DARK MATTER AND X-RAY CLOUDS

O. B. Firsov

Kurchatov Institute, Russian Scientific Center, 123181 Moscow, Russia


In two previous papers defending the notion of baryonic matter in the universe, the author showed that despite the high temperature of x-ray clouds, T~10^7 K, the presence of condensed dark matter is compatible with the baryonic nature of the dark matter by virtue of the clouds’ low density. The present paper discusses the origin of baryonic dark matter.

1. INTRODUCTION

In two previous papers, the author showed that in contrast to the opinion held by most astrophysicists, dark matter in the universe is most likely baryonic, and is probably described by a mass distribution function consistent with that established for meteors, meteorites, and asteroids. The problems entailed by the idea of baryonic dark matter are probably no worse than those incurred by as yet unobservable particles with m~10^-12 eV (masses m~10^-7 eV or higher are forbidden). These problems, however, can probably be overcome.

In the present paper, I show that the existence of dark matter in x-ray clouds with T~10^7 K is compatible with its baryonic nature, by virtue of the low density of these clouds.

Cosmic x rays were detected in 1962, as soon as the Einstein Observatory (HEAO-2) into low Earth orbit, and full-time studies began with the 1978 launch of the author, showing that in contrast to the opinion held by most astrophysicists, dark matter consists of condensed baryonic matter than luminous material.

2. GENERAL CONSIDERATIONS

It would be natural to suppose that the outside of an x-ray cloud is an order of magnitude cooler than its inside. If the radius of a spherical cloud is R, and the photon mean free path is l, then upon emerging from the central region of the cloud, a photon will undergo an average of (R/l)^2 collisions with electrons. The energy exchanged (usually lost) in each of these collisions is T/m_c^2, where m_c is the mass of the electron and c is the speed of light. For a photon not to “cool off,” then, we require that R~(m_c^2/T)^(1/2).

Noting that l=1/m_e, where m_e=0.66·10^-34 cm^3 is the Thomson cross section, n~m_e R^3, and the well known equilibrium relation R~GMm_e/T holds (G is the gravitational constant and m_e is the mass of the proton), it can easily be shown that

M~T^1/2(m_c^2/2p_e)(G/m_e)^2 = 10^{-6} g for

T=0.5·10^{-8} erg = 7·10^4 K.

The masses of x-ray clouds should thus be quite high, just as observed.

It will be desirable in what follows to specify a definite electron density distribution n and dark matter distribution n as a function of radius r:

n=n_0 exp[−πr^2/R_s^2],

ρ=ρ_0 exp[−πr^2/R_s^2].

These are not fundamental equations; they are merely a matter of convenience.

3. MASS M

The distribution (1) yields

M=M_0 R_s^3 + ρ_0 R_s^2,

where the first term on the right-hand side is the mass of luminous matter and m is the mass of such matter impinging on a single electron; the second term is the mass of dark matter—missing mass—than luminous material.

In the opening sections of this paper, internal parameters of a spherical plasma cloud such as the electron density n, the concentration of the ith element. The dark electron density distribution n, and dark matter distribution ρ as a function of radius r:

n=n_0 exp[−πr^2/R_s^2],

ρ=ρ_0 exp[−πr^2/R_s^2].

These are not fundamental equations; they are merely a matter of convenience.
4. TEMPERATURE T

The temperature can be determined via the plasma pressure
\[ P = n' T, \]
where \( n' \) is the total number of particles—electrons plus nuclei—per unit volume of plasma. Here we take \( n' = 1.9n. \)

\[ T = \frac{P}{n'} = \frac{mG}{1.9} \int_0^R \exp(-\pi R'^2 R^2) \int_0^\pi \exp(-\pi R'^2 R^2) dR'R^2 \]

\[ \times \int_0^\theta \exp(-\pi R'^2 R^2) (m n_0 + \exp(-\pi R'^2 R^2))\rho_0 \frac{4 \pi R'^2 dR'}{R^2} \]

at \( R = 0 \), i.e., \( T(0) = T_0 \). The integral in (3) can be evaluated analytically, and

\[ T_0 = \frac{mG}{1.9} \left( 1 - \frac{\pi}{4} \right) \left( 1 - \frac{\pi}{4} \right) \left( 1 - \frac{\pi}{4} \right) \left( 1 - \frac{\pi}{4} \right) \frac{m n_0 + \rho_0}{R^2} \]

\[ \times \left( 1 + \frac{S \arctg R_0}{R_0} \right) = 0.226 \left( 1 + \frac{S \arctg R_0}{R_0} \right) \frac{mG}{R_0} \]

\[ \times \left( 1 + \frac{S \arctg R_0}{R_0} \right) = 0.226 \left( 1 + \frac{S \arctg R_0}{R_0} \right) \frac{mG}{R_0} \]

