THE ASTROPHYSICAL JOURNAL, 456:L97–L100, 1996 January 10

THE ROLE OF THERMAL INSTABILITIES ON THE THEORETICAL MASS FUNCTION OF MACHOs

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Received 1995 June 21; accepted 1995 October 19

ABSTRACT

We analyze the primordial protogalactic gas to investigate the role of thermal instabilities in the formation of condensations that could contribute to the MACHO mass function. We find that the previrialized protogalaxy could have undergone a thermally unstable regime, which defines a peak in mass condensations starting around 0.1 $M_\odot$ after being reprocessed by the environment that may be identified as giving origin to the MACHOs. Effects of magnetic fields on this instability-driven MACHOs scenario are also addressed.

Subject headings: dark matter — Galaxy: halo — stars: low-mass, brown dwarfs

1. INTRODUCTION

Two years after the first reports of candidate microlensing events (Alcock et al. 1993; Aubourg et al. 1993; Udalski et al. 1993), the field is rapidly growing and becoming a major tool for investigating the morphology and history of our Galaxy (see Zylberajch 1994 and Paczyński 1995 for reviews). Particularly intriguing is the fact that a massive halo composed of MACHOs seems *not* to be favored by the observations (Alcock et al. 1995; Gates et al. 1995). The nature of dark objects so far responsible for the events toward the LMC needs also to be addressed. It is not clear as yet if they belong to a disk of our own Galaxy, to the Galactic halo itself, or even to the LMC (Sahu 1994). We shall adopt in this work the hypothesis that MACHOs are indeed located in the halo, and we shall attempt to outline a scenario that could have led to their formation.

A number of works have stressed the possible relevance of *thermal instabilities* (TI) for globular cluster formation during the collapse of the protogalactic cloud (e.g., Fall & Rees 1985; Vietri & Pesce 1995 and references therein). While these scenarios address an epoch in which the gas has already been shock heated to its virial temperature, $T \sim 10^6$ K, we will be mainly interested in a previous, optically thin stage of the cloud at $T \sim$ few times 10^4 K. We also note some previous works dealing with the protogalaxy that have addressed the question of Population III star formation (e.g., Silk 1983). The motivation for a reassessment of such scenarios is stronger than ever, particularly now that an initial spectrum can be calculated and its reprocessing characterized observationally by the growing body of data. To stress this point, let us consider, as in those previous works, a quasi-static, isothermal protogalactic cloud with an estimated virial temperature $T \sim \mu v_r^2/2k_{\rm B}$, or $\sim 10^6$ K. It is known that at earlier stages the faster cooling at the center would have destroyed the dynamical equilibrium and induced the collapse of the cloud, mainly because the equilibrium requires the opacity of the gas to be large enough in order to delay radiative losses. On the other hand, an optically thin gas will cause the cloud to cool more quickly, so that at least for timescales of the order of the free-fall time $t_{\rm ff} \sim (G\rho)^{-1/2}$, the protogalactic gas will remain at a temperature much smaller than the virial one. There is also an observational motivation to consider a previous epoch of the galaxy; namely, that a substantial fraction of the halo mass is in the form of field stars, not in globular clusters at a ratio of 1/100. Furthermore, the most metal-poor clusters have [Fe/H] $\simeq -2.5$, while some field stars have been detected with [Fe/H] $\simeq -4$, which suggests that they were formed *before* globular clusters.

2. THERMAL INSTABILITY

Motivated by the above discussion, we consider a previrialized collapsing cloud with mass $M \simeq 10^{12}\ M_\odot$, radius $r \gtrsim 50$ kpc (corresponding to an average density $\rho \simeq 10^{-25}$ g cm^{-3}), and $T =$ few times 10^4 K. Such a cloud would have started to collapse at $z \leq 10$ when its mean density was ~ 5 times the background density (Rees & Ostriker 1977).

