Massive star formation in 100,000 years from turbulent and pressurized molecular clouds

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Massive stars (with mass $m_* > 8$ solar masses $M_\odot$) are fundamental to the evolution of galaxies, because they produce heavy elements, inject energy into the interstellar medium, and possibly regulate the star formation rate. The individual star formation time, t_{*f}, determines the accretion rate of the star; the value of the former quantity is currently uncertain by many orders of magnitude[1–6], leading to other astrophysical questions. For example, the variation of t_{*f} with stellar mass dictates whether massive stars can form simultaneously with low-mass stars in clusters. Here we show that t_{*f} is determined by the conditions in the star's natal cloud, and is typically $\sim 10^5$ yr. The corresponding mass accretion rate depends on the pressure within the cloud—which we relate to the gas surface density—and on both the instantaneous and final stellar masses. Characteristic accretion rates are sufficient to overcome radiation pressure from $\sim 100 M_\odot$ protostars, while simultaneously driving intense bipolar gas outflows. The weak dependence of t_{*f} on the final mass of the star allows high- and low-mass star formation to occur nearly simultaneously in clusters.

Massive stars form in dense molecular clumps inside molecular clouds[7]. These regions are highly turbulent and are in approximate virial equilibrium, with comparable values of the gravitational energy and the kinetic energy. Observed star-forming regions have a virial parameter $\alpha_{\rm vir} = 5\langle\sigma^2\rangle R_{\rm cl}/GM_{\rm cl}$ of order unity[8], where $\langle\sigma^2\rangle$ is the mass-averaged one-dimensional velocity dispersion, $R_{\rm cl}$ the radius, and $M_{\rm cl}$ the mass of the clump. The regions of high-mass star formation studied in ref. 7 are characterized by masses $M_{\rm cl} \approx 3{,}800 M_\odot$ and radii $R_{\rm cl} \approx 0.5$ pc. The corresponding mean column density and visual extinction are $\Sigma \approx 1\,{\rm g\,cm^{-2}}$ and $A_V = (N_{\rm H}/2 \times 10^{21}\,{\rm cm^{-2}})$ mag $= 214\,\Sigma$ mag, with a dispersion of a factor of a few. These column densities are far greater than those associated with regions of low-mass star formation, but are comparable to the central surface density in the Orion nebula cluster[9], which is about $1\,{\rm g\,cm^{-2}}$. These regions have very high mean pressures $\bar{P}$:

$$\frac{\bar{P}}{k_{\rm B}} = \frac{3M_{\rm cl}\langle\sigma^2\rangle}{4\pi R_{\rm cl}^3 k_{\rm B}} = \left(\frac{3\pi f_{\rm gas}\alpha_{\rm vir}}{20k_{\rm B}}\right) G\Sigma^2 \approx 2.28 \times 10^8 f_{\rm gas}\alpha_{\rm vir}\Sigma^2\ {\rm K\,cm^{-3}} \qquad (1)$$

where $k_{\rm B}$ is Boltzmann's constant and $f_{\rm gas}$ is the fraction of the cloud's mass that is in gas, as opposed to stars. They have power-law density profiles[10] $\rho \propto r^{-k_\rho}$, with $k_\rho \approx 1.5 \pm 0.5$.

Molecular clouds are observed to be inhomogeneous on a wide range of scales, and numerical simulations show that this is a natural outcome of supersonic turbulence[11]. Dense, self-gravitating inhomogeneities that are destined to become stars are termed cores. We assume that the distribution of core masses determines the resulting distribution of stellar masses (the initial mass function), and we take this distribution as given. Because of the disruptive effects of protostellar outflows, only a fraction $\epsilon_{\rm core}$ of the total core mass $M_{\rm core}$ ends up in the star: $m_{*f} = \epsilon_{\rm core} M_{\rm core}$, where m_{*f} is the final stellar mass (for binary and other multiple star systems, m_{*f} is the total mass of stars in the system). To estimate $\epsilon_{\rm core}$ for massive stars, we assume that their outflows are scaled versions of the outflows from low-mass stars, which is qualitatively consistent with observation[12]. For low-mass stars, protostellar outflows typically eject 25–75% of the core mass[13]. A similar analysis applied to high-mass stars yields comparable results (J.C.T. and C.F.M., manuscript in preparation), so we adopt $\epsilon_{\rm core} = 0.5$ as a typical value.

