Turbulent Clustering of Particles and Radiation-Induced Mechanism of Dust Explosions

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Abstract

We show that clustering of dust particles in a turbulent flow ahead of the primary flame gives rise to a significant increase of the radiation absorption length. This ensures that clusters of dust particles even far ahead of the flame are sufficiently exposed and heated by the radiation from the primary flame to make them multiple ignition kernels capable to ignite the secondary explosion in a large volume of fuel-air mixtures. The multiple radiation-induced ignitions of fuel-air mixtures by ignition kernels in a large volume ahead of the primary flame increase efficiently the total flame area, resulting in a strong increase of the effective combustion speed, defined as the rate of reactant consumption of a given volume. The extent, to which the rate of combustion and overpressures increase, depends on the parameters of the turbulent flow ahead of the flame. It is shown that the radiation-induced multi-point ignitions of the secondary explosion give rise to the high rate of combustion and overpressures of the order required to account for the observed level of damages in unconfined dust explosions, e.g. such as the 2005 Buncefield vapor-cloud explosion.

1. Introduction

Dust explosions occur when an accidentally ignited flame propagates through a cloud of fine particles suspended in gaseous fuel-air mixtures [1-3]. Dust explosions have been significant hazards for centuries in the mining industry and in grain elevators. Currently the danger of dust explosions is a permanent threat in all those industries in which powders of fine particles are involved. Despite intense investigations over more than 100 years the mechanism of dust explosions still remains one of the main unresolved problem. It is well established that unconfined dust explosions consist of a relatively weak primary explosion followed by much more severe secondary explosions. While the hazardous effect of the primary explosion is relatively small, the secondary explosions may propagate with a speed of up to 1000 m/s producing overpressures of

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over 8-10 atm, which is comparable to the pressures produced by a detonation. However, analysis of the level of damages indicates that a detonation is not involved, while normal deflagrations are not capable of producing such high velocities and observed overpressures and damages. Of special interest is the 2005 Buncefield fuel storage depot explosion. In this particular case investigators and forensics teams were able to collect a large amount of data and evidences, providing a unique valuable information about the timings and damage data in the event [4]. The consequent analysis of damages, combustion rates and associated tests [5, 6] was performed in attempt to understand possible mechanisms leading to the unusually high rate of combustion. Among different mechanisms, the hypothesis proposed earlier by Moore and Weinberg [7] that the anomalously high rate of the flame propagation in dust explosions can be through the radiative ignition of large fibrous particles a few millimeters size heated and ignited ahead of the flame front, attracted much attention.

In normal practice emissivity of combustion products and radiation absorption in a fresh unburnt gaseous mixture are small and do not influence the flame propagation. The situation is drastically changed for flames propagating through a cloud of fine particles suspended in a gas mixture. Microns size dust particles suspended in the gas mixture ahead of the flame efficiently absorb radiation emitted by the flame with heat being transferred to the surrounding gas by thermal conduction, so that the gas temperature lags behind that of the particle. For the uniformly distributed dust particles with a typical dust cloud mass density, $\bar{\rho}_d = m_p \bar{N} = (0.01 - 0.03)\,\mathrm{kg/m^3}$, the radiation absorption length, $L_a = 1/\bar{\kappa} \approx (\rho_p / \rho_d) r_p$, is of the order of a few centimeters. Here $\bar{\kappa}$ is the radiation absorption coefficient of particles, ρ_p, m_p, $\bar{N}$, and r_p are the particle material density, mass, mean number density, and radius, correspondingly. If dust particles are evenly dispersed, the radiation emanated from the flame do not increase substantially the rate at which the flame propagates, because the flame consumes the heated unburned fuel before the fuel temperature will have risen up appreciably, so that radiation cannot become a dominant process of the heat transfer [8]. On the contrary, if there is a transparent gap between the optically thick particle layer (or a particle clump) and the flame, then the layer ahead of the flame front can be sufficiently heated by the flame-emanated radiation to ignite new combustion modes in the surrounding fuel-air mixture [8].