At temperatures below $T \sim 3 \times 10^4$ K, the primordial gas of the cloud (which has a metal abundance $Z \sim 0$) cooling through radiative thermal processes (collisional excitation and recombination) is stable to condensation formation by isobaric thermal instability (e.g., Field 1965; Corbelli & Ferrara 1995). We assume that the onset for the thermal instability occurs at $T \gtrsim 3 \times 10^4$ K. We also assume a perfect gas equation of state $P = R\rho T/\mu$, where $\mu = \langle m \rangle / m_{\rm H}$ and ρ, T are the equilibrium density and temperature of the cloud. A perturbed gas element is subject to an isobaric thermal instability growing to the first order at a rate (Field 1965; see also Begelman 1990; de Gouveia Dal Pino & Opher 1989a, b, 1990, 1991, 1993)

$$n = \left(\frac{\partial \mathscr{L}}{\partial T}\bigg|_\rho - \frac{\rho}{T}\frac{\partial \mathscr{L}}{\partial \rho}\bigg|_T + \frac{k^2 K}{\rho} \right) \frac{(1-\gamma)\mu}{\gamma R}, \qquad (1)$$

where μ is the mean molecular weight (= 0.59 for an ionized gas of primordial composition), $\gamma = 5/3$, $K = 0.77 \times 10^{-6}\ T^{5/2}$ is the thermal conductivity term, and $\mathscr{L} = L(\rho, T) - G(\rho, T)$, i.e., cooling minus heating, is the net energy loss per unit time and mass, which vanishes at equilibrium. Note that, strictly speaking, equation (1) holds for the short wavelengths of the perturbation. A very important (and uncertain) point is the nature and dependence of the heating function $G(\rho, T)$. The heating function may include the UV flux of an active galactic nucleus (AGN) phase (likely to have been present in the early Galaxy because of the observational evidence collected from the bulge); turbulence (the Reynolds number Re $= 2r_c v/\eta$ is known to be very much more than 1 for the primordial gas, v

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being the gas velocity and η the shear viscosity); and MHD waves among the most important sources. Given the uncertainties of the nature of the dominant heating mechanisms in the protogalactic gas, we assume as in previous works (e.g., Field 1965; Eilek & Caroff 1979; de Gouveia Dal Pino & Opher 1989a, b, 1990, 1991, 1993) that the heating function $G(\rho, T)$ is due to an externally generated source and its dissipation is locally independent of the density and temperature, a condition that is known to hold for several of the above-mentioned sources. Thus, the derivatives of $\mathscr{L}$ in equation (1) will depend only on the cooling function L.

For a strongly ionized collisional H + He plasma of primordial composition ($X \simeq 0.76$, $Y \simeq 0.24$, and $Z \simeq 0$) at 2.5×10^4 K $\lesssim T \lesssim 10^6$ K, the net cooling is

$$L \simeq \Lambda(T) n_{\rm H}^2/\rho \text{ ergs g}^{-1}\text{ s}^{-1}, \quad (2)$$

where $n_{\rm H}$ is the hydrogen number density, and the cooling function $\Lambda(T)$, evaluated by, e.g., Fall & Rees (1985), can be fitted by

$$\log \Lambda(T) = -22.9 + 2.45\phi(T) \exp[-\phi(T)]^2 + \psi(T) \exp[-\psi(T)]^2 + \theta(T) \exp[-\theta(T)]^2, \quad (3)$$

with $\phi(T) = 1.1 \log T - 4$, $\psi(T) = 3.7 \log T - 17.6$, and $\theta(T) = 1.4 \log T - 9.2$. This rather complicated fit can be replaced by the much simpler expression $\log \Lambda(T) = -24.45 + 1.04 \log T - 0.11(\log T)^2$ in the restricted region of interest of this work, 2.5×10^4 K $\lesssim T \lesssim 10^5$ K, with sufficient accuracy. Furthermore, since $|T\partial\mathscr{L}/\partial T| \ll |\rho\partial\mathscr{L}/\partial\rho|$ (see eq. [1]), we can safely neglect $\partial\mathscr{L}/\partial T$ in our analysis and adopt a constant $\Lambda(T) \simeq 9.6 \times 10^{-23}$ ergs cm^3 s^{-1} in equation (2); this extremely simple form introduces errors of, at most, 10% in the calculated quantities.