Our fundamental assumption is that star-forming clumps and the cores embedded within them are each part of a self-similar, self-gravitating turbulent structure that is virialized—that is, in approximate hydrostatic equilibrium—on all scales above the thermal Jeans mass. The density and pressure are then power laws in radius (for the pressure, $P \propto r^{-k_P}$), so that the cloud is a polytrope with $P \propto \rho^{\gamma_p}$. In hydrostatic equilibrium[14], $k_\rho = 2/(2-\gamma_p)$ and $k_P = \gamma_p k_\rho = 2\gamma_p/(2-\gamma_p)$. Let $c \equiv (P/\rho)^{1/2} \propto r^{(1-\gamma_p)/(2-\gamma_p)}$ be the effective sound speed; if the pressure is dominated by turbulent motions, then c is proportional to the line width. Molecular clouds and cloud cores satisfy a line width–size relation in which c increases outwards[15], corresponding to $\gamma_p < 1$. The equation of hydrostatic equilibrium gives $M = k_P c^2 r/G$ for the mass inside r, and $\rho = Ac^2/(2\pi G r^2)$ for the density at r, where $A = \gamma_p(4-3\gamma_p)/(2-\gamma_p)^2$. It is then possible to determine the properties of a core in terms of the pressure at its surface and the mass of the star that will form within it (see Fig. 1). The radius of a core is then $r_{\rm s} = 0.074(m_{*f}/30M_\odot)^{1/2}\Sigma^{-1/2}$ pc; recall that the typical clump observed in ref. 7 has a radius of 0.5 pc, which is set by the angular resolution of those observations. The mean density in a core is $\bar{n}_{\rm H} = 1.0 \times 10^6 (m_{*f}/30M_\odot)^{-1/2}\Sigma^{3/2}\,{\rm cm^{-3}}$, and the r.m.s. velocity dispersion is $1.65(m_{*f}/30M_\odot)^{1/4}\Sigma^{1/4}\,{\rm km\,s^{-1}}$. The thermal sound speed at temperature T is $0.3(T/30{\rm K})^{1/2}\,{\rm km\,s^{-1}}$, and so a high-mass core is dominated by supersonic turbulence and should be very clumpy.

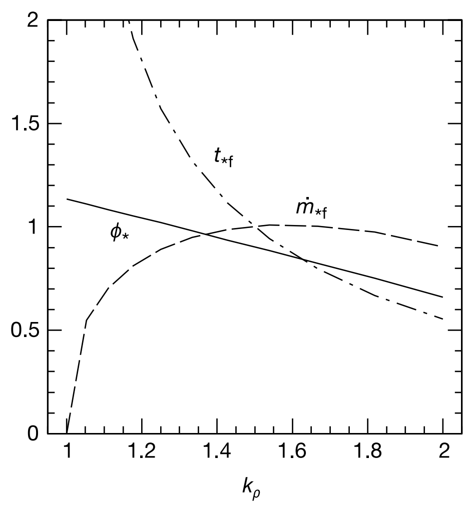

Figure 1 Variation of model parameters and results with k_ρ. Let $P_{\rm s} = \phi_P \bar{P}$ be the surface pressure of a core. The properties at the surface of the core are given by $r_{\rm s} = [AGm_{*f}^2/(2\pi k_P^2 \epsilon_{\rm core}^2 P_{\rm s})]^{1/4}$, $\rho_{\rm s} = [Ak_P^2\epsilon_{\rm core}^2 P_{\rm s}^3/(2\pi G^3 m_{*f}^2)]^{1/4}$, and $c_{\rm s} = [2\pi G^3 m_{*f}^2 P_{\rm s}/(Ak_P^2\epsilon_{\rm core}^2)]^{1/8}$. We anticipate that the overall star formation efficiency in the clump will be relatively high, so in equation (1) we adopt $f_{\rm gas} = 2/3$ as a fiducial value. We estimate $\phi_P = P_{\rm s}/\bar{P} \simeq 2$ (C.F.M. and J.C.T., manuscript in preparation), and we set $\alpha_{\rm vir}$ equal to unity[7]. We take $k_\rho = 1.5$ as a typical value, corresponding to $\gamma_p = 2/3$, $k_P = 1$ and $A = 3/4$. Following equation (3), we evaluate ϕ_* using the results of ref. 2: $\phi_* = \pi\sqrt{3}[(2-\gamma_p)^2(4-3\gamma_p)^{(7-6\gamma_p)/2}8^{(3\gamma_p-5)/2}m_0]^{1/(4-3\gamma_p)}$, where m_0 is tabulated in ref. 2. Over the entire range of γ_p and k_ρ relevant to molecular clouds $(0 \le \gamma_p \le 1, 1 \le k_\rho \le 2)$, $\phi_* \approx 1.62 - 0.48k_\rho$ to within about 1%. The star formation time decreases from $3.5t_{\rm ff}$ to $1.5t_{\rm ff}$ as γ_p varies from 0 to 1. The variation of t_{*f} and $\dot{m}_{*f}$ relative to the $k_\rho = 1.5$ case is also shown. Note that the singular polytropic model in hydrostatic equilibrium breaks down for $k_\rho = 1$, $\gamma_p = 0$, because then the pressure gradient vanishes $(k_P = 0)$.