In turbulent flows ahead of the primary flame ($Re \approx 10^4 - 10^5$) dust particles assemble in small clusters. The turbulent eddies, acting as small centrifuges, push particles to the regions between eddies, where the pressure fluctuations are maximum and the vorticity intensity is minimum. This effect is known as inertial clustering ([9] and references therein). Recent analytical [10] and experimental [11] studies have shown that clustering of particles is much more effective in the presence of a mean temperature gradient. In this case correlations between fluctuations of fluid temperature and velocity, and, therefore, correlations between fluctuations of pressure and fluid velocity, cause additional pressure fluctuations. As a result, the particle concentration in clusters may rise by a few orders of magnitude compared to the mean particle concentrations of the evenly dispersed particles [10, 12].

We show that the mechanism of the secondary explosion is the multiple ignitions of the large volume of te fuel-air mixture ahead of the primary flame by the ignition kernels formed of the optically thick particle clusters, which are exposed and heated by the radiation emanated from the primary flame. We show that due to the strong increase of the radiation penetration length caused by the particle clustering, the particle clusters are sufficiently exposed and heated by the radiation in a large volume ahead of the flame. The increase of the radiation penetration length due to an inhomogeneous spatial distribution of the density concentrations has been studied analytically [13, 14], and the radiative heat transfer in a particle-laden flow was studied using Monte Carlo modeling taking into account the inertial particle clustering [15, 16].

In order for the secondary explosion followed by a shock wave occurred, the pressure, produced by the mixture ignited by the radiatively heated clusters ahead of the flame, must rise faster than it can be equalized by sound waves. This means that the penetration length of radiation, L_{rad}, in the dust cloud must be sufficiently large to ensure that the clusters of particles even far ahead of the flame are sufficiently exposed and heated by the radiation, to become ignition kernels, e.g. $L_{rad} >> \tau_{ign} c_s$, where c_s is the speed of sound in the mixture ahead of the flame, and τ_{ign} is the characteristic time-scale of fuel-air ignition by the radiatively heated particle clusters. The severity of the secondary explosion depends on the strength of the inequality $L_{rad} >> \tau_{ign} c_s$, which in its turn depends on turbulence parameters in the flow ahead of the flame.

2. The radiative heat transfer. Mean-field approach

We assume that the particle absorption coefficient is much larger than the scattering coefficient and that the gas phase is transparent for the radiation. Then, the radiative transfer equation for the radiation intensity $I(\mathbf{s},\mathbf{r})$ is reduced to [17]:

$$(\mathbf{s}\cdot\nabla)I(\mathbf{s},\mathbf{r}) = \kappa(\mathbf{r})\left(I_b(\mathbf{r}) - I(\mathbf{s},\mathbf{r})\right) \tag{1}$$

where $I_b(\mathbf{r})$ is the black body radiation.

For obtaining the effective radiation absorption coefficient for the condition of turbulent particle clustering we apply a mean-field approach. All quantities are decomposed into the mean and fluctuating parts: $I = \bar{I} + I'$, $I_b = \bar{I}_b + I'_b$, $\kappa = \bar{\kappa} + \kappa'$. The fluctuating parts have zero mean values, and overbars denote averaging over an ensemble of fluctuations. The instantaneous particle number density is $n = \bar{N} + n'$, where $\bar{N}$ and n' are the mean particle number density and fluctuations of number density. The instantaneous absorption coefficient is $\kappa = \sigma_a n$, and the mean particle absorption coefficient is $\bar{\kappa} = \sigma_a \bar{N}$, so that the fluctuations of absorption coefficient is $\kappa' = n'\sigma_a = n'\bar{\kappa}/\bar{N}$. By averaging Eq.(1) over the ensemble of the particle number density fluctuations and assuming that the characteristic correlation scales for the fluctuations are much smaller than the scales of the mean field variations, we obtain the mean-field equation for the mean irradiation intensity

$$(\mathbf{s}\cdot\nabla)\bar{I}(\mathbf{s},\mathbf{r}) = \bar{\kappa}\left(\bar{I}_b - \bar{I}\right) + \langle \kappa' I'_b \rangle - \langle \kappa' I' \rangle \tag{2}$$

The equation for fluctuations I' is:

$$(\mathbf{s}\cdot\nabla + \bar{\kappa} + \kappa')I'(\mathbf{r},\mathbf{s}) = I_{source}\,. \tag{3}$$

