatus large enough to study
flames at \( l_0 \) without substantial conductive loss,
where maximum length scales are about \((\sqrt{\alpha_{\text{rad}}} ) \Sigma \), are
insufficient for the weaker flames attainable at \( \mu \),
where maximum length scales are about \((\sqrt{\alpha_{\text{rad}}} ) \Sigma \).

Laminar Gas-Jet Flames

A fuel jet issuing into an oxidizing environment is
one of the simplest types of flames. Jost [47] and
Roper [48] estimated the flame height (\( L_f \)) and resi-
dence time from jet exit to flame tip (\( t_{\text{jet}} \)) by deter-
mining the height (\( y \)) where the transverse diffusion
time (\( d(y)^2/D \), where \( d(y) \) is the stream tube width)
equals the convection time (\( U(y)y \), where \( U(y) \) is
the axial velocity). When buoyancy and viscosity ef-
fects are negligible (momentum-controlled jets),
\( U(y) \) is constant and equals the jet exit velocity (\( U_0 \)),
whereas when buoyancy effects dominate, \( U(y) \sim
(\gamma y)^{1/2} \). In either case, mass conservation requires
that \( d(y)^2 U(y) = d^3 U_0 \), constant for round jets or
\( d(y) U(y) = d_d U_0 \) = constant for slot jets. The re-
resulting estimated scalings for \( L_f \) and \( t_{\text{jet}} \) are given in
Table 2. Transition from buoyancy-controlled to mo-
tum-controlled conditions occurs when the time
scale for the former exceeds the latter, which cor-
responds to \( U_0 \gtrsim \gamma d^2/D \) for either round jets or slot
jets. The scalings for momentum-dominated flames
presume constant \( U \), which is reasonable for collo-
wing Burke-Schumann flames, but for nonbuoyant jet
flames without coflow, the jet spreads and decceler-
ates. For this situation [49]

\[
L_f \sim \frac{U_d d_o}{D} \frac{1}{\sqrt{2 Sc \gamma}} \left( \frac{1}{1-c_o} \right)^{1/2} \tag{7}
\]

\[
t_{\text{jet}} \sim \frac{d_d}{U(y)} \frac{d^3}{D} \left( \frac{1}{1-c_o} \right) \tag{8}
\]

Because \( Sc \approx 1 \), the scalings of \( L_f \) and \( t_{\text{jet}} \) are similar
with or without viscosity. Fig. 9 shows measurements
of \( L_f \) for CH\(_4\) flames [50]. Note that, as the scalings
predict, \( L_f/d_o \sim Re /U_d d_o \) at both \( l_0 \) and \( \mu \), and
only small differences exist between \( l_0 \) and \( \mu \) flame
lengths.

All \( \mu \) studies show larger flame widths (\( w \)) at \( \mu \)
than \( l_0 \) due to lower \( U \) and longer \( t_{\text{jet}} \) [50]. Also, \( w \)
is larger at \( \mu \) because the temperatures are lower
(see below) and \( D \sim T^{1.75} \). Because \( w \) depends on
whether \( U \) is accelerating (buoyant jets), constant
(nonbuoyant Burke-Schumann flames), or decceler-
ating (nonbuoyant jets), \( w \) is more difficult to predict
than is \( L_f \) [2]. The difference between \( l_0 \) and \( \mu \)
widths decreases as \( Re \) (thus, \( U_d \)) increases [Fig. 10].
The nonbuoyant widths increase slightly with \( Re \),

![Graph](https://via.placeholder.com/150)

**Fig. 8** Measured and predicted extinction strain rates for strained non-premixed N\(_2\)-diluted CH\(_4\) versus air counterflow flames at \( \mu \) showing dual-limit response analogous to premixed flames (Fig. 3).

**Table 2** Predicted scalings of flame heights (\( L_f \)) and residence times (\( t_{\text{jet}} \)) for non-premixed round-jet and slot-jet flames under momentum-dominated and buoyancy-dominated conditions.

<table>
<thead>
<tr>
<th>Geometry</th>
<th>Flow</th>
<th>Mechanism</th>
<th>( L_f )</th>
<th>( t_{\text{jet}} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Round-jet</td>
<td>Momentum</td>
<td>( U_d d_o/d )</td>
<td>( d^2/D )</td>
<td>( d_d U_0 )</td>
</tr>
<tr>
<td>Round-jet</td>
<td>Buoyant</td>
<td>( U_d d_o/d )</td>
<td>( (U_d d_o D)^{1/2} )</td>
<td>( d_d U_0 )</td>
</tr>
<tr>
<td>Slot-jet</td>
<td>Momentum</td>
<td>( U_d d_o/d )</td>
<td>( d^2/D )</td>
<td>( d_d U_0 )</td>
</tr>
<tr>
<td>Slot-jet</td>
<td>Buoyant</td>
<td>( (U_d d_o D)^{1/3} )</td>
<td>( (U_d d_o^2 D)^{1/3} )</td>
<td>( d_d U_0 )</td>
</tr>
</tbody>
</table>
This suggestion is supported by the results shown in Fig. 10 which show that $w/d_o$ is lowest for lowest $Re (=20)$, where axial diffusion is most significant, but asymptotes to fixed values at high $Re$.

Compared to yellow $lg$ flames, $\mu g$ gas-jet flames are more red [49,50], indicating lower blackbody soot temperatures and presumably lower maximum flame temperatures. This occurs because $t_w/\nu$ is larger at $\mu g$ and thus radiative loss effects ($\sim k\nu /
u$) are greater. Drop-tower [51] and space [52] experiments indicate surprisingly large and consistent radiative loss fractions (0.45–0.60) at $\mu g$ compared to 0.07–0.09 at $lg$, for various fuels, pressures, $O_2$ mole fractions, and flow rates. Thus, differences in $t_w$ at $lg$ and $\mu g$ result in widely varying characteristics even for flames having nearly the same $L_T$.