An isobaric thermal instability grows only if the pressure in the perturbed region is equalized in a timescale that is small compared to the cooling time, defined as $t_c \equiv (3\rho^2 k_{\rm B} T)(2\mu m_{\rm H} L)^{-1}$. This implies an upper limit for the wavelength of the unstable region (e.g., Field 1965) of

$$\lambda \leq \lambda_{\rm max} \simeq 2\pi v_s t_c = 1.1 \times 10^{20} (\rho_{-25})^{-1} (T_{4.7})^{3/2} \text{ cm}, \quad (4)$$

where $v_s = (\gamma k_{\rm B} T/\mu m_{\rm H})^{1/2}$ is the speed of sound, $\rho_{-25} = \rho/(10^{-25}$ g cm$^{-3})$, and $T_{4.7} = T/(5 \times 10^4$ K).

On the other hand, it is easy to show that, because of the ability of thermal conduction to suppress the growth of a thermally unstable region, a critical wavenumber $k_{\rm max} = 2\pi/\lambda_{\rm min}$ exists over which $n = 0$, and it is given by $k_{\rm max} = [k_K (k_\rho - k_T)]^{1/2}$, where

$$k_\rho = [(\gamma - 1)\mu\rho(\partial\mathscr{L}/\partial\rho)_T](R v_s T)^{-1},$$

$$k_K = (R v_s \rho)[\mu(\gamma - 1)K]^{-1},$$

$$k_T = [(\gamma - 1)\mu(\partial\mathscr{L}/\partial T)]_\rho (R v_s)^{-1}. \quad (5)$$

After an elementary manipulation with these expressions we find $\lambda_{\rm min} = 2\pi/k_{\rm max}$, which is given by

$$\lambda_{\rm min} \simeq 2.1 \times 10^{18} (\rho_{-25})^{-1} (T_{4.7})^{7/4} \text{ cm}. \quad (6)$$

Since n is small for small k (large λ) and vanishes for $k > k_{\rm max}$, a k exists for which n reaches a maximum. This value is $k(n_{\rm max}) \simeq [k_\rho^2 (k_\rho - k_T) k_K]^{1/4}$, or in terms of $\lambda(n_{\rm max})$,

$$\lambda(n_{\rm max}) \simeq 1.5 \times 10^{19} (\rho_{-25})^{-1} (T_{4.7})^{13/8} \text{ cm}. \quad (7)$$

The most rapidly growing condensations are those at $\lambda(n_{\rm max})$. They define a typical mass $M_{\rm blob} \sim \rho_b \lambda^3(n_{\rm max})$ or

$$M_{\rm blob} \sim 1.8 \times 10^{-1}(1 + f_b e)(\rho_{-25})^{-2}(T_{4.7})^{39/8} M_\odot, \quad (8)$$

where $f_b < 1$ is the initial density perturbation relative to the ambient density ρ. This is the most probable mass of the condensations, which is bracketed between the limiting values

$$M_{\rm blob}(\lambda_{\rm min}) \sim 4.4 \times 10^{-4}(1 + f_b e)(\rho_{-25})^{-2}(T_{4.7})^{21/4} M_\odot \quad (9)$$

and

$$M_{\rm blob}(\lambda_{\rm max}) \sim 73(1 + f_b e)(\rho_{-25})^{-2}(T_{4.7})^{9/2} M_\odot, \quad (10)$$

which define the condensation mass spectrum at early times.