The solution of Eq. (3) is

$$I'(\mathbf{r},\hat{\mathbf{s}}) = \int_{-\infty}^{\infty} I_{source} \exp\left[-\left|\int_{s'}^{s}\left[\bar{\kappa} + \kappa'(s'')\right]ds''\right|\right]ds',. \tag{4}$$

where $s = \mathbf{r}\cdot\hat{\mathbf{s}}$ and

$$I_{source} = -\kappa'\left(\bar{I} - \bar{I}_b\right) + \langle \kappa' I' \rangle - \langle \kappa' I'_b \rangle + (\bar{\kappa} + \kappa')I'_b\,. \tag{5}$$

Expanding exponent in (4) in Taylor series, multiplying Eq. (4) by κ' and averaging over the ensemble of fluctuations, we obtain equation for the one-point correlation function $\langle \kappa' I' \rangle$:

$$\langle \kappa' I' \rangle = -\bar{\kappa}\left(\bar{I} - \bar{I}_b\right)\frac{2\beta J_1}{1 + 2\beta J_2}\,. \tag{6}$$

where

$$J_1 = \int_0^{+\infty} \Phi(Z)\exp(-\beta Z)\,dZ\,, \tag{7}$$

$$J_2 = \beta\int_0^{+\infty} \left(\int_0^Z \Phi(Z')dZ'\right)\exp(-\beta Z)\,dZ\,, \tag{8}$$

Here $\Phi(\mathbf{R}) = \langle n'(\mathbf{x})n'(\mathbf{x}+\mathbf{R})\rangle$ is the two-point instantaneous correlation function of the particle number density fluctuations, $\beta = \overline{\kappa}\ell_D$, $\ell_D = a\ell_\eta / Sc^{1/2}$, $Sc = \nu / D$ is the Schmidt number, ν is the kinematic viscosity, D is the coefficient of the Brownian diffusion of particles, $\ell_\eta = \ell_0 / Re^{3/4}$ is the Kolmogorov turbulent length. In derivation of Eq. (6) we applied the quasi-linear approach, i.e., we neglected in Eq. (6) the third and higher moments in fluctuations of κ' due to a small parameter, $\kappa'\ell_\eta / Sc^{1/2} << 1$ and we neglected by the correlation between the particle number density fluctuations and the black body temperature fluctuations ($\langle \kappa' I_b' \rangle = 0$). Substituting Eq. (6) into Eq. (2) we obtain the mean-field equation for the mean radiation intensity:

$$(\mathbf{s}\cdot\nabla)\,\overline{I}(\mathbf{s},\mathbf{r}) = \kappa_{eff}\left(\overline{I}_b - \overline{I}\right), \tag{9}$$

where the effective absorption coefficient, κ_{eff}, is

$$\kappa_{eff} = \overline{\kappa}\left(1 - \frac{2\beta J_1}{1+2\beta J_2}\right). \tag{10}$$

Integrals J_1 and J_2 in Eq. (10) are calculated using the two-point correlation function of the particle number density fluctuations for temperature stratified turbulence derived in [10]. Finally, we obtain for the effective penetration length of radiation $L_{eff} = 1/\kappa_{eff}$

$$L_{eff} / L_a = 1 + \frac{2a}{Sc^{1/2}}\left(\frac{n_{cl}}{\overline{N}}\right)^2\left(\frac{\ell_\eta}{L_a}\right)\left(1 + \frac{\mu - 1}{(\mu-1)^2 + \alpha^2}\right). \tag{11}$$

Here we used the following notations: $\mu = \frac{1}{2(1+3\sigma_T)}\left[3 - \sigma_T + \frac{20\sigma_v(1+\sigma_T)}{1+\sigma_v}\right]$; $\alpha = \frac{3\pi(1+\sigma_T)}{(1+3\sigma_T)\ln Sc}$, $\sigma_v = aSt^2K^2 / \left(3 + 3bSt^2K^2\right)$ is the degree of compressibility of the particle velocity field, $K^2 = 1 + Re\cdot\left(L_*\nabla\overline{T}/\overline{T}\right)^2$, $\sigma_T = (\sigma_{T0}^2 + \sigma_v^2)^{1/2}$ is the degree of compressibility of the turbulent diffusion tensor, $\sigma_{T0} = \sigma_T(St = 0)$; $St = \tau_p / \tau_\eta$, is the Stockes number, $\tau_p = m_p / (6\pi\rho\nu r_p)$, and $\tau_\eta = \tau_0 / Re^{1/2}$ is the Kolmogorov turbulent turnover time; σ_v is the degree of compressibility of the particle velocity field, $Re = \ell_0 u_0 / \nu$ is the Reynolds number; u_0, ℓ_0, ν are the turbulent velocity, the turbulence integral scale and the kinematic viscosity; $L_* = c_s^2\tau_\eta^{3/2} / 9\nu^{1/2}$, c_s is the speed of sound.