**Turbulent Flames**

In turbulent non-premixed jet flames, $D$ is not fixed but rather is nearly proportional to $u'/L_T$. Because $u' \sim U_o$ and $L_T \sim d_o$, $L_T \sim U_o u'/u' L_T \sim d_o$ for round jets; thus, $L_T$ is independent of $U_o$. This prediction is supported by classical $lg$ experiments [53] as well as recent $\mu g$ experiments [54] (Fig. 11). Note that $L_T(\mu g)/L_T(lg)$ is practically constant even beyond the transition to turbulent conditions (high $U_o$ and $Re$). Note also that the maximum $Re$ at which flame exists (the “blow-off” limit) is different at $lg$ and $\mu g$. This is surprising because blow-off conditions are typically controlled by behavior near the flame base [55], where buoyancy effects are often considered insignificant. This suggests that blow-off is partially affected by convection induced by the buoyant plume far above the jet exit, even at very high $U_o$. One would intuitively conclude that the $lg$ flames should blow-off at lower $U_o$ because buoyant flow would induce higher “effective” $U_o$ which is consistent with Fig. 11. This shows that buoyancy effects are quite ubiquitous even under conditions commonly thought to be unaffected by buoyancy.
Soot Formation Processes

Compared with $lg$ flames $\mu g$ gas-jet flames have much greater tendencies to emit soot [49,56], indicating that increases in $t_{\text{ext}}$ (thus, greater time for soot formation) plus broader regions in which composition and temperature are favorable for soot formation [56,57] outweigh lower temperatures at $\mu g$, which decreases soot formation [58]. Recent quantitative measurements [56] show peak soot-volume fractions about twice as high at $\mu g$ than $lg$ for 50% $C_2H_6/50% N_2$-air flames.

Surprisingly, $\mu g$ gas-jet flames exhibit "smoke points," corresponding to critical $U_c$ below which soot is consumed within the flame, and above which soot is emitted from the flame [57]. Smoke points are expected for buoyant round-jet flames because $t_{\text{sm}} \sim U_c^{-3/2}$; thus, increasing $U_c$ increases the time available for soot formation. But for nonbuoyant flames, $t_{\text{sm}} \sim U_c^2$, suggesting no smoke point should exist. Fully elliptic numerical computations [59] show that, for some circumstances, $t_{\text{sm}}$ does increase monotonically with $U_c$, which could explain smoke points for nonbuoyant flames. This behavior was suggested to result from axial diffusion effects [59], but in this case, $t_{\text{sm}}$ should asymptote to constant values for large $L_f$, where axial diffusion is negligible. Thus, simple explanations of $\mu g$ smoke points remain elusive. Residence-time considerations alone may be misleading; soot precursor temperature-composition-time history effects are apparently also important.

With weak convection, thermophoretic forces, which move particles toward lower temperatures, are an important effect. If convection and temperature gradients are in the same direction, the convective and thermophoretic forces may balance at some location. This leads to soot accumulation inside the flame front at $\mu g$ (Fig. 12) [60], where an annulus of soot accumulation is convected by axial flow through the flame tip. Downstream, for reasons not yet explained, the soot annulus fragments, creating crown-like structures. These effects are only observable at $\mu g$ where convection velocities are comparable to thermophoresis velocities, ($\sim 5$ mm/s for the conditions of Fig. 12 [60]).

Condensed-Phase Combustion

Droplet Combustion

The first microgravity combustion experiments were isolated fuel-droplet tests conducted by Kuma-gai and Isoda [61]. At $lg$, experimental measurements of droplet burning rates are compromised by buoyant convection, which destroys the spherical symmetry and inducing additional heat and mass transport, which alters burning rates and complicates modeling. Classical theory [62,63] predicts that burning rates for spherically symmetric buoyancy-free droplet burning are given by

$$d^3 = d_t^3 = Kt,$$

where $K = (8\lambda \rho_c C_p \ln(1 + B))$. In experimental studies, the droplet diameter $d_t$ could be fixed by forcing fuel through a porous sphere at rates that balances evaporation, leading to a steady mass burning rate ($m^\prime = (\pi/4) d_t^3 K$) but with most experiments employ fuel droplets where $d_t$ decreases with time. Somewhat surprisingly, many fuel-droplet results follow equation 8 well despite unsteadiness, heat losses, soot formation, and water absorption effects discussed below.

As with flame balls, steady solutions exist for spherical droplet flames even in infinite domains, with the flame front located at

$$d_t = d_t \ln(1 + B) / (ln(1 + f))$$

Whereas flame balls are convection-free, droplet flames exhibit Stefan flow due to fuel vaporization at the droplet surface. Mass conservation dictates that the Stefan velocity decays as $1/r^2$, causing temperature and concentration profiles to vary with radius in proportion to $(1 + B)^{-d/d^2}$ rather than $1/r$ as in flame balls, although these profiles, when normalized by flame radius, are indistinguishable at large $r$. Unlike flame balls, heat losses are not required for stable droplet flames because the flame location cannot move away from the stoichiometric contour.

The characteristic time scale for droplet combustion is $t_{\text{droop}} \sim d^2/\alpha$. This leads to two extinction limits [64,65], one for small $d_t$ where $t_{\text{droop}} < t_{\text{cond}}$ and thus fuel and oxidant cannot react before interdiffusing, and one for large $d_t$ where $t_{\text{droop}} > t_{\text{cond}}$ and thus the temperature decrease from radiative loss...
I requires - 100 s to reach steady state even in flame or occurs in flame balls [41], where radiative loss changing radiative loss over time. Analogous behavior into the far field, may be significant; this would cause increases of thermal energy and radiant combustion products, suggesting that nonsteady effects, decreases with increasing \( d_{\text{do}} \). The quasi-steady model (equation 8) predicts, but rather decrease in \( d_{\text{do}} \), apparently because of increased accumulation of soot and gas-phase radiant species for larger \( d_{\text{do}} \) [70].

renders the reaction rate sufficiently to cause extinction. The former occurs when \( d_f < \left( \frac{\text{art}_{\text{chem}}}{\text{art}_{\text{chem}}/2} \right) \) and the latter when \( d_f > \left( \frac{\text{art}_{\text{chem}}}{\text{art}_{\text{chem}}/2} \right) \). Recent space experiments have reported radiative extinction of large droplets in air [66] and \( \text{O}_2-\text{He} \) [67] atmospheres. Even when radiation does not cause extinction, it causes a decrease in \( K \), especially when soot formation is significant [68,69].