3. PRIMORDIAL MAGNETIC FIELD EFFECTS

If a small primordial field is present in the gas $\boldsymbol{B} \neq 0$ (which seems unavoidable in the pregalactic medium; see, e.g., Rees 1995) in the direction $\parallel$ to $\boldsymbol{B}$, the condensation is not affected by the presence of the magnetic field, and its growth is given by equation (1). However, in the direction $\perp$ to $\boldsymbol{B}$, the thermal conduction K is greatly reduced (because of the e^- spiraling between collisions), and the initial collapse of the condensation is more efficient (see, e.g., Field 1965; de Gouveia Dal Pino & Opher 1989a, b, 1990, 1991, 1993). We can therefore determine two scales in this direction $\lambda_{\perp\,\rm min}$ and $\lambda_{\perp\,\rm max}$ for the instability growth with meanings analogous to those of the former section. Defining an effective pressure given by the sum of the gas pressure and the magnetic pressure, we have (e.g., Field 1965)

$$\lambda_{\perp\,\rm min} \sim 4.1 \times 10^{17}(B_{-12})^{-1}(T_{4.7})^{1/4}[1 + (\gamma v_{\rm A}^2/v_s^2)]^{1/2} \text{ cm}, \quad (11)$$

where $v_{\rm A} = B/(4\pi\rho)^{1/2}$ is the Alfvén speed, $B_{-12} = B/(10^{-12}$ G), and (e.g., de Gouveia Dal Pino & Opher 1990)

$$\lambda_{\perp\,\rm max} \sim 1.1 \times 10^{20}(\rho_{-25})^{-1}(T_{4.7})^{3/2}[1 + (v_{\rm A}^2/v_s^2)]^{1/2} \text{ cm}. \quad (12)$$

Note that for small $\boldsymbol{B}$ the ratio $v_{\rm A}^2/v_s^2 \ll 1$ for typical values of protogalactic ρ and T. The masses that bound the spectrum $[\lambda_\parallel = \lambda(n_{\rm max})]$ can be estimated as $M_{\rm blob}(\lambda_{\perp\,\rm min}) \simeq \rho_b \pi \lambda_{\perp\,\rm min}^2 \lambda_\parallel$, which results in

$$M_{\rm blob}(\lambda_{\perp\,\rm min}) = 4.1 \times 10^{-4}(1 + f_b e)(T_{4.7})^{17/8}(B_{-12})^{-2} \times [1 + (\gamma v_{\rm A}^2/v_s^2)] M_\odot, \quad (13)$$

and $M_{\rm blob}(\lambda_{\perp\,\rm max}) \simeq \rho_b \pi \lambda_{\perp\,\rm max}^2 \lambda_\parallel$, yielding

$$M_{\rm blob}(\lambda_{\perp\,\rm max}) = 31(1 + f_b e)(\rho_{-25})^{-2}(T_{4.7})^{37/8}[1 + (v_{\rm A}^2/v_s^2)] M_\odot. \quad (14)$$

4. EVOLUTION OF THE BLOBS

From equation (1), we can evaluate the growth time τ_g of the condensations in the absence of $\boldsymbol{B}$:

$$\tau_g = [1 \times 10^{-13}(\rho_{-25})(T_{4.7})^{-1} - 1.2 \times 10^{22}(T_{4.7})^{5/2}(\rho_{-25})^{-1}k^2]^{-1} \text{ s} \quad (15)$$

for $k \simeq k(n_{\rm max})$ (eq. [7]) and T, ρ close to our scaling values; the first term dominates and gives $\tau_g \simeq 2.9 \times 10^5$ yr. This initial growth timescale is much smaller than the free-fall time of the condensed blobs $\tau_{{\rm ff},b} \sim (G\rho_b)^{-1/2}$ and indicates that initially the condensations will not collapse gravitationally but rather will contract and cool because of the thermal instability itself. (In the presence of $\boldsymbol{B}$, the time is given by

$\tau_{g,B} \simeq 9 \times 10^{12} T_{4.7}(\rho_{-25})^{-1}[1 + (v_A^2/v_s^2)]$ s, which for $v_A^2/v_s^2 \ll 1$ is also of the order of τ_g.) In fact, the condensed blob will become gravitationally unstable if its mass exceeds the Jeans mass $M_J = 1.2(k_B T_b/\mu m_H)^2 G^{-3/2} P^{-1/2}$ (where T_b is the initial temperature of the blob and P is the pressure of the unperturbed medium), which is much larger than $M_{\rm blob}$ for temperatures $\simeq 10^4$ K.