Fig.1 (left) shows temporal evolution of the particle concentration inside the cluster for 5 μm particles. The ratio of the effective radiation penetration length, L_{eff}, to the radiation penetration length of uniformly distributed particles, L_a, versus the Stockes number is shown in the middle frame ($\ell_0 = 1\text{m}$, $u_0 = 1\text{m/s}$, $\nu = 0.2\text{cm/s}^2$, $\text{Re} = 5\cdot 10^4$, $c_s = 450\text{m/s}$).

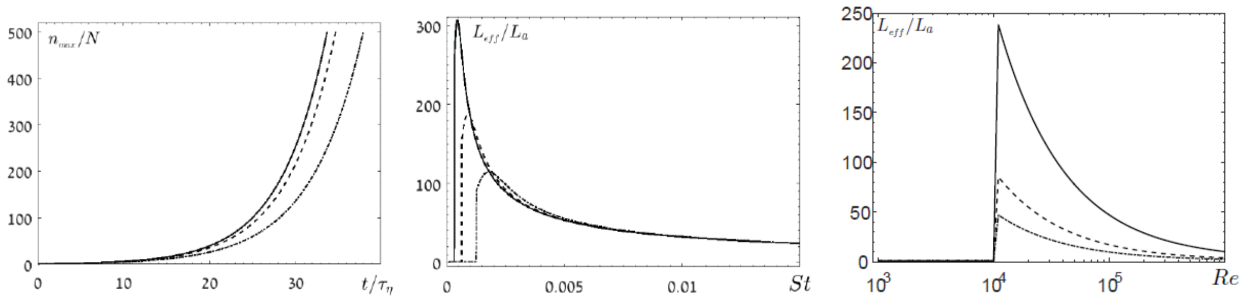

Figure 1. Left: Temporal increase of $(n_{cl}/\bar{N})$ during the particle clustering. Middle: L_{eff}/L_a versus Stockes number, Right: L_{eff}/L_a versus Reynolds number. (Solid -0.9 K/m; dashed -0.45K/m; dashed-dotted -0.225K/m).

2. Discussion

At the early stage of dust explosions the combustion mode is an accidentally ignited deflagration. Pressure waves produced by the advancing flame run away, giving rise to turbulence in the flow ahead. As velocity of the accelerating flame and the flame surface increase turbulence parameters in the flow ahead change continuously. From the moment when parameters of turbulence fall within the interval corresponding to the “transparent window” where, $L_{eff}/L_a >> 1$, the particle clusters ahead of the flame front are sufficiently exposed and heated by the radiation to become ignition kernels. The primary flame is a deflagration, which propagates with a slow velocity, therefore the time-scales during which clusters are exposed to radiation are sufficiently long for the increase of temperature of microns size particles up to the ignition level ($\Delta T \geq$ 1000K). Given a radiative flux from the primary flame, $S \approx 400\text{kW/m}^2$, the temperature of 5μm particles increases up to the ignition level in less than 10ms. If the effective penetration length of radiation increases up to 10-20m (Fig.1, middle) and for ignition time-scales of the fuel-air mixture by the radiatively heated particle cluster of the order 20ms, we obtain the rate of combustion due to the secondary explosion ~ 1000 m/s. Such intensity of the secondary explosion corresponds to the intensity of shock waves with Mach numbers M=2-3, producing overpressures of 8-10 atm, which explains the level of damages observed in unconfined dust explosions. It is obvious that parameters of turbulence are changed significantly after the secondary explosion, so that the condition $L_{eff}/L_a >> 1$ is no longer satisfied. This means that during the next phase the combustion rate will be low, until the turbulent parameters ahead of the growing and accelerating flame will fall again within the interval corresponding to the “transparent window” appropriate for the next secondary explosion. This agrees fairly well with the Buncefield explosion scenario [5]: “*The high overpressures in the cloud and low average rate of flame advance can be reconciled if the rate of flame advance was episodic, with periods of very rapid combustion being punctuated by pauses when the flame advanced very slowly. The widespread high overpressures were caused by the rapid phases of combustion; the low average speed of advance was caused by the pauses*.”

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