We now consider the timescale for a star to form in such a core. On dimensional grounds, we expect the protostellar accretion rate to be[16]:

$$\dot{m}_* = \phi_* \frac{m_*}{t_{\rm ff}} \tag{2}$$

so long as radiation pressure is not important. Here m_* is the instantaneous stellar mass, $t_{\rm ff} = [3\pi/(32G\rho)]^{1/2}$ is the free-fall time and ϕ_* is a dimensionless constant of order unity. This equation could be violated in the sense that $\phi_* \gg 1$ only in the unlikely case that the star forms from a coherent spherical implosion; if the star formation is triggered by an external increase in pressure, ϕ_* could be increased somewhat, but deviations from spherical symmetry in the triggering impulse and in the protostellar core will generally prevent ϕ_* from becoming too large. It could be violated in the opposite sense that $\phi_* \ll 1$ if the core is magnetically dominated, so that collapse could not begin until the magnetic field diffused out of the core. However, magnetic fields are not observed to be dominant in molecular cores[17]. Equation (2) implies that the accretion rate, and thus the star formation time $t_{*{\rm f}} \propto m_*/\dot{m}_*$, depends weakly on the properties of the ambient medium, $\dot{m}_* \propto \rho^{1/2}$.

We now show that if the collapse is spherical and self-similar, then ϕ_* is quite close to unity provided that radiation pressure does not disrupt the flow. Although we assume the collapse is spherical far from the star, it will naturally proceed via a disk close to the star owing to the angular momentum of the accreting material; we assume this does not limit the flow of matter onto the star as otherwise the disk would become very massive and gravitationally unstable. Shielding by the disk reduces the importance of radiation pressure on the accretion flow, and allows the formation of massive stars provided that the accretion rate is sufficiently large[18]. Under the assumption of a polytropic structure in hydrostatic equilibrium, equation (2) implies:

$$\dot{m}_* = \phi_* \epsilon_{\rm core} \frac{4}{\pi\sqrt{3}} k_P A^{1/2} \frac{c^3}{G} \tag{3}$$

where the value of ρ entering $t_{\rm ff}$ is evaluated at the radius in the initial cloud that just encloses the gas that goes into the star when its mass is m_*. For the isothermal case ($\gamma_p = 1$) and for $\epsilon_{\rm core} = 1$, this reduces to the classic result of ref. 19 with $\phi_* = 0.975\pi\sqrt{3/8} = 0.663$. The accretion rate $\dot{m}_*$ is[2] proportional to $t^{3-3\gamma_p}$, so that for $\gamma_p < 1$, as observed, the accretion rate accelerates[2,4]. As discussed in ref. 2, termination of the accretion breaks the self-similarity once the expansion wave encloses $m_{*{\rm f}}$, which occurs at $t \approx 0.45 t_{*{\rm f}}$. Thereafter, the relation $\dot{m}_* \propto t^{3-3\gamma_p}$ becomes approximate, but the approximation should be reasonably good[2]. The star formation time is then given by:

$$t_{*{\rm f}} = \frac{(4-3\gamma_p)}{\phi_*} t_{\rm ff} \tag{4}$$

Evaluating ϕ_* (see Fig. 1), we conclude that $\dot{m}_* \approx m_*/t_{\rm ff}$ to within a factor 1.5 for spherical cores in which the effective sound speed increases outward.