A blob in pressure equilibrium with the ambient gas will have a density $\rho_b = (T\mu_b/T_b\mu)\rho$ (where $\mu_b \sim 1.22$ for a weakly ionized gas inside the blob). Once the temperature of the blob dropped below 10^4 K, the cooling will be dominated by the collisional excitation of the rotational and vibrational transitions of the H_2 molecule (the heavy elements also contribute, but for very small abundances $Z/Z_\odot$, they do not dominate the cooling; see, e.g., Fall & Rees 1985). Thus, just after its formation the blob will cool in a time $\tau_{c,i} \sim (3\rho_b k_B T_b)(2\mu_b m_H L_{H_2})^{-1}$, with $L_{H_2} \simeq 10^{-22}\, n_b(H_2)$ ergs cm^{-3} s^{-1} for an initial hydrogen number density in the blob $n_b(H) \sim 1$ cm^{-3} (e.g., Lepp & Shull 1983). The number density of molecular hydrogen $n_b(H_2)$ can be evaluated from statistical equilibrium, taking into account the relevant reaction rates and the effect of UV dissociation resulting from the hotter ambient gas; we find $n_b(H_2) \simeq 3 \times 10^{-2}\, x n_b(H)$ at $T_b \sim 9000$ K, where $x \equiv n_b(H^+)/n_b(H) < 1$ is the fractional ionization (see, e.g., Fall & Rees, eq. [18], and references therein). In this estimate we have assumed a photoionization rate $k_{UV} \simeq 2.4 \times 10^{-13}(r/{\rm kpc})^{-1}$ s^{-1}, appropriate for our physical conditions. This cooling time $\tau_{c,i} \simeq 2.0 \times 10^4\, x^{-1}(T_b/9000$ K) yr is now larger than the condensation formation time τ_g (eq. [15]) for $x \lesssim 10^{-2}$ (see, e.g., Dalgarno & McCray 1972; Fall & Rees 1985) but is still much smaller than the protogalactic free-fall time $t_{\rm ff} \sim 4 \times 10^8(\rho/10^{-25}$ g cm$^{-3})^{-1/2}$ yr, so that the blobs will continue to cool within the protogalactic infall gas.

A cold blob moving relative to the ambient gas may be subject to the Kelvin-Helmholtz (K-H) instability at the interface between blob and ambient gas with a characteristic growth time $\tau_{\rm K\text{-}H} \simeq \lambda_b(\rho_b/\rho)^{1/2}/u$, where λ_b is the radius of the blob and u is the relative velocity (Murray et al. 1993). Buoyancy effects may accelerate some blobs to speeds comparable to the sound speed, for which $\tau_{\rm K\text{-}H} \sim \lambda_b/(\gamma k_B T_b/\mu_b m_H)^{1/2}$. Imposing the "survival" condition for a blob against K-H instability $\tau_{c,i} < \tau_{\rm K\text{-}H}$ gives a lower limit to the size of the blobs $\lambda_b > 6.3 \times 10^{17}\, x^{-1}(T_b/9000$ K$)^{3/2}$ cm. Thus, blobs moving with relative speed close to the sound speed are expected to break up, since they do not satisfy this condition for $x \lesssim 10^{-2}$ (eqs. [7], [11], and [12]). On the other hand, those "moving with" the collapsing protocloud will be stable to the K-H instability. The presence of magnetic fields will further inhibit the K-H instability effects, although the surviving fraction of blobs *must not* be expected to be large, since in that case the observed MACHO fraction would also be large, which is not the case (Alcock et al. 1995; Gates, Gyuk, & Turner 1995). A detailed calculation of the K-H effects is out of the scope of this work and will be discussed elsewhere.