The assumption (1) is probably tenable if \( R_0/R_0 \) is not too different from 1. But even if \( R_0/R_0 \) ranges from 0.8 to 1.25, the second term only ranges from 0.6 to 1.65. For the most part, we assume below that \( R_0/R_0 \) is.

Then \( T_0 = 0.226 \frac{mG}{R_0} \).

If, on the other hand, \( R = R_0 \) and \( n' = n_0 \), then for \( R = R_0 \),

\[ T(R) = T_0 \cdot \frac{mG}{1.9} \exp(-\pi R^2 R_0^2) \]

\[ \times \int_0^\theta \exp(-\pi R^2 R_0^2) dR^2 \int_0^\pi \exp(-\pi R^2 R_0^2) dR^2 \]

\[ \times 4 \pi n R^2 dR = T_0 \exp(-0.553 \pi R^2 R_0^2), \]

or

\[ T(R) = T_0 \left( \frac{n(R)}{n_0} \right)^{0.553}, \]

(5)

which corresponds to a polytropic index 0.553. (The polytropic index for an adiabat is 1.5.) In general, the density distribution (1) is largely consistent with the stellar density distribution predicted by polytropic models with indices between 1.5 and 3.5. If \( R_0 = R_0 \) for \( \rho_0 = \rho_0 \), the factor 0.226 in (a) will remain constant as the distribution function (1) is modified within some reasonable limits. Inasmuch as \( T, M, \) and \( R \) are observable, Eq. (4) provides us with an indication of the departure of the last factor from unity (i.e., information about \( R_0/R_0 \) when it is.

5. LUMINOSITY L

The source of x-ray cloud luminosity is almost exclusively electron bremsstrahlung on ions, principally protons. Only at \( T < 10^6 \) K does recombination radiation begin to make inroads, and it becomes the dominant contribution at \( T < 3 \times 10^5 \) K. The ratio of bremsstrahlung to recombination radiation is proportional to the temperature \( T \).

For bremsstrahlung, the emission rate per unit volume per unit time is

\[ W = \frac{2 \Phi_0}{137 \pi} \sum \frac{n_z m_e c^2}{\nu}, \]

(6)

where \( \Phi_0 = (8 \pi/3) (e^2/m_e c^2)^3 \times 0.66 \times 10^{-7} \) cm\(^2\) is the Thomson cross section for light, \( n_z \) is the number of nuclei with atomic number \( Z \) per unit volume, and \( v_s = (2T/m_n)^{1/2} \) is half the mean thermal velocity of the electrons. Here we have \( \Sigma n_z Z^2 = 1.15n. \) The total luminosity of the cloud due to bremsstrahlung is

\[ L = 4 \pi \int_0^\infty \frac{W}{1} R^2 dR = 4 \cdot 10^{-3} \Phi_0 n_z R_0^2 m_e c^2 \]

\[ \times \frac{4 \pi}{\nu} \exp\left(-2.28 \pi \frac{4 \pi}{\nu} \right), \]

(7)

The bremsstrahlung spectral distribution was derived in Ref. 10:

\[ dW/d\omega = \frac{W}{27} \exp\left(-\frac{\hbar \omega}{T}\right). \]

(8)

For \( \hbar \omega \) at most of order \( T \), \( dW/d\omega \) is close to a linear function of \( T \) and \( \omega \) over a relatively small temperature range. We can therefore write for the total luminosity over a frequency interval \( \Delta \hbar \omega \)

\[ \Delta L = L \frac{\hbar \omega}{T} \frac{27}{\hbar \omega} \Delta \hbar \omega, \]

(9)

substituting for \( T \) the temperature at the maximum of the integrand in (7), i.e., \( \omega = \pi/5 \sqrt{R_0} \).