In experiments, \( K \) is not constant as the quasi-steady model (equation 8) predicts, but rather decreases with increasing \( d_{\text{do}} \) (Fig. 13) [70]. This suggests that nonsteady effects, specifically the diffusion of thermal energy and radiant combustion products into the far field, may be significant; this would cause changing radiative loss over time. Analogous behavior occurs in flame balls [41], where radiative loss requires \( \sim 100 \) s to reach steady state even in flame balls much smaller than typical droplet flames. Another indication of unsteadiness in droplet flames is that constant \( d_f/\text{art}_{\text{chem}} \) values are not generally achieved, especially for large droplets at \( \mu g \) [69,70], in contrast to the predictions of the quasi-steady theory. Unsteadiness effects in droplet flames were recently analyzed by King [71]. The effect of \( d_{\text{do}} \) on \( K \) has also been proposed [68] to result from soot accumulation, which is more significant for larger \( d_{\text{do}} \), and acts to decrease net heat release and increase radiation. As with non-premixed gas-jet flames, soot particles in \( \mu g \) droplet flames exhibit thermophoresis effects [72,73], leading to soot agglomeration between the droplet and flame front (Fig. 14). The agglomerates may break apart suddenly, leading to multiple burning fragments. Because velocity and temperature gradients are readily modeled in spherical droplet flames, such experimental observations enable assessment of thermophoresis effects on soot particle transport.

Another complicating factor arises in fuels that are miscible in water. The fuel may absorb water vapor from the combustion products, causing significant departure from equation 8 and reducing flammability. Such behavior is found [74,75] in methanol flames, where extinction diameters depend substantially on \( d_{\text{do}} \) due to water absorption during combustion.

Finally, in recent space experiments, flame oscillations with amplitudes comparable to the mean flame diameter have been observed [76]. The oscillation amplitude grew with time with extinction typically occurring after eight cycles. The oscillation frequency was approximately 1 Hz. Cheatham and Matalon [65] predicted oscillations of roughly the correct frequency in droplet flames of pure fuels under near-extinction conditions when the reactant Lewis numbers are sufficiently high. To date, oscillations have been reported for methanol/dodecanol (80/20 mass fraction) bicomponent droplets burning in air; it is uncertain whether oscillations also occur in pure fuels. Moreover, \( \mu g \) experiments in \( \text{O}_2-\text{He} \) atmospheres [67], with much higher \( Le \), did not exhibit oscillations. One possible explanation is that droplet support fibers were used in Ref. [76] but not Ref. [67]. The fiber could increase conductive and radiative losses, which encourage oscillations [65]. Another explanation is that, because oscillations occur only near extinction conditions, depletion of oxygen (leading to extinction) was much more significant in the smaller combustion chamber used in the methanol/dodecanol experiments. Such oscillatory instabilities are discussed further in the following section.

**Candle Flames**

An excellent example of the differences between \( \text{lg} \) and \( \mu g \) flames is seen in perhaps the most common and familiar of all combustion processes—candle flames. At \( \text{lg} \), candle flames are supported by air
entirely due to buoyant flow, which generates a self-sustaining flame and flow configuration. Obviously, this mechanism cannot apply at \( \mu \text{g} \). The spherical diffusion equation admits steady solutions for flame balls and droplet flames without forced convection. An interesting question is whether candle flames, which are not strictly spherical, can behave similarly. Space experiments indicate that candle flames can be steady for >45 min, with flame shapes typically hemispherical (Fig. 15). Eventually, the flames always extinguished, whereas the spherical flame model predicts that the flame would burn indefinitely. Protective screens used in the experiments may have limited the \( \text{O}_2 \) supply, eventually allowing sufficient \( \text{O}_2 \) depletion to cause extinction.

Before extinction, the candle-flame edge frequently advanced and retreated periodically. The oscillation amplitude increased over time, and on one retreating cycle, the entire flame extinguished. With larger wicks (thus, larger flame diameters), oscillations started spontaneously, whereas with smaller wicks, oscillations occurred only when solid objects were placed near the flame. Only a few cycles before extinction were observed in the Spacelab experiments, whereas hundreds of cycles were observed in the Mir experiments. This is probably because the protective screen was much more permeable in the Mir experiments, thereby decreasing the \( \text{O}_2 \) depletion rate and maintaining the flame at near-extinction conditions much longer.

At least two possible explanations for these oscillations have been advanced. Cheatham and Matalon [65] showed that, near extinction, oscillatory instabilities occur in spherically symmetric droplet flames with radiative loss at sufficiently high Lewis numbers. Their predicted oscillation frequencies (0.7–1.4 Hz) are comparable to experimental observations, however, the differences between spherically symmetric droplet flames and roughly hemispherical candle flames were noted [55]. Alternatively, Buckmaster [79] showed that the flame "edge" separating burning and nonburning regions of non-premixed flames exhibit oscillatory behavior at (for quasi-stationary edges) \( \text{Le} > 1 + \sqrt[3]{\beta (1 - (T_f / T))} \approx 2 \) when \( \text{Le} \) for the other reactant is unity. \( \text{Le} \approx 1 \) for \( \text{O}_2 \) in \( \text{N}_2 \), but \( \text{Le} \) for fuel vapors is probably closer to 2. Thus, edge-flame instabilities could explain the observed oscillations. While neither instability mechanism has been definitively linked to the candle-flame experiments, both predict greater propensities for oscillation with greater heat losses, which is consistent with the observation that oscillations occur near extinction.

**Flame Spread over Solid Fuel Beds**

Flame spread over solid fuel beds is typically classified as opposed-flow, where convection opposes flame propagation, or concurrent-flow. Because upward buoyant flow opposes flame spread, \( 1\text{g} \) downward flame spread is opposed-flow, whereas upward flame spread is concurrent-flow. At \( \mu \text{g} \) without forced flow, flame spread is always opposed-flow because the flame spreads toward the fresh oxidant with a self-induced velocity equal to the spread rate (\( S_f \)). At \( 1\text{g} \), self-induced convection is negligible because buoyancy-induced flows are typically \( (\rho \text{g}_{\text{d}})^{1/3} \approx 0.10 \text{ m/s} \) \( \gg S_f \). Very few concurrent-flow flame spread have been conducted at \( \mu \text{g} \) [80]; consequently, this section focuses on opposed-flow spread.

