Plasma-Redshift Cosmology: A Review

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Abstract. The newly discovered and experimentally verified plasma-redshift cross section of photons penetrating hot sparse plasma leads to a new cosmology, which is radically different from the conventional big-bang cosmology. The plasma-redshift cross section is deduced from conventional axioms of physics without any new assumptions. It has been overlooked, because it is insignificant in ordinary laboratory plasmas; but it is important in sparse hot plasmas, such as those in the corona of the Sun, stars, quasars, galaxies, and intergalactic space. The energy that the photons lose in plasma redshift heats the plasma. The deduction of plasma redshift requires that we take into account the dielectric constant more accurately than is usually done. In the Sun, the plasma redshift predicts the observed densities and the temperatures in both the transition zone and in the corona. Plasma redshift predicts the observed intrinsic redshifts of the Sun, stars, quasars, and galaxies, the cosmological redshifts, cosmic microwave background, and cosmic X-ray background. There is no need for: Einsteins cosmological constant Lambda, Big Bang, Cosmic Inflation, Dark Energy, Dark Matter, Black Holes, and Cosmic Time Dilation. Plasma redshift shows also that contrary to general belief, the gravitational redshift in the Sun is reversed when photons move from the Sun to the Earth. This is a quantum mechanical effect. This means also that the photons are weightless in local system of reference. All the many experiments, which have been assumed to prove photons have weight, are meaningless, because in all cases the researchers disregarded the quantum mechanical uncertainty principle. It is essential to use quantum mechanical concepts for deducing plasma redshift and weightlessness of photons. Plasma redshift cannot be derived using classical physics methods. It would, therefore, not exist in the conventional plasma cosmology. Plasma-redshift cosmology, which besides the plasma redshift cross section includes the newly discovered weightlessness of photons shows that there are no black holes (BHs) or super-massive BHs (SMBHs), because the weightless photons accumulate at the centers of BH candidates (BHCs) and SMBH candidates (SMBHCs) and prevent formation of BHs, as shown in the related poster session paper at this conference.

1. Deduction of Plasma Redshift

1.1. The Equation of Motion

In references [Brynjolfsson(2004a-e, 2006a,b] it is shown that plasma-redshift theory gives simpler and more accurate descriptions of cosmological phenomena than the contemporary Big-Bang cosmology.

Plasma redshift of photons interacting with hot sparse plasma has been overlooked in the past. It is in some respect analogous to the energy loss of

fast charged particles through Cerenkov radiation. (Enrico Fermi discovered the explanation for the Cerenkov radiation in 1939 and 1940 long after it had been observed, see Fermi (1939).) In both cases, it is essential that the dielectric constant be taken properly into account. We can treat the interactions of photons with plasma semi-classically, because the exchange effects are not important. At the position of $r = 0$, the equation of motion for a plasma electron, acted upon by the electrical field's Fourier component $E \exp(i\omega t) = (A/\varepsilon) \exp(i\omega t)$, is

$$m\ddot{r} + m\alpha\dot{r} - m\beta_p\dddot{r} + m\omega_q^2 r = e\frac{A}{\varepsilon} \exp(i\omega t), \qquad (1)$$

where on the right side $A = D$ is the modulus of the displacement field, and A/ε is the modulus of the electrical field, see Appendix A. This correct inclusion of the ε in the dynamical equation leads to the plasma redshift. Usually, it has been surmised that we could set $\varepsilon = 1$. The first term on the left side is due to the acceleration of the electron with charge e and mass m. The second term on the left side accounts for the collision damping, where $\alpha = 2/\tau$ and τ is the time between collisions. The third term, the radiation damping term, accounts for the emitted radiation by the electron when it is accelerated in the external field on the right side. We note it by $-m\beta_p\dddot{r}$ rather than $-m\beta_0\dddot{r}$ to make it clear that β_p could deviate from $\beta_0 = 2e^2/(3mc^3)$. If a field with only the frequency ω_0 acts on the electron, we have that $-m\beta_p\dddot{r} = -m\beta_0\dddot{r} = m\beta_0\omega_0^2\dot{r}$. But if several frequencies of the field act on the electron simultaneously, then $-m\beta_p\dddot{r}$ can deviate from $m\beta_0\omega^2\dot{r}$. The collision field can be replaced with the Fourier harmonics of the field of the fast moving electrons in the hot plasma. These Fourier fields cause the plasma electrons to oscillate and lose the radiation energy in a similar way as the incident photon field on the right side of Eq. (1). We can then add the collision term to $-m\beta_p\dddot{r}$, and replace the sum with $-m\beta\dddot{r} = m\beta\omega^2\dot{r}$. The fourth term, $m\omega_q^2 r$, accounts for any "elastic" force that binds the electron to a certain equilibrium position with "eigenfrequency" ω_q.