Finally, the blobs must survive to a much more "dangerous" later epoch in which the protocloud has been heated to $T \sim 10^6$ K (for some radius) and $\rho \sim 10^{-24}$ g cm^{-3} (see, e.g., Fall & Rees 1985). The blobs that should have by now cooled to $T_b \lesssim 1000$ K since $\tau_{c,f} \simeq 13 n_b(H)/n_b(H_2)$ yr [where we have assumed $L_{H_2} \simeq 5 \times 10^{-22}\, n_b(H_2)$ ergs cm^{-3} s^{-1} for $n_b(H) \sim 10^3$ cm^{-3}; see Lepp & Shull 1983] will be heated by X-rays from the hot ambient protocloud. The effective cross section for photoionization of the blob gas due to $\sim$0.1 keV photons is $\sigma \simeq 1.8 \times 10^{-20}\, (h\nu/150$ eV$)^{-3}$ cm^2. Since the radiation is mainly thermal bremsstrahlung with an intensity $\propto \exp(-h\nu/k_B T)$, $\sim$90% of the heat will be deposited in a layer of width $l_X \simeq [n_b(H)\sigma]^{-1} \simeq 1.3 \times 10^{17}(T/10^6$ K$)^2(\rho/10^{-24}$ g cm$^{-3})^{-1}(T_b/1000$ K) cm by photons having $\langle h\nu\rangle \sim 2.3 k_B T$. This length scale defines a mass $M_X \simeq 2.1 \times 10^{-3}(T/10^6$ K$)^7(\rho/10^{-24}$ g cm$^{-3})^{-2}(T_b/1000$ K$)^2\, M_\odot$. Therefore, blobs with $M_{\rm blob} \leq M_X$ will be heated throughout, and the more massive ones will be heated only in a thin layer $\sim l_X$ near their surfaces. This X-ray heating will kill the smallest blobs from the original spectrum (eqs. [9] and [13]), providing a possible explanation to the lack of low-mass ($M \lesssim 10^{-3}\, M_\odot$) candidate events for MACHOs (Zylberajch 1994). These figures must be considered as lower limits to the heating rates since any other sources such as supernova explosions or cosmic rays would increase l_X. However, we must stress that the blobs are in principle able to cool quickly to $\lesssim 100$ K in a timescale of $\sim 3.3 \times 10^3[n_b(H)/n_b(H_2)]^2$ yr [for $L_{H_2} \simeq 2 \times 10^{-29} n_b(H_2)^2$ ergs cm^{-3} s^{-1} for $n_b(H) \sim 10^4$ cm^{-3}; see Capuzzo-Dolcetta, Di Fazio, & Palla 1995], which is shorter than the protogalactic $t_{\rm ff}$ at that epoch. These low temperatures are enough for the blobs to become gravitationally unstable, and therefore it is not unlikely that a number of them can survive until today.

Several other processes (not discussed here) may be also important for the evolution of the blob (including ambipolar diffusion at $T_b \leq 10^4$ K [e.g., Shu et al. 1987] and $\sim 10^8$ collisions between blobs since their formation), which may bias the final mass distribution in a yet unknown manner.

The main result of these calculations is that a thermally unstable regime of the protogalactic gas before the standard stellar formation trigger induces the formation of condensations or blobs in a suitable mass interval. If the physical conditions of the protogalactic gas are not very different from the standard models of the protogalaxy, the most abundant surviving blobs have typical masses $M \gtrsim 10^{-1}\, M_\odot$, which is consistent with the microlensing candidates reported by the MACHO and EROS collaborations. In this scenario, the MACHOs may be identified with cold clouds that started to collapse from condensed blobs because of the thermal instability of the gas. It has also implications for the generation of low-mass Population III luminous stars, a subject that has received some attention in the recent past (e.g., Silk 1983; Bessel & Norris 1984; Beers, Preston, & Shectman 1985; Gass, Liebert, & Wehrse 1988).