The range of interest has \( \hbar \omega \) in the visible \((1.7 \text{ eV} < \hbar \omega < 3 \text{ eV})\), yielding

\[ \Delta L = L \left(0.5 \times 10^{-4} T_0\right) \times 1.14 \times 10^{-1} T_0, \]

(10)

where \( T_0 \) is in kelvins (here we average over \( T^2 \)); in arbitrary units,

\[ \Delta L = L \left(1.48 T_0 / \hbar \omega \right) \Delta \hbar \omega / 1.32 T_0. \]

O. B. Firsov

676 JETP 79 (5), November 1994
6. PHOTON AND ELECTRON MEAN FREE PATH. TRANSPARENCY, COOLING

The photon mean free path is \( l_{\gamma} = (\pi n_{\gamma} e^2) \), or, according to (1),
\[
l_{\gamma} = (n_{\gamma} e^2) \exp(\pi p^2 / R_{\gamma}^2).
\]
(11)

The transparency can be determined on the basis of the fraction of the radiative flux passing through the cloud at an impact parameter \( p \) without being scattered:
\[
\gamma = \exp \left(-\Phi_{\gamma} \int_{-\infty}^{\infty} \exp(-\pi(p^2 + x^2)) \, dx \right).
\]
(12)
where \( \Phi_{\gamma} = 0 \). Information about the radial distribution of intensity is only available if \( \theta = \theta(0) \leq 1 \).

The electron and proton mean free path is
\[
\alpha_{e} = 10^{16} \text{ cm}^{-2},
\]
(13)
so that if \( R_{e} / l_{\gamma} \) is normally at most of order 1, \( R_{e} / l_{\gamma} \) will always be much greater than 1, and \( R_{e} > 10^{-10} \text{ cm} \), \( n_{e} > 10^{-3} \text{ cm}^{-3} \).

The cooling of x-ray clouds will usually take place essentially only via bremsstrahlung at a rate \( L \). The rate of conductive losses may be comparable. Setting \( L = 10^{-13} \text{ erg} \cdot \text{cm}^{-2} \cdot \text{s}^{-1} \), for \( T = 10^{6} \text{ K} \), \( \alpha_{e} = 10^{16} \text{ cm}^{-2} \),
\[
\alpha_{e} \approx 10^{16} \text{ cm}^{-2},
\]
and as a result of that radiation, rather than cooling down, the cloud heats up (evaporative energy losses are much smaller, hence: \( L / \alpha_{e} \approx 10^{13} \text{ K} \)).

The electron and proton mean free path is
\[
\alpha_{e} = 10^{16} \text{ cm}^{-2},
\]
(13)
so that if \( R_{e} / l_{\gamma} \) is normally at most of order 1, \( R_{e} / l_{\gamma} \) will always be much greater than 1, and \( R_{e} > 10^{-10} \text{ cm} \), \( n_{e} > 10^{-3} \text{ cm}^{-3} \).

The photon mean free path is
\[
\alpha_{e} = 10^{16} \text{ cm}^{-2},
\]
(13)
so that if \( R_{e} / l_{\gamma} \) is normally at most of order 1, \( R_{e} / l_{\gamma} \) will always be much greater than 1, and \( R_{e} > 10^{-10} \text{ cm} \), \( n_{e} > 10^{-3} \text{ cm}^{-3} \).

The neutral x-ray clouds will usually take place essentially only via bremsstrahlung at a rate \( L \). The rate of conductive losses may be comparable. Setting \( L = 10^{-13} \text{ erg} \cdot \text{cm}^{-2} \cdot \text{s}^{-1} \), for \( T = 10^{6} \text{ K} \), \( \alpha_{e} = 10^{16} \text{ cm}^{-2} \),
\[
\alpha_{e} \approx 10^{16} \text{ cm}^{-2},
\]
and as a result of that radiation, rather than cooling down, the cloud heats up (evaporative energy losses are much smaller, hence: \( L / \alpha_{e} \approx 10^{13} \text{ K} \)).