We can express the accretion rate in terms of the mean pressure of the clump in which the star forms, $\bar{P}$ (equation (1)), and the final mass of the star, $m_{*{\rm f}}$:

$$\dot{m}_* = 4.75 \times 10^{-4} \epsilon_{\rm core}^{1/4} (f_{\rm gas}\phi_P \alpha_{\rm vir})^{3/8} \left(\frac{m_{*{\rm f}}}{30 M_\odot}\right)^{3/4} \Sigma^{3/4} \left(\frac{m_*}{m_{*{\rm f}}}\right)^{j} M_\odot\,{\rm yr}^{-1} \tag{5}$$

where $\phi_P \equiv P_s/\bar{P}$ is the ratio of the core's surface pressure to the mean pressure in the clump, and $j \equiv 3(2-k_\rho)/[2(3-k_\rho)]$; for $k_\rho = 3/2$, $j = 1/2$. This typical accretion rate is large because the core is embedded in a high-pressure environment, so that in hydrostatic equilibrium it has a high density and short free-fall time. Because massive cores are turbulent and clumpy, we expect the accretion rate to exhibit large fluctuations. Whereas clumpiness is an attribute of massive star-forming regions in our model, it is not a prerequisite, as suggested in ref. 3. The star formation time is:

$$t_{*{\rm f}} = 1.26 \times 10^5 \epsilon_{\rm core}^{-1/4} (f_{\rm gas}\phi_P \alpha_{\rm vir})^{-3/8} \left(\frac{m_{*{\rm f}}}{30 M_\odot}\right)^{1/4} \Sigma^{-3/4}\,{\rm yr} \tag{6}$$

The weak dependence of the star formation time on the final stellar mass means that low-mass and high-mass stars can form approximately at the same time. Furthermore, $t_{*{\rm f}}$ is small compared to the estimated timescale ($\sim 10^6$ yr) for cluster formation[20]. With the fiducial values of the parameters, a $100 M_\odot$ star forming in a clump with $\Sigma = 1\,{\rm g\,cm^{-2}}$ has a final accretion rate of $1.1 \times 10^{-3}\,M_\odot\,{\rm yr}^{-1}$ and a star formation time of 1.8×10^5 yr.

A necessary condition for massive star formation is that the ram pressure associated with accretion exceed the radiation pressure at the point where the dust in the infalling gas is destroyed[21]. For a $100 M_\odot$ star, this requires[18] $\dot{m}_* \gtrsim 6 \times 10^{-4}\,M_\odot\,{\rm yr}^{-1}$, which is satisfied for $\Sigma \gtrsim 0.5\,{\rm g\,cm^{-2}}$. Once the core has formed a star, the star can continue to grow by Bondi-Hoyle accretion[22] provided $m_* \lesssim 10 M_\odot$ so that radiation pressure does not prevent the focusing of gas streamlines in the wake of the star. The Bondi–Hoyle accretion rate $\dot{m}_* = 2.2 \times 10^{-7} (M_{\rm cl}/4 \times 10^3 M_\odot)^{-5/4} \Sigma^{3/4} (m_*/10 M_\odot)^2\,M_\odot\,{\rm yr}^{-1}$ (C.F.M. and J.C.T., manuscript in preparation) is so low that this process does not significantly alter the stellar mass under the conditions observed in the clumps of ref. 7.