The principal solution of Eq. (1) is (see Eq. (A15) in Appendix A of Brynjolfsson (2004a))

$$\begin{aligned}
r &= \frac{e}{m} \frac{A/\varepsilon}{\left\{\omega_q^2 - \omega^2 + i\left(\alpha + \beta_p\omega^2\right)\omega\right\}} \exp(i\omega t)) \\
&= \frac{e}{m} \frac{A/\varepsilon}{\left\{\omega_q^2 - \omega^2 + i\,\beta\omega^3\right\}} \exp(i\omega t).
\end{aligned} \qquad (2)$$

This equation may be compared with Eq. (21-19) of reference (Panofsky & Pillips 1956), but take note of the subtle differences.

> **Comment 1.** In Eq. (21-19) of (Panofsky & Pillips 1956), which corresponds to Eq. (2) above, Panofsky and Phillips do not have any collision term and not the factor $1/\varepsilon$. The same applies to Eq. (15.2a) by Becker (Becker 1949). But in Eq. (15.10) of (Becker 1949), Becker considers the collision damping but disregards $1/\varepsilon$ in the numerator. These conventional approximations in references (Panofsky & Pillips 1956) and (Becker 1949) prevent us from discovering the plasma redshift. We emphasize that for deducing the plasma redshift it is essential to use the correct form of Eq. (1) and (2).

1.2. The Dielectric Constant

The polarization in the plasma is given by

$$P(\omega) = N_e\, e\, r, \qquad (3)$$

where N_e is the number of plasma electrons per cm^3, e the electronic charge, and where r, the displacement of each of the electrons is given by Eq. (2). The dielectric constant is defined as

$$\varepsilon(\omega) = 1 + 4\pi \frac{P(\omega)}{(A/\varepsilon)\exp(i\omega t)} = 1 + \frac{4\pi\, N_e\, e\, r}{(A/\varepsilon)\exp(i\omega t)}. \qquad (4)$$

When we insert Eqs. (2) in this expression for the dielectric constant, we get

$$\varepsilon = 1 + \frac{4\pi N_e e^2/m}{\left\{\omega_q^2 - \omega^2 + i\left(\alpha + \beta_p\omega^2\right)\omega\right\}} = 1 + \frac{\omega_p^2}{\left\{\omega_q^2 - \omega^2 + i\left(\alpha + \beta_p\omega^2\right)\omega\right\}}, \qquad (5)$$

where the cyclic plasma frequency is

$$\omega_p = 2\pi\nu_p = \sqrt{\frac{4\pi N_e e^2}{m}} = 5.642 \cdot 10^4 \sqrt{N_e}, \qquad (6)$$

From Eq. (5), we see that for $\omega_q \approx 0$ and $\omega < \omega_p$ the dielectric constant is very different from one. The conventional assumptions that set $\varepsilon = 1$, are, therefore, impermissible approximations for small frequencies of the Fourier harmonics. By contrast, Eqs. (1) and (5) are much more accurate. In addition to the dielectric constant ε, Eqs. (1) and (5) also take the collision damping in the imaginary part, $i(\alpha + \beta\omega^2)\omega$, properly into account.

If we write the complex dielectric constant on the form $\varepsilon = (n - i \cdot \kappa)^2$, we get from Eq. (5) that

$$\frac{2n\kappa\omega}{\varepsilon\bar{\varepsilon}} = \frac{\left(\alpha + \beta_p\omega^2\right)\omega_p^2\omega^2}{\left(\omega_q^2 + \omega_p^2 - \omega^2\right)^2 + \left(\alpha + \beta_p\omega^2\right)^2\omega^2} = \frac{\omega_p^2\beta\omega^4}{\left(\omega_q^2 + \omega_p^2 - \omega^2\right)^2 + \beta^2\omega^6}, \qquad (7)$$

where in the last form, we replaced the damping constant $\left(\alpha + \beta_p\omega^2\right)$ by $\beta\omega^2$. Eq. (7) is deduced for calculating in subsection 1.3 the attenuation of the photons in a plasma. This form of the damping constant is the same as that in section 3 of reference (Brynjolfsson 2004a). The value of β, estimated in subsection 1.5, can be very large in a hot plasma. We show that β is on the order of $(3kT/\hbar\omega_p)\,\beta_0$.