7. QUASISTEADY LIFETIME OF AN X-RAY CLOUD

The radiation from an x-ray cloud is fueled by its conduction and as a result of that radiation, rather than cooling down, the cloud heats up (evaporative energy losses are much smaller, hence: \( L / \alpha_{e} \approx 10^{13} \text{ K} \)).

The cooling of x-ray clouds will usually take place essentially only via bremsstrahlung at a rate \( L \). The rate of conductive losses may be comparable. Setting \( L = 10^{-13} \text{ erg} \cdot \text{cm}^{-2} \cdot \text{s}^{-1} \), for \( T = 10^{6} \text{ K} \), \( \alpha_{e} = 10^{16} \text{ cm}^{-2} \),
\[
\alpha_{e} \approx 10^{16} \text{ cm}^{-2},
\]
and as a result of that radiation, rather than cooling down, the cloud heats up (evaporative energy losses are much smaller, hence: \( L / \alpha_{e} \approx 10^{13} \text{ K} \)).

8. EXPRESSIONS FOR \( n_{0} \), \( S \), \( T \), AND \( \theta \) IN TERMS OF THE OBSERVABLES \( M \), \( T_{0} \), \( L \), AND \( R_{0} \)

We have Eqs. (2), (4), (4a), (7), and (17) with which to determine these eight quantities, as well as the definition of \( S \) and as the ratio of missing mass to plasma mass. At \( t_{0} = R_{0} \), this enables us to express all of these quantities in terms of the observables \( M \), \( T_{0} \), \( L \), and \( R_{0} \), and in certain instances \( R_{0} / t_{0} \). Taking \( M = (1 + S) n_{0} R_{0}^{2} \) [cf. (2)], \( L_{0} = 0.226 G m(M/R_{0}) \) [cf. (4a)], and
\[
L_{0} = 1.25 \times 10^{-5} \Phi_{\gamma} R_{0}^{2} \, T_{0}^{4} / m_{e} c^{2} - \text{cf. (7)}.
\]
we obtain in cgs units
\[
n_{0} = 2.84 \times 10^{-17} L_{0} T_{0}^{4} / m_{e} c^{2},
\]
(18)

9. THE NATURE OF THE DARK MATTER

It was shown in Refs. 1 and 2 that the hypothetical particles invoked to explain dark matter should have a mass of order \( 10^{19} \text{ eV} \). Particles of mass \( 10^{-10} \text{ eV} \) should either decay or annihilate with their antiparticles before the universe reaches critical density. Particle decay, however, is inconsistent with the existence of the solar system, and furthermore with the observed \( \text{C}^{57}/\text{C}^{57} \) ratio, leading to too small a cosmological time \( t < 10^{19} \text{ yr} \) (see also Ref. 4). Annihilating particles should have a mass \( m_{p} > 0.1 \text{ eV} \), and such particles would constitute most of a star's mass (indeed, if the particle products of annihilation interacted strongly with baryonic matter, the stars would explode).

Particles of mass \( 10^{-10} \text{ eV} \) have a gravitational instability radius greater than that of baryonic matter. By the time of radiation decoupling from matter, \( t_{0} = 10^{5} \text{ sec} \), perturbations of the photon gas (which prior to \( t_{0} \) is tightly bound to matter (plasma)), and thus of these same particles, which interact gravitationally with photons and matter, are necessarily no greater than \( 3 \times 10^{-3} \text{ cm} \) on scales \( L \). At that time, \( t_{0} = 3 \times 10^{-3} \text{ cm} \).

Perturbations cannot develop on a scale less than \( R_{0} \), and oscillatory perturbations exist for conventional matter. Perturbations of hypothetical particles on a scale less than or equal to \( R_{0} \), are lacking, however, since they collide neither with one another nor with baryonic matter; the pressure concept is speculative here, at best. All perturbations are rapidly
smoothed out. As long as perturbations of baryonic matter are at most of order 1, no hypothetical particles take part. From that point on, their development slows. It is scarcely likely that this sort of dark matter could show up in x-ray clouds. In any event, despite diligent ongoing searches, no hypothetical particles that might account for the putative dark matter have thus far been detected. Before embracing exotic hypotheses, therefore, we ought to attempt to rule out inconsistencies relating to the baryonic nature of dark matter, as was done in Ref. 1. Here we also adopt the baryonic explanation for dark matter.