\( S_f \) is estimated by equating the conductive heat flux to the fuel bed \(\left(= \lambda(T_f - T_*)\delta \right)\) to the thermal transport zone thickness and \( U \) is the opposed-flow velocity (forced, buoyant and/or self-induced). To rate of fuel bed enthalpy increase \(\left(= \rho C_p \tau(T_f - T_*)\right)\). Assuming mixing-limited reaction (infinite-rate chemistry), for thermally thin fuels, where heat conduction through the solid is negligible, \( S_f \) is predicted to be [81,82]

\[
S_f = \frac{\pi}{4} \frac{\lambda}{\rho C_p \tau} \frac{T_f - T_*}{T_f - T_*} \tag{9}
\]

Note that \( S_f \) is independent of \( U \) and \( P \). For thermally thick fuels, where heat conduction through the solid fuel dominates, \( \tau \) is the thermal penetration depth into the solid, estimated by equating the conductive heat flux to the fuel bed to the heat flux through the fuel \(\left(= \lambda(T_f - T_*)/\rho C_p \tau \right)\). Assuming the subscripts \( f \) refers to the direction normal to the fuel surface. This leads to the exact solution [81]

\[
\tau = \frac{\lambda(T_f - T_*)}{\lambda(T_f - T_*)} \frac{\lambda(T_f - T_*)}{\lambda(T_f - T_*)} \frac{\alpha}{\lambda(T_f - T_*)} \frac{1}{\tau} \frac{1}{\lambda(T_f - T_*)} \frac{1}{\rho C_p} \frac{1}{\rho C_p} \frac{1}{\rho C_p} \frac{1}{\rho C_p} \tag{10}
\]

Note that, unlike the thin-fuel case, for thick fuels \( S_f \sim U^3 P^3 \).
FIG. 16. Minimum mole percent O\textsubscript{2} in N\textsubscript{2} supporting flame spread over a thin solid fuel bed, as a function of the opposed flow velocity (U) [83], showing dual-limit behavior, that is, residence-time limited extinction at high U and radiative loss extinction at low U.

Dual-limit extinction behavior is observed in \textmu g flame-spread experiments (Fig. 16) [83]. The time for thermal energy to diffuse across the convection-diffusion zone (t\textsubscript{diff}) is \( \delta U = \alpha U^2 \), thus, high-U extinction occurs when \( t_{\text{diff}} < t_{\text{chem}} \) or \( U > (\alpha t_{\text{chem}})^{1/2} \) and radiative extinction occurs when \( t_{\text{diff}} > t_{\text{rad}} \) or \( U < (\alpha t_{\text{rad}})^{1/2} \). (Surface radiative loss may also be important, particularly at moderate and higher U [84].) Interestingly, the minimum O\textsubscript{3} concentration supporting combustion (X\textsubscript{O\textsubscript{2},lim}), and thus the greatest hazard, corresponds to \( U \sim 0.1 \text{ m/s} \), which is lower than buoyant convection at 1g, and might correspond to ventilation drafts in manned spacecraft.

A radiative loss parameter can be defined as \( H = t_{\text{diff}}/t_{\text{rad}} = \alpha U^2 / t_{\text{rad}} \). Because \( H \sim U^{-2} \), \( S_f \) is lower at \( \mu g \), for \( U \) is lower and thus \( H \) is higher. Experiments [85] (Fig. 17) show that an imposed forced flow at \( \mu g \) increases \( S_f \); because \( U \) (sum of forced flow and \( S_f \)) increases and thus \( H \) decreases, whereas at higher \( U \) (whether buoyant or forced), \( S_f \) decreases as the high-\( U \) limit is approached. For 21\% O\textsubscript{2} or lower, the infinite-rate chemistry prediction of equation 9, \( S_f \sim U^n \), is never achieved. Only at 30\% O\textsubscript{2} is \( T_f \) high enough that this condition is achieved.

Because \( \alpha \sim P^{-1} \) and \( t_{\text{rad}} \sim P^4, H \sim P^{-4} \), thus for thin fuels, \( S_f \) should increase with \( P \) toward the ideal (adiabatic) value (equation 9). This is confirmed by quiescent thin-fuel space experiments [86,87], which show \( S_f \) increasing from 3.2 to 5.9 mm/s as \( P \) increases from 1.0 to 2.0 atm with fixed O\textsubscript{2} mole fraction (0.50). For these conditions, \( H \) decreases from 24 to 3.5, thus even at the highest \( P \), radiative effects are probably still important. This is consistent with computations [86] that predict \( S_f = 12 \text{ mm/s} \) (almost independent of \( P \)) for adiabatic conditions for this fuel/atmosphere combination.

Neither N\textsubscript{2} nor O\textsubscript{2} emit thermal radiation; thus, for flames in O\textsubscript{2}-N\textsubscript{2} atmospheres, only H\textsubscript{2}O and CO\textsubscript{2} combustion products radiate significantly. For these cases, typically \( a_p \sim 1.2 \text{ m} \gg \alpha \) and thus, radiative transfer is optically thin (negligible reabsorption). When \( a_p \gg 1 \), reabsorption effects cannot be neglected. With reabsorption, some radiation is not lost and may augment conduction to increase \( S_f \) above that without radiation. This behavior is seen experimentally [88] using strongly emitting/absorbing CO\textsubscript{2} and SF\textsubscript{6} diluents (Fig. 18), where \( S_f \) is higher and X\textsubscript{O\textsubscript{2},lim} is lower at \( \mu g \) than 1g, whereas the opposite (conventional) behavior is found in nonradiant diluents (not shown). These data indicate that, for nonradiating diluents, \( \mu g \) is less hazardous because X\textsubscript{O\textsubscript{2},lim} is higher at \( \mu g \) than at 1g (0.21 vs. 0.16

FIG. 17. Flame-spread rate over a thin solid fuel bed as a function of the opposed-flow velocity (U), for three values of the mole percent O\textsubscript{2} [85], showing dual-limit response. Note that the infinite-rate kinetics prediction [81,82] that the spread rate is independent of \( U \) is only satisfied at O\textsubscript{2} mole fractions higher than that in air.
unsteady solid-phase conduction is considered, \( \tau = (\alpha t)^{1/2} \), which results in \( S_f \sim t^{-1.2} \) [89]. Consequently, all fuel beds at quiescent \( \mu g \) conditions eventually become thermally thin (penetration depth greater than the bed thickness) unless the radiative effects, discussed later, are considered. Of course, flames may extinguish due to large \( \delta \) (thus large radiative loss) before reaching steady-state, thermally thin conditions.