Direct comparison of our results with observation is difficult

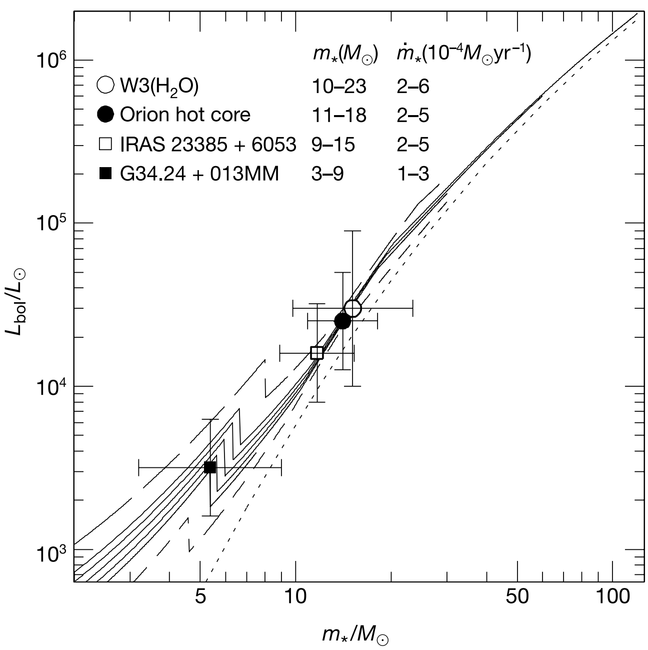

Figure 2 Derived properties of nearby massive protostars. Solid lines show the predicted evolution in luminosity of protostars (including their accretion disks, but allowing for the powering of protostellar outflows so that $f_{\rm acc} = 0.5$; J.C.T. and C.F.M., manuscript in preparation) of final mass 7.5, 15, 30, 60 and $120 M_\odot$ accreting from cores with $k_\rho = 1.5$ embedded in a $\Sigma = 1\,{\rm g\,cm^{-2}}$ clump, typical of Galactic regions[7]. The luminosity step occurring at around 5–$7 M_\odot$, depending on the model, corresponds to the onset of deuterium shell burning, which swells the protostellar radius by a factor of about two and thus reduces the accretion luminosity by the same factor. The dashed and long-dashed lines show a $30 M_\odot$ star forming in a clump with mean pressure ten times smaller and larger than the fiducial value, respectively. The dotted line shows the zero-age main-sequence luminosity from ref. 26. Four observed hot molecular cores are shown: G34.24+0.13MM[27], IRAS23385+6053[28], Orion hot core[29] and W3(H$_2$O)—the Turner–Welch object[30]. The vertical error bar illustrates the uncertainty in the bolometric luminosity. The horizontal error bar shows the corresponding range of allowed values of m_* for the $\Sigma = 1\,{\rm g\,cm^{-2}}$ models. These values and the constraints on $\dot{m}_*$ are listed here for each source.

because the actual masses of massive protostars are poorly determined. Our approach is to predict the properties of some well studied massive protostars in terms of their bolometric luminosities. The bolometric luminosity L_{bol} has contributions from main-sequence nuclear burning L_{ms}, deuterium burning L_D, and accretion L_{acc}. The accretion luminosity $L_{acc} = f_{acc} G m_* \dot{m}_* / r_*$, where f_{acc} is a factor of order unity accounting for energy radiated by an accretion disk, advected into the star or converted into kinetic energy of outflows, and where the stellar radius r_* may depend sensitively on the accretion rate $\dot{m}_*$. Massive stars join the main sequence during their accretion phase at a mass that also depends on the accretion rate[23]. To treat accelerating accretion rates, we have developed a simple model for protostellar evolution based on that of refs 6 and 24. The model accounts for the total energy of the protostar as it accretes and dissociates matter and, if the central temperature $T_c \gtrsim 10^6$ K, burns deuterium. We have modified this model to include additional processes, such as deuterium shell burning, and we have calibrated these modifications against the more detailed calculations of refs 23 and 25.

Our model allows us to make predictions for the masses and accretion rates of embedded protostars that are thought to power hot molecular cores (C.F.M. and J.C.T., manuscript in preparation). Figure 2 compares our theoretical tracks with the observed bolometric luminosities of several sources. We find that uncertainties in the value of the pressure create only small uncertainties in m_* for L_{bol} in excess of a few times 10^4 solar luminosities.

The infrared and submillimetre spectra of accreting protostars and their surrounding envelopes have been modelled in ref. 5, modelling the same sources shown in Fig. 2. We note that uncertainties in the structure of the gas envelope and the possible contributions from additional surrounding gas cores or diffuse gas will affect the observed spectrum. Comparing results, our inferred stellar masses are similar, but our accretion rates are systematically smaller by factors of ~2–5. The modelled[5] high accretion rates of $\sim 10^{-3} M_{\odot} \mathrm{yr}^{-1}$ for stars with $m_* \approx 10 M_{\odot}$ would be difficult to achieve unless the pressure was increased substantially; for example, if the stars are destined to reach $m_{*f} \approx 30 M_{\odot}$, pressure increases of a factor ~40 are required. □

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Coherent emission of light by thermal sources

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A thermal light-emitting source, such as a black body or the incandescent filament of a light bulb, is often presented as a typical example of an incoherent source and is in marked contrast to a laser. Whereas a laser is highly monochromatic and very directional, a thermal source has a broad spectrum and is usually quasi-isotropic. However, as is the case with many systems, different behaviour can be expected on a microscopic scale. It has been shown recently[1,2] that the field emitted by a thermal source made of a polar material is enhanced by more than four orders of magnitude and is partially coherent at a distance of the order of 10 to 100 nm. Here we demonstrate that by introducing a periodic microstructure into such a polar material (SiC) a thermal infrared source can be fabricated that is coherent over large distances (many wavelengths) and radiates in well defined directions. Narrow angular emission lobes similar to antenna lobes are observed and the emission spectra of the source depends on the observation angle—the so-called Wolf effect[3,4]. The origin of the coherent emission lies in the diffraction of surface-phonon polaritons by the grating.

It is usually taken for granted that light spontaneously emitted by