10. PLAUSIBILITY OF BARYONIC DARK MATTER IN X-RAY CLOUDS

According to our previous work,2 dark matter of baryonic origin consists principally of large-scale parcels of condensed matter. The evidence provided by meteorites and asteroids3 is that the mass distribution of particles of condensed matter looks like f(M) ∝ M⁻². If the mean density of matter in the universe is 5 × 10⁻³ g cm⁻³ and the mean mass of particles is 10⁻¹⁹ g, then 25% of the mass of dark matter resides in particles whose mass is less than 20 g. 50% is less than 10⁴ g and 75% is less than 10⁸ g. Out of all photons that are absorbed, half are absorbed by particles of dark matter with mass less than 10⁻¹² g, and 75% by particles of mass less than 10⁻⁹ g. Thus, most of the mass in dark matter is essentially unobservable optically.

The condensed matter probably consists mainly of C, N, and O. Their compounds with one another and with H, and to a lesser extent Ne and heavier elements, particularly S, Si, Fe, and Mg. At the prevailing temperatures in space, T < 10⁴ K, all are either in the solid or liquid state. Hydrogen probably constitutes ~10% of all dark matter by mass (H₂O, hydrocarbons starting with CH₄, NH₃, and other compounds).

At this point, we may ask whether such dark matter can survive temperatures ~10⁷ K in the plasma of an x-ray cloud. To find out, we must calculate the equilibrium temperature of a surface irradiated by bremsstrahlung and electrons at the center of the cloud.

The photon flux at the center of the cloud is

\[ q_\text{ph}\left(\theta, \phi\right) = \frac{4\pi}{v_e} \int_0^{\pi/2} r^2 W(\theta) \sin \theta d\theta, \]

where \( r = \sqrt{2Tr/M} \), \( W(\theta) = \int_0^\theta \frac{W(\theta) \sin \theta d\theta}{\int_0^\pi W(\theta) \sin \theta d\theta} \), and \( R = 1.15 \times 10^{15} \text{cm} \). The calculation was carried out as in Ref. 13, but assuming that the dark matter is baryonic. At the present time, the power per unit area transferred by the electron flux is

\[ q_\text{e} = 27T r_\text{ph} \sqrt{2Tr/M} n_\text{e}, \]

where \( r_\text{ph} = (3.55 \times 10^{11}) \text{erg cm}^2 \text{g}^{-2} \text{s}^{-1} \) for the electron energy spectrum in the very last expression. The power per unit area transferred by the electron flux is

\[ q_\text{e} = 0.7 \times 10^{31} \text{erg cm}^{-2} \text{s}^{-1}, \]

Adopting the values \( L = 10^4 \text{erg cm}^{-2} \text{s}^{-1} \), \( T = 2 \times 10^9 \text{K} \), \( M = 10^6 \text{M}_\odot \), and \( q_\text{ph} = 10^3 \text{erg cm}^{-2} \text{s}^{-1} \) as before, we obtain \( q_\text{e} = 4 \times 10^{-3} \text{erg cm}^{-2} \text{s}^{-1} \) and \( q_\text{ph} = 0.84 \text{erg cm}^{-2} \text{s}^{-1} \).

In this example, electrons produce most of the heating. Note, however, that the electrons do not give up all their energy—much less than that, in fact, so that the numerical code with \( q_\text{e} = 0.84 \text{erg cm}^{-2} \text{s}^{-1} \) yields too high a surface temperature, which is obtained by using \( q_\text{e} = 5.7 \times 10^{-7} \text{erg cm}^{-2} \text{s}^{-1} \), where \( T \) is in kelvins (Stefan-Boltzmann law). The dark-matter particle temperature obtained is 11 K. Since we are taking a fourth root here, the particle temperature remains low over a wide range of \( L \), \( T \), and \( M \). Thus, baryonic matter can easily exist in x-ray clouds.