A difficulty in comparing space experiments to two-dimensional model predictions is that the fuel bed width \( W \) (30 mm for thin fuels [86,87] and 6.2 mm for thick fuels [89]) is smaller than the thermal transport zone thickness \( \delta \). Consequently, these experiments can hardly be considered two-dimensional. Both lateral heat loss, which retards spread, and lateral \( O_2 \) influx, which enhances spread, are probably important, thus their effects may partially cancel. Some authors [86] suggest that radiative losses decrease \( \delta \) to values much smaller than \( \alpha / U \), but the oxygen transport zone thickness \( \delta_{O_2} \) is still \( D_{O_2}/U \), because no analog to radiative loss exists for \( O_2 \) transport. Because \( Le = \alpha / D_{O_2} \approx 1 \), \( \delta_{O_2} \) is nearly the same as \( \delta \) in adiabatic flames. Thus, the \( \mu g \) flames mentioned above have probably benefited substantially from lateral \( O_2 \) influx, especially for lower pressures and \( O_2 \) mole fractions, where \( \delta W / \alpha S_f W \) is largest. In fact, \( S_f \) might be higher at smaller \( W \) due to lateral \( O_2 \) influx. Space experiments using cylindrical fuel rods are planned [90] to examine truly two-dimensional spread.

Recently, a surprising observation of fingering fronts was found in space experiments using paper samples treated to inhibit flaming combustion but allow smoldering propagation (Fig. 19) [91]. Fingering was observed at \( \mu g \) when \( U < 50 \text{ m/s} \), whereas smooth fronts were observed at \( 1 g \) for all \( U \). This was proposed [91] to result from limited \( O_2 \) mass transport at \( \mu g \) with low \( U \), which caused the \( O_2 \) consumption regions to become localized spots instead of continuous fronts. This proposition does not explain why heat conduction does not smooth out potential fingers as it does (for example) in premixed-gas flames with \( Le > 1 \). The following alternative explanation is proposed here. Gas-phase heat transport occurs on the length scale \( \delta = \alpha / U \), and solid-phase transport occurs on the scale \( \delta_s = \alpha_s / \alpha \), where \( \alpha \) is the smolder front velocity and \( \alpha_s \) the solid thermal diffusivity. Oxygen transport occurs only through the gas phase on the scale \( D_{O_2}/U = \alpha / \delta \). Radiative loss can suppress heat transport through the gas, but no corresponding effect on \( O_2 \) transport can occur. Thus, at low \( U \), the effective \( Le \) is \( \alpha / D_{O_2} \ll 1 \). At higher \( U \) or at \( \mu g \), \( \delta \) is smaller, gas-phase heat transport dominates, and radiative effects are weaker; thus, the effective \( Le \) is \( \alpha / D_{O_2} \approx 1 \). These assertions are consistent with estimates [91].
Flame spread over liquid fuel pools encompasses practically all solid-fuel flame spread phenomena discussed above, plus liquid-phase flow effects. Typically, $\mu_g$ is smaller for liquid than solid fuels, thus $S_f$ is higher. Also, if $T_s - T_f$ is small, some fuel pre vaporization occurs even at $T = T_f$, thus partially premixed gas-phase combustion phenomena may occur. Because of the fuel surface temperature gradient upstream of the flame, surface tension gradients are produced that cause the surface layer to move upstream (away from the flame), which increases $S_f$. At $\mu_g$, this heated liquid layer must lie near the surface, whereas at $\mu_l$, no limitation exists.

Experiments at $\mu_l$, summarized in [95], show that at low fuel temperatures, the average $S_f$ is small (typically 10 mm/s) and the spread alternates between a fast "jump" velocity and a slow "crawl" velocity. At higher fuel temperatures, $S_f$ is faster and steady. For the conditions exhibiting pulsating spread at $\mu_l$, flame spread cannot be maintained, whereas for the conditions exhibiting uniform spread at $\mu_l$, steady spread is also exhibited at $\mu_l$ [96]. Pulsating spread has never been observed at $\mu_l$. No definitive explanation for these observations has been advanced. For the conditions exhibiting pulsating spread at $\mu_l$ and no spread at $\mu_l$, flame spread is still different at $\mu_l$ and when forced flows comparable to buoyancy-induced convection ($U = 0.30$ m/s) are imposed (Fig. 20) [97]. Specifically, $\mu_l$ spread is almost unaffected by the imposed flow, but $\mu_l$ spread is steady with $S_f$ being lower (15 mm/s) than either the $\mu_l$ jump velocity ($\approx 100$ mm/s) or crawl velocity ($\approx 22$ mm/s).

Detailed numerical modeling [98] predicts pulsating spread at $\mu_l$ for the conditions of Fig. 20 and values of $S_f$ much closer to the measured $\mu_l$$S_f$. Remarkably, if thermal expansion is artificially suppressed, good agreement between the model and $\mu_l$ experiments is found. It is proposed [97] that this agreement results from three-dimensional effects; specifically, in the experiment, flow induced by thermal expansion is relaxed in the lateral dimension, whereas the two-dimensional model does not permit this. That three-dimensional effects might dominate is surprising considering that, for this flame, $\alpha W = \alpha_{W'} = 0.02$, thus, $\delta \ll W$. Also, this hypothesis does not explain why pulsating flame spread is observed at $\mu_l$ but not $\mu_l$. Consequently, in the case of liquid-fuel flame spread, $\mu_l$ experiments have identified limitations in our current understanding of combustion processes at $\mu_l$.

**Recommendations for Future Studies**

**Reabsorption Effects**

The $\mu_l$ studies described here suggest new unresolved issues and opportunities for further improvements in understanding. Perhaps the most important are the effects of reabsorption of emitted radiation, including both reabsorption by the emitting gas and, in two-phase combustion, absorption by the condensed phase. All radiative effects discussed above are critically dependent on the degree of reabsorption. To study reabsorption effects requires radiatively active diluents (CO$_2$, SF$_6$), high pressures and/or large systems. All of these conditions lead to higher $G_{rd}$ at $\mu_l$ and thus turbulent flow. Hence, $\mu_l$
experiments enable study of reabsorption effects without the additional complications due to turbulence.