After the completion of this Letter, we were informed of a recent work (De Paolis et al. 1994) that discusses the issue of thermal instabilities forming $M \sim 10^6\, M_\odot$ hot clouds that later fragment into low-mass objects identified with the MACHOs. Two comments are in order. First, as stated by the authors, clustering of the lensing objects is expected, and this prediction can be checked against accumulated observed events. Second, the clustered MACHOs formed inside those massive clouds should in any case add up to the fraction resulting from the previous epoch as discussed here; thus, we may be observing different dark populations originated closely in time but by different physics.

We are indebted to the referee A. Ferrara, for his fruitful comments and suggestions. Discussions with R. Opher, S. Viegas, A. C. S. Friaça, A. Dal Pino, Jr., and G. A. Medina-Tanco about various subjects of this work are also acknowledged. J. E. H. acknowledges the warm hospitality given to him during a visit to DAPNIA-Saclay and interesting discussions about the MACHO events maintained with the EROS group. This work was partially supported by CNPq and FAPESP grants.

REFERENCES

Alcock, C., et al. 1993, Nature, 365, 621
———. 1995, Phys. Rev. Lett., 74, 2867
Aubourg, E., et al. 1993, Nature, 365, 623
Beers, T. C., Preston, G. W., & Shectman, S. A. 1985, AJ, 90, 2089
Begelman, M. C. 1990, in The Interstellar Medium of Galaxies, ed. H. A. Thompson, Jr., & J. M. Shull (Dordrecht: Kluwer), 287
Bessel, M. S., & Norris, J. 1984, ApJ, 285, 622
Capuzzo-Dolcetta, R., Di Fazio, A., & Palla, F. 1995, in Physical Processes in Fragmentation and Star Formation, ed. R. Capuzzo-Dolcetta et al. (Dordrecht: Reidel)
Corbelli, E., & Ferrara, A. 1995, ApJ, 447, 708
Dalgarno, A., & McCray, R. A. 1972, ARA&A, 10, 375
de Gouveia Dal Pino, E. M., & Opher, R. 1989a, ApJ, 342, 686
———. 1989b, MNRAS, 240, 573
———. 1990, A&A, 231, 551
———. 1991, A&A, 242, 319
———. 1993, MNRAS, 263, 687
De Paolis, F., et al. 1995, Comments Astrophys., in press
Eilek, J. A., & Caroff, L. J. 1979, ApJ, 233, 463
Fall, S. M., & Rees, M. J. 1985, ApJ, 298, 18
Field, G. 1965, ApJ, 142, 531
Gass, H., Liebert, J., & Wehrse, R. 1988, A&A, 189, 194
Gates, E. I., Gyuk, G., & Turner, M. S. 1995, Phys. Rev. Lett., 74, 3724
Lepp, S., & Shull, J. M. 1983, ApJ, 270, 578
Murray, S. D., et al. 1993, ApJ, 407, 588
Paczyński, B. 1995, preprint
Rees, M. J. 1995, in IAU Symp. 168, Big Bang and Cosmic Background Radiation, ed. M. Kafatos (Dordrecht: Kluwer), in press
Rees, M. J., & Ostriker, J. P. 1977, MNRAS, 179, 541
Sahu, K. 1994, Nature, 370, 275
Shu, F., Adams, F., & Lizano, S. 1987, ARA&A, 25, 23
Silk, J. 1983, MNRAS, 205, 705
Udalski, A., et al. 1994, Acta Astron., 44, 165
Vietri, M., & Pesce, E., 1995, preprint
Zylberajch, S. 1995, in Proc. Vulcano Workshop, ed G. Mannocchi & F. Giovanelli (Bologna: Edit. Compositori), 105