We still need to assess the sputtering of dark-matter particles by protons with energies ~1 keV and \( n_\text{e} \sim 1 \). Proton velocities at \( n_\text{e} = 1 \) are ~3 × 10⁶ cm s⁻¹, and the sputtering coefficient is of order unity. One atomic layer, containing some 10⁸ atoms cm⁻², will erode in ~3 × 10⁻¹⁷ sec. Since half the dark-matter particles have mass greater than 10⁻⁸ g, containing more than 10⁸ atomic layers, more than 10³⁸ yr would be required. The first to be eroded would be dust particles with masses up to 10⁻⁴ g, which would be destroyed in 10⁶ yr. It must be borne in mind, however, that previously sputtered atoms would be redepositing on the dust particles at the same time, so that there should be no significant change in the amount of dark matter in the plasma.

11. ORIGIN OF DARK MATTER. THE FIRST STARS

We assume that at the present time, the mean density of matter in the universe is close to the critical value \( \rho_0 = 5 \times 10^{-30} \text{g cm}^{-3} \), and that it is essentially all of baryonic origin. At the present time, then, \( \rho_0 = (6 \pi G \rho_0)^{1/3} = 1.27 \text{ Gyr} \).

By the time radiation decoupled from matter \( (t = 10^{11} \text{sec}) \), at which point the matter, due to the recombination of plasma \( (T = 3200 \text{K}) \), became transparent to photons, the minimum gravitationally unstable mass of baryonic matter was \( (5 \times 10^{-4} - 10^{-5}) \text{M}_\odot \) (in previous work, due to the enormous heat capacity of the photon gas, the Jeans mass was identified with the reciprocal of the gravitational instability wave number,

\[ R_j = \frac{1}{M_j^{1/2}} = \frac{1}{v_r^{1/2}} \text{ M}_\odot, \]

where \( v_r = (T/M)^{1/2} \) is the isothermal speed of sound. The corresponding result obtained in that work was \( M_j = 2 \times 10^{10} \text{M}_\odot \). A better figure, however, is probably \( R_j = \frac{1}{2} \text{M}_\odot \), which yields \( M_j = 5 \times 10^9 \text{M}_\odot \). That value remained unchanged as long as the temperature of the baryonic matter was approximately equal to that of the photon gas which, despite the very low degree of ionization, continued to heat the baryonic matter for a long time.

The heat capacity of the photon gas is 10⁶ times that of the baryonic matter. Nevertheless, the baryonic matter ultimately cools more rapidly (in the limit as \( T_j \approx 300 \text{K} \), \( T_j = 1.36 \times 10^{-7} T_\odot \), with \( T \) in kelvins). The calculation was carried out as in Ref. 13, but assuming that the dark matter is baryonic and that the mean density of the universe at the
present time is close to the critical value \( \rho_c = \rho_0 = 5 \times 10^{-30} \text{ g cm}^{-3} \); by contrast, it was assumed in Ref. 13 that \( \rho = 10^{-31} \text{ g cm}^{-3} \).

Fluctuations with \( M > M_c = 5 \times 10^{-11} \text{ M}_\odot \) cannot develop prior to \( t_c = 10^{13} \text{ sec} \). The time required for small perturbations to develop in the primordial matter (12H+He) is \( t > a^2 \times 10^{13} \text{ sec} \), where \( a \) is the magnitude of the relative density perturbation \( \delta \rho / \rho - 1 \) prior to cessation of the initial Hubble expansion.

This is followed by a contraction stage and the formation of stars. If the mass of a nascent star very slightly exceeds \( M_c \), the contraction and heating of the star may not precede the temperature at which nuclear reactions are ignited, nor even at which plasma is produced, and expansion may begin anew. A description of this process is outside the scope of the present paper: the theory was worked out beginning with the work of Ritter (1883), Emden (1907), and others.14 In any event, the time for a star to form is limited at the lower range,15 and also it should be proportional to \( a^{-1} \).