Reabsorption effects are important not only to \( \mu \) studies but also to combustion at high pressures and in large combustors. For example, at 40 atm, typical of premixed-charge internal combustion engines, \( \alpha_p \approx 18 \text{ m}^{-1} \), thus \( \alpha_p^{-1} \approx 0.045 \text{ m} \), for stoichiometric combustion products. Because this length scale is comparable to cylinder radii, reabsorption effects within the gas cannot readily be neglected. Simple estimates [99] indicate radiative loss may influence flame quenching by turbulence in lean mixtures. Similarly, reabsorption cannot be neglected in atmospheric-pressure furnaces larger than \( \alpha_p^{-1} \approx 2.2 \text{ m} \). Moreover, many combustion devices employ exhaust-gas or flue-gas recirculation; for such devices, the unburned mixtures contain significant amounts of absorbing \( \text{CO}_2 \) and \( \text{H}_2\text{O} \).

Although reabsorption of emitted radiation could affect practically all types of flames reviewed here, to date, reabsorption effects have been studied only for propagating premixed-gas flames [24,26], flame balls [45], and flame spread over thermally thin fuels [88]. All have shown substantial differences from optically thin behavior. Two examples of effects expected for other flames are given below.

Reabsorption effects could be substantially more important for droplet combustion than for flame spread over solid fuels because for droplets, the Stefan flow severely limits heat conduction to the droplet surface. This is why heat release (\( B \)) affects burning rates (\( K \)) only weakly (logarithmically) (equation 8). Radiative transfer is unaffected by the Stefan flow. Equation 8 is readily extended to include surface radiative flux (\( q_s \)).

\[
\Omega = \ln \left(1 + \frac{B}{1 - R/\Omega}ight), \quad R = \frac{q_s d C_p}{2 L_a}, \quad \Omega = \frac{K p d C_p}{\lambda}\]  

(11)

Figure 21 shows the predictions of equation 11. (Although equation 11 has apparently not been presented previously, numerical studies [100,101] have shown qualitatively similar predictions. Moreover, these studies show that typical radiative absorption lengths for liquid fuels at relevant wavelengths are on the order of 1 mm, thus large droplets could absorb most incident radiation.) For spherical shells of radiant combustion products having thickness \( \delta \ll d_p, q_s = \delta \lambda / 4 \), then for typical values \( B = 8.5, \lambda = 2 \times 10^6 \text{ W/m}^2, \delta = 10 \text{ mm}, d_p = 5 \text{ mm}, C_p = 1400 \text{ J/kgK}, \lambda = 0.07 \text{ W/mK}, \) and \( L_a = 400 \text{ kJ/kg} \), equation 11 predicts \( R = 0.63 \) and \( \Omega = 1.11 \); thus, moderate effects of radiative transfer are expected. For droplets in radiatively active diluents such as \( \text{CO}_2 \), the effect could be much stronger. Using the P1 approximation, for a sphere of unit emissivity in an infinite gray gas, \( q_s = [4/(2 + 3 \alpha_p)]\sigma(T^4 - T^4) \) [102]. Using volume-averaged properties \( T = 1000 \text{ K} \) and \( \alpha_p(11) = 20 \text{ m}^{-1} \), \( R = 18 \) and thus \( \Omega = 8.0 \), indicating radiation dominates heat transport. As discussed later, at high pressures radiative effects may prevail even in \( \text{O}_2-N_2 \) atmospheres.

Flame spread over thermally thick fuel beds in quiescent atmospheres at \( \mu \) is typically considered inherently unstable [59], however radiative transfer to the fuel bed could enable steady spread. If the flame is modeled by an isothermal volume with dimensions \( \delta \times \delta \times W \), radiation induces a radiative flux \( A \delta^2 W = \lambda (\alpha_s^2 S) \delta W \) that augments the conductive flux \( \lambda (\delta W)/(T_f - T) \delta \). Equating this total heat transfer to \( \rho C_f \delta (T_f - T) \delta W \) yields

\[
S_f = \left[ \frac{\lambda A \delta^2}{\alpha_s^2 S_0} (T_f - T) - \lambda \frac{\delta W}{\delta T} (T_f - T) \right]^{1/2}
\]

(12)

which vanishes without gas radiation (\( \lambda = 0 \)). Thick-fuel space experiments in \( \text{O}_2-\text{CO}_2 \) or \( \text{O}_2-\text{SF}_6 \) atmospheres could be employed to check for steady spread and test the accuracy of equation 12.

Whereas optically thin radiation modeling is a reasonably straightforward, modeling of spectrally dependent emission and absorption is challenging because local fluxes depend on the entire radiation field, not just local scalar properties and gradients. Some relevant computations have employed gray-gas models [74], but recent studies [26,103] show that these methods are probably inaccurate because of the wide variation in spectral absorption coefficient with temperature, species, and wavelength. Comparisons of various radiative treatments for small one-dimensional non-premixed flames has been made [103]. Comparisons for larger, multidimensional systems would be valuable. Moreover, recent studies of \( \mu \) soot formation [52,56,57] may enable improved modeling of soot radiation at \( \mu \).

**High-Pressure Combustion**

All practical combustion engines operate at pressures much higher than atmospheric. The impact of buoyancy for premixed-gas flames scales as \( n = v_{\text{avg}} / (\rho g S) \delta^2 \delta^2 \), which is the overall reaction order \( (S_1 = \rho g S \delta^2 \delta^2) \). Because, typically, \( n < 4/3 \) for weak mixtures [104], where buoyancy effects are most important, the impact of buoyancy increases with pressure. Also, as discussed earlier, radiation effects are more difficult to assess at higher pressure due to increased interference from buoyant transport. Nevertheless, few high-pressure \( \mu \) combustion experiments have been performed. High-pressure droplet combustion experiments [105] revealed substantial but different increases in \( K \) with \( P \) at 1 atm and \( \mu \). Radiative effects were not discussed but could have been important, because in equation 11,
Imerous investigations of round-jet flames at \( q_g \) and residence times (t) for buoyant and non-buoyant gas-jet flames can neither predict the number of flame balls produced from an ignition source be predicted. Modeling using three-dimensional effects (Fig. 5), but the transition from splitting into the two-dimensional model to account for lateral heat and mass transport and chamber sizes employed to date; near-wall species for lateral heat losses and lateral diffusion of each fuel bed, into the two-dimensional model to account for the coordinates parallel and perpendicular to the coordinate system.