The foregoing all applies to masses greater than \( 5 \times 10^9 \text{ M}_\odot \). As for \( a_1 \), it has only recently become possible to detect fluctuations in the temperature of the microwave background—reaching back to \( t_c = 10^{10} \text{ sec} \)—at the 3-\( \sigma \) level, and only in 1992 were fluctuations detected at \( \Delta T / T = 10^{-6} \) [\( \Delta \phi / \Phi = 4 \Delta T / T \)] over angular distances of 0.18°, which for \( t_c = 10^{13} \text{ sec} \) corresponds to linear distances of 3 \( \times 10^{22} \text{ cm} \) (and which is less than \( R_p = 10^{25} \text{ cm} \) for \( m_p e^2 / 10^7 \text{ eV} \)). This result is consistent with the photon mean free path for \( t_c = 10^{13} \text{ sec} \). Fluctuations are obviously impossible at smaller scales. Since radiation and matter were tightly coupled at \( t_c < 10^{13} \text{ sec} \), these same fluctuations also affected the density of matter.

The next question that arises has to do with just how perturbations depend on the linear size of a region or its mass. The answer is unclear—possibly as \( M^{-1/2} \), possibly as \( M^{-1} \) (or have other possibilities been ruled out)?15 The first answer seems more natural, but the second is somewhat more consistent with reality, insomuch as there has been no indication of the past existence of stars with \( M > 10^6 \text{ M}_\odot \), and based upon other considerations as well. At \( t_c < 10^{13} \text{ sec} \), the mass corresponding to \( \Delta \phi = M = 10^{10} \text{ g} \). If the corresponding perturbations were detected at a level \( \Delta \phi > 10^{-13} \), and if the perturbations went as \( \Delta \phi \propto M^{-2} \), then for a minimal Jeans mass \( M = 5 \times 10^9 \text{ M}_\odot = 10^{10} \text{ g} \), we might expect \( a_1 \), which would lead to a huge number of stars of that mass in the early universe, 10^36 years ago (\( t_e = 10^{13} \text{ sec} \)). If \( a_1 \propto M^{-1/2} \), then for \( M \) we obtain \( a_1 \propto 4 \times 10^{-4} \), although this too leads to similar stars by \( t_e = 10^{13} \text{ sec} \). Perturbations probably depend on mass even more weakly. But by that time, the temperature of baryonic matter would have fallen to 3.4 K, and \( M \) to 2 \( \text{M}_\odot \). This then means that somewhat earlier (\( t_e = 10^9 \text{ yr} \) perturbations led to the formation of stars with masses of order \( 10^{-10} \text{ M}_\odot \), since \( a_1 \) would be at least 0.1 for such masses. Consequently, the first generation consisted almost entirely of massive stars (\( \sim 100 \text{ M}_\odot \)). Such stars would have fully evolved after \( 10^9 \text{ yr} \), and ended their lives in a supernova explosion, contaminating the interstellar medium with heavier elements (mainly C, N, and O) primarily in the condensed state (see below). The temperature of the microwave background would have been approximately 70 K at that point.

Subsequent generations of stars would then, to a large extent, have been formed in gas and dust clouds, and would also have been more massive than the sun (\( > 2 \text{M}_\odot \)). They evolved in about \( 10^7 \text{ yr} \), and also for the most part became supernovae. After the explosion, a neutron star with 1 \( \text{M}_\odot \) was ejected as dust and gas (mostly H, He, and H_2), i.e., about 90% of the star’s initial mass. The expansion velocity of this shell was more than \( 10^3 \text{ cm sec}^{-1} \), corresponding to an initial temperature of order \( 10^3 \text{ K} \). The assumed initial radius of the ejected shell was approximately \( 10^{13} \text{ cm} \), and its volume was \( 10^{23} \text{ cm}^3 \), corresponding to an initial heavy-element concentration of \( n_e = 10^{28} \text{ cm}^{-3} \). The expansion velocity of the shell can be assumed constant at \( v_{\text{esc}} = 10^9 \text{ cm sec}^{-1} \).

Now let the concentration fall off as \( R^{-3} \), and the temperature as \( R^{-5} \), where \( R \) is the mean radius of the expanding shell. We then obtain for the radius of the ball of condensed elements

\[
4 \pi^2 \rho_l (\text{d}R / \text{d}t) = 4 \pi^2 R v_{\text{esc}}^3 \rho_l,
\]

where \( a^3 \) is the volume ascribed to each atom, \( v_{\text{esc}} \) is the mean thermal velocity of the condensed elements, and \( n \) is their number density. Obviously \( n = n_l (R_l / R)^3 \), \( v_{\text{esc}} = v_{\text{esc}} (R_l / R) \), so

\[
\frac{\text{d}r}{\text{d}t} = \frac{v_{\text{esc}} a^3}{r} = n_l a^3 \left( \frac{R_l}{R} \right)^4,
\]

\[
r = n_l a^3 R_l^4 \left( \frac{R_l}{R} \right)^4 \left( \frac{1}{R_l} \right)^3 R_l^3 / R^3
\]

Taking \( n_l = 10^{24} \text{ cm}^{-3} \), \( R_l = 10^{10} \text{ cm} \), \( a^3 = 10^2 \text{ cm}^3 \), and \( R_l = 10^{15} \text{ cm} \), corresponding to \( T = 10^3 \text{ K} \) at the beginning of condensation, we obtain \( a = 0.3 \text{ cm} \).

This would all be the case if the number of condensation centers were less than 1 per 10^23 atoms. It is in fact likely that the number is much greater, and that fine dust is formed much earlier. Almost all of the precipitable matter is exhaled, and dust grains proceed to stick together to form more massive objects, possibly in accordance with the mass distribution function of Ref. 3. Very fine grains are slowed by the interstellar gas (12H+He) and remain within the galaxy itself. Particles above \( \sim 10^{-2} \text{ g} \) (at a mean gas density of \( \sim 0.1 \text{ cm}^{-3} \)) leave for the halo, thereby becoming dark matter. This dark matter probably then comes to thermal equilibrium with the microwave background, and subsequently heats the interstellar gas to the background temperature.

By virtue of the foregoing, most condensed matter—the missing mass—winds up in the galactic halo. Second and third generation late-type stars spawned by clouds of gas and dust (cf. Ref. 1) actually contain 80–90% heavy elements (C, N, O, ...) by mass; the ~2% by mass observed spectroscopically would seem to be a surface effect, since the dust con-
tracts first when a cloud of gas and dust contracts to form a star. Accordingly, the estimated lifetime of a solar-type star should probably be reduced by an order of magnitude, from $-10^{11}$ yr to $-10^{10}$ yr. As for the sun, this argument is favored by neutrino measurements that indicate the $p-p$ cycle ("boron" neutrinos) can account for at most 40% of the solar luminosity. The rest of the energy is derived from the Bethe CNO cycle, with perhaps another $-3\%$ at the center of the sun coming from $\text{He} + \text{C} \rightarrow \text{O} + \gamma 7.16$ MeV. Globular cluster stars, most of which reside outside the galactic disk, may be an exception. These were probably formed earlier in regions not yet contaminated by heavy elements, out of large-scale perturbations ($-10^6 \text{M}_\odot$) that fragmented at $t_r\approx10^{16}$ sec, with $T_r\approx3$ K and $\text{M}_r/\text{M}_\odot$.

Many unsolved problems in physical cosmology are related to the idea of baryonic dark matter or missing mass. This problem has not yielded even to suggestions of unobservable exotic particles, which, more than likely, do not exist. Solving the missing mass problem in a baryonic context instills us with a great deal of hope. In the $D/H$ problem, for example, in addition to the sources of deuterium proposed in Refs. 1 and 2, we now have new information on deuterium production in the solar corona (and therefore other stars). Cosmic rays of stellar origin probably fragment He nuclei into deuterons and neutrons; the latter are absorbed by protons to form new deuterons, releasing characteristic 2.15-MeV $\gamma$-rays, which can then be detected by satellite-borne instruments. The evidence here seems to favor a secondary origin for deuterium in the interstellar medium. To summarize, the idea that dark matter is baryonic in nature is not inconsistent with its being detected in x-ray clouds, and would seem to provide a more attractive way of accounting for dark matter than would hypothetical unobservable particles.

I thank V. A. Vlasov, V. I. Kogan, S. Yu. Luk’yanyov, Yu. V. Martynenko, and B. A. Trubnikov for useful discussions during the course of this work, and T. I. Zhukova for assistance in preparing the paper.

5. A. A. Suchkov, Galaxies: Familiar...and Puzzling [in Russian], Nauka, Moscow (1988).

Translated by Marc Damashek