Three-Dimensional Effects

In earlier sections, effects of lateral heat and mass transport on flame spread were discussed. To assess three-dimensional effects, \( \mu \)g experiments with varying fuel bed width (W) are needed. Complementary three-dimensional modeling using codes such as those developed by NIST [106], extended to include gas-phase radiation, would be instructive. An approximate but much less expensive approach would be to incorporate volumetric terms \( 6 \rho (T(x, y) - T_a) / W^2 \) and \( 6 \rho \phi (Y(x, y) - Y_a) / W^2 \), where \( x \) and \( y \) are the coordinates parallel and perpendicular to the fuel bed, into the two-dimensional model to account for lateral heat losses and lateral diffusion of each species.

Another three-dimensional effect is found in the development of flame balls from ignition kernels. Currently, it is known that large flame balls are linearly unstable to three-dimensional disturbances for weak loss (Fig. 5), but the transition from splitting flame balls to stable flames is not well understood, nor can the number of flame balls produced from an ignition source be predicted. Modeling using three-dimensional premixed flame codes [107] is needed.

Gas-Jet Flames

Table 2 shows predicted scalings of flame lengths (\( L_f \)) and residence times (t) for buoyant and non-buoyant round-jet and slot-jet flames. Despite numerous investigations of round-jet flames at \( \mu \)g, no \( \mu \)g slot-jet results are available to test those predictions. Currently, it is unknown whether slot-jet flames at \( \mu \)g would exhibit smoke points or whether this information could be used to explain smoke points in \( \mu \)g round-jet flames. Because \( L_f \) depends on \( g \) for buoyant slot-jet but not round-jet flames, \( L_f \) should be quite different at \( 1g \) and \( \mu \)g for slot-jets but not round-jets. Residual accelerations in aircraft \( \mu \)g experiments will be more problematic for slot-jet than round-jet flames because \( t_{\text{ext}} \sim g^{-1/2} \) for buoyant round-jet flames whereas \( t_{\text{ext}} \sim g^{-1} \) for slot-jet flames.

There has been little investigation of blow-off behavior of laminar gas-jet flames at \( \mu \)g. Dual-limit behavior might occur for flames of fixed-mass flow rate but varying \( d_o \), with short residence time extinction at small \( d_o \), \( t_{\text{ext}} \sim d_o / U_o \) and radiative extinction at large \( d_o \), (thus, large \( t_{\text{ext}} \)). Experiments should be conducted by diluting the fuel rather than increasing \( U_o \) to obtain blow-off without transition to turbulence. In this way, dual-limit behavior has been observed at \( 1g \) [92] with short residence time and conductive loss (to the burner rim) extinction branches.

Quasi-Steady Spherical Diffusion Flames

As discussed previously, comparing predicted radiative extinction limits of droplet flames to experiments is problematic because quasi-steady conditions may not be obtained, because extinction occurs for sufficiently large droplet and flame diameters, but the droplet diameter decreases throughout its life. Numerical models can account for transient effects, but the multidimensional ignition process is difficult to model quantitatively. Comparisons of droplet experiments and computations to corresponding results obtained using fuel-wetted porous spheres would be most interesting. The fuel should be forced through the porous sphere at slowly increasing rates until extinction occurs when \( d_o > (\alpha_{\text{rad}})^{1/3} \), thus obtaining truly quasi-steady extinction. Long \( \mu \)g durations would be required to establish steady diffusion-dominated far-field temperature and composition profiles (thus, steady radiative loss). Some preliminary results have been obtained in aircraft experiments [108] but were severely compromised by the residual \( g \) levels. Candle flames are similar to wetted porous spheres, though without true spherical symmetry or any means to control or measure the instantaneous \( d_o \).

A simpler related experiment employs porous spheres through which gaseous fuel is forced at prescribed \( m \), resulting in a flame diameter \( d_f \sim m^P / 2\pi \phi (1 + f) \). Some drop-tower experiments using this configuration have been reported [109,110], although steady-state conditions were not obtained due to short \( \mu \)g durations. Steady-state behavior appears unlikely given the \( m \) (thus, air consumption rates) and chamber sizes employed to date; near-wall
The discrepancies result largely from differences in mixtures away from extinction limits very faithfully. These chemical models predict lower [114], these chemical models predict smaller balls, balls [45] and lg strained premixed H2-air flames losses could explain. In contrast, for H2-air flame than experimental uncertainty or unaccounted heat observations [22,33,113]. The discrepancy seems larger and caner flammability limits than experimental ob-
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CatalyticCombustion

Catalytic combustion holds promise for reduced emissions and improved fuel efficiency in many combustion systems [111,112]. Because catalysis occurs at surfaces, it is inherently multidimensional and/or unsteady, requiring reactant transport to the surface and heat and products transport from the surface. Although boundary-layer approximations can be invoked, probably the only truly one-dimensional steady catalytic configuration is a spherical surface immersed in nonbuoyant quiescent premixed gas—a "catalytic flame ball." In this case, \( r_s \) is fixed but the surface temperature \( T_s \), and fuel concentration \( Y_F \) are unknown. These are related through the energy conservation (including surface radiation) and diffusion equations to obtain the surface reaction rate in moles per second (\( \dot{\Theta} \)):

\[
\dot{\Theta}(T_s, Y_F) = \rho D \gamma_F \left( (1 - Y_F) / \epsilon \right) \gamma_F \frac{Y_F}{Y_F} = 1 - \frac{T_s - T_0}{T_m - T_0} \left( 1 + \frac{\sigma \epsilon \gamma_F (T_s^4 - T_0^4)}{\lambda (T_s - T_0)} \right)
\]

where \( \epsilon \) is the surface emissivity and properties with the subscript \( s \) are evaluated at \( T = T_s \). By varying \( r_s, Y_F, \) pressure, and diluent gas, \( \dot{\Theta}(T, Y) \) can be inferred from equation 13 and the measured \( T_s \). Of course, conditions must be unfavorable for initiation of propagating flames or flame balls that stand off from the surface.

Chemical Models

An important contribution of \( \mu \) combustion experiments has been improved understanding of extinction processes, which are inherently related to finite-rate chemistry. To obtain closure between experiments and computations, accurate chemical models are needed. For lean premixed hydrocarbon-air flames, most models [21,26] predict higher \( S_L \) and leaner flammability limits than experimental observations [22,33,113]. The discrepancy seems larger than experimental uncertainty or unaccounted heat losses could explain. In contrast, for H2-air flame balls [45] and lg strained premixed H2-air flames [114], these chemical models predict smaller balls, lower \( S_L \), and richer flammability limits than experiments. These chemical models predict \( S_F \) in mixtures away from extinction limits very faithfully. The discrepancies result largely from differences in rates for H + O2 + M reactions, particularly the Chaperon efficiencies of various M species [45]. These reactions are extremely important in near-limit flames due to competition between chain-branching and chain-inhibiting steps near limits [21] but are much less important away from limits. Further scrutiny of the rates for these reactions at intermediate temperatures (1100–1400 K) would be welcomed.

Conclusions

Our understanding of combustion fundamentals has been broadened by \( \mu \) experiments into regimes not previously explored. In particular, these experiments have helped integrate radiation into flame theory. Although flame radiation has long been recognized as an important heat-transfer mechanism in large fires [115], its treatment has largely been ad hoc because of the difficulty of predicting soot formation. Also, large-scale fires at Ig are inevitably turbulent, leading to complicated flame-flow interactions. Small-scale \( \mu \) flames are laminar, often soot-free, and have significant influences of radiation. As a result of radiation effects, both premixed and non-premixed flames have exhibited dual-limit extinction behavior, with residence time-limited extinction at high strain, or curvature and radiative loss-induced extinction at low strain or curvature. The high-strain limit is readily observed at Ig; when forced flow is absent, buoyant flow causes this strain. For weak mixtures, these limits converge, but the convergence and the entire low-strain extinction branch can only be seen at \( \mu \). This dual-limit behavior is observed for stretched and curved premixed-gas flames, strained non-premixed flames, isolated fuel droplets, and flame spread over solid fuels. Besides radiative effects, \( \mu \) studies have enabled observation and clarification of numerous other phenomena, for example, thermophoresis effects in soot formation, spherically symmetric droplet burning, diffusion-controlled premixed flames (flame balls), and flame instabilities in droplets and candle flames. Considering the rapid progress made recently, further advances are certain to occur. It is hoped that this report on the current state of understanding will help motivate and inspire such advances.

Nomenclature

- \( \alpha_p \): Planck mean absorption coefficient
- \( A \): flame surface area
- \( B \): transfer number (chemical enthalpy generation/enthalpy needed for fuel evaporation)
- \( C_p \): stoichiometric molar ratio of fuel to air constant-pressure heat capacity
INVITED PLENARY LECTURE

characteristic flow-length scale or tube diameter

droplet flame diameter
droplet diameter
droplet initial diameter
jet exit diameter (round jets) or slot width
(slot jets)

mass diffusivity

overall activation energy of the heat-release reactions

stoichiometric fuel to air mass ratio

acceleration of gravity

Grashof number based on characteristic length scale (d) = \( gd^2/\nu^2 \)

heat transfer coefficient in a cylindrical tube = \( 16\beta/d^2 \)

radiative loss parameter for flame spread

\[ t_{\text{flam}}/t_{\text{rad}} = \alpha/\gamma_{\text{rad}} \]

droplet burning rate constant (equation 8)

latent heat of vaporization of liquid fuel

Lewis number (\( \alpha CD = \text{thermal diffusivity}/\text{reactant mass diffusivity} \))

mass burning rate

fuel molecular weight

pressure

radial coordinate

flame radius

scaled flame radius (eq 6); radiation parameter (eq 11)

gas constant

jet Reynolds number = \( Uod/\nu \)

flame spread rate over solid fuel bed

remixed laminar burning velocity

burning velocity at the flammability limit

Schmidt number = \( v/D \)

chemical time scale

droplet flame time scale = \( d^2/\nu \)

inviscid buoyant transport time scale

residence time of non-premixed jet flame

radiative loss time scale

viscous buoyant transport time scale

temperature

adiabatic flame temperature

inertia density

convective velocity

local axial velocity in counterflow configuration

jet exit velocity

gas-jet flame width

solid fuel bed width

axial coordinate

fuel mass fraction

thermal diffusivity

nondimensional activation energy = \( \frac{E}{RT_f} \)

minimum oxygen concentration supporting combustion

flame thickness

gas specific heat ratio

thermal conductivity

radiative heat loss per unit volume = \( -4\pi\alpha(T_1 - T_2) \)

kinematic viscosity

density

fuel bed half-thickness (thin fuel) or thermal penetration depth (thick fuel)

Stefan-Boltzmann constant

Subscripts

d droplet surface condition

f flame front condition

s solid-fuel or solid-surface condition

v solid- or liquid-fuel vaporization condition

\( \infty \) ambient conditions

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REFERENCES


H. S. Mukunda, Indian Institute of Science, Bangalore, India. In your presentation, you credited the computational demonstration of the result on the absence of a flammability limit in adiabatic flames to Giovanangli and Smooke [1]. This result was first obtained in 1988 [2]. You might like to amend the citation suitably.

REFERENCES


Author's Reply. In this earlier paper by the questioner and his collaborators [2], it was shown that the solutions of the unsteady planar one-dimensional adiabatic premixed flame equations do not predict any flammability limit for lean CH4-O2-N2 mixtures—the burning velocity decreases asymptotically to zero as the mixture strength is decreased. In Smooke and Giovanangli (Ref. [1] of the comment), the steady version of these equations was solved for CH4-air and H2-air mixtures, and a similar conclusion was reached. Both are important findings, because it is possible in principle that unsteady effects could suppress limits of steady flames or cause limits to occur for mixtures that are flammable as steady flames. Together they show that one must identify loss mechanisms of the appropriate magnitude to explain flammability limits, which was one of the key points of my presentation.

J. P. Gore, Purdue University, USA. You seem to have drawn a conclusion based on your time-scale analysis that “radiation is not important at 1 g.” For combustion under high-pressure conditions, for fires (large and small), and for all flames as far as high activation energy reaction steps are involved, radiation can be potentially important. Could you please re-examine and clarify your conclusion regarding radiation not being important at 1 g?

Author's Reply. Certainly radiation is an important mechanism of heat transfer in flames even at 1 g—large
Industrial furnaces and boilers depend almost entirely on radiative transfer. The same is true of many large accidental fires (e.g., in buildings, over liquid fuel pools). The message I intended to convey was that the radiation transport time scale is longer than the buoyant transport time scale at earth gravity, and thus radiation does not play an important role on local flame structure and extinction conditions at earth gravity compared to the role that buoyant convection plays. Even on the large scale, however, buoyancy is probably a more dominant transport mechanism at 1 g—while a building fire would be very different with buoyancy but no radiation, imagine how different it would be with radiation but no buoyancy.