Comment 2. The form of the dielectric constant in Eq. (5) above is similar to the form of the dielectric constant given by Becker (1949) (see Eqs. (25.5) and (25.9) of that source), except that Becker uses the form $\beta_0\omega^2 = 2e^2\omega^2/(3mc^3)$ instead of $[\alpha + \beta_p\omega^2] \approx [\alpha + 2e^2\omega^2/(3mc^3)]$ in the denominator of Eq. (5). (We draw attention to that Eq. (25.6) of (Becker 1949) applies to ponderable matter and does not apply to plasmas.) Becker then proceeds to integrate over the Fourier harmonics as if $2e^2\omega_q^2/(3mc^3)$, was small and a constant (see his Eq. (26.8) of reference (Becker 1949)). That is, he disregards the solutions that could lead to the plasma redshift and the scattering on the plasma frequency. Panofsky and Phillips do the same in reference (Panofsky & Pillips 1956) (see Eq. (21-26) of that source). We, on the other hand, will integrate more accurately over the Fourier harmonics and take into account that the damping factor could be large and that the damping constant varies with the frequency ω as $\beta\omega^2$, because only by using the exact form for the damping of the electrons and the exact integration can we obtain the cross section for the plasma redshift. In the conventional literature on plasmas it is even customary to disregard altogether the imaginary part in Eq. (5) by assuming that $\varepsilon = 1 - \omega_p^2/\omega^2$. *Such approximations can never lead to plasma redshift.*

1.3. The Attenuation of the Photon

We can normalize the energy flux $\breve{S}$ to one photon per second and per cm^2, $\breve{S} = \hbar\omega$, in vacuum.

When using the solutions of Maxwell's equations given in Appendix A, we find that the normalized photon flux at the distance x cm from the source is given by (see Eq. (A21) of section A3 of Appendix A in reference (Brynjolfsson 2004a))

$$\breve{S} = \hbar\omega_0 \frac{\gamma}{2\pi} \int_{-\infty}^{\infty} \frac{n}{\varepsilon\bar{\varepsilon}} \frac{[\exp(-2\kappa\omega x/c)]}{\left\{\gamma^2/4 + (\omega - \omega_0)^2\right\}} d\omega, \tag{8}$$

where ω_0 is the center frequency of the photon in vacuum at x = 0.

The decrease, $d\breve{S} = d\hbar\omega$, in the photon's energy flux $\breve{S} = \hbar\omega$ per dx is given by

$$\frac{d\breve{S}}{dx} = \frac{d\hbar\omega}{dx} = -\frac{\hbar\omega_0\gamma}{2\pi c} \int_{-\infty}^{\infty} \frac{2n\kappa\omega}{\varepsilon\bar{\varepsilon}} \frac{[\exp(-2\kappa\omega x/c)]}{\left\{\gamma^2/4 + (\omega - \omega_0)^2\right\}} d\omega. \tag{9}$$

The quantity γ in Eq. (8) and (9) represents the true quantum mechanical photon width and not the classical photon width $\gamma_0 = 2e^2\omega_0^2/(3mc^3)$, (see Becker's Eq. (14.8) of reference (Becker 1949)). For the dielectric constant given by Eq. (5), this expression shows how the photon energy flux decreases when the photon penetrates the plasma. Eq. (9) is equal to Eq. (A22) in section A4 of Appendix A in reference (Brynjolfsson 2004a). When inserting Eq. (7) into Eq. (9), we have at a distance $x = 0$ that

$$\frac{d\hbar\omega_0}{dx} = -\frac{\hbar\omega_0\gamma}{2\pi c}\int\limits_{-\infty}^{\infty}\frac{\omega_p^2\beta\omega^4}{\left[\left(\omega_q^2+\omega_p^2-\omega^2\right)^2+\beta^2\omega^6\right]\left\{\gamma^2/4+(\omega-\omega_0)^2\right\}}\frac{d\omega}{\quad}, \quad (10)$$

which gives the photon's energy loss per cm at $x = 0$.

When evaluating the integral of Eq. (10) with focus on hot plasmas, we usually assume that:

$$1)\ \omega_q = 0;\quad 2)\ \beta \gg \beta_0 = 6.266\cdot 10^{-24};\quad 3)\ \beta\omega_p \ll 1;$$

$$4)\ \omega_0 \gg \omega_p;\quad \text{and}\quad 5)\ \gamma \ll \omega_0. \tag{11}$$

These conditions for the plasma redshift are usually fulfilled for the plasmas of main interest. For example, for $T \approx 500,000$ K and $N_e \approx 10^9\ cm^{-3}$ in the transition zone to the solar corona, the value of $\beta = (3kT/\hbar\omega_p)\,\beta_0 \approx 1.1\cdot 10^8\beta_0$, $\beta\omega_p \approx 1.23\cdot 10^{-6}$, $\omega_p \approx 1.784\cdot 10^9\ s^{-1}$, and $1/\beta \approx 1.45\cdot 10^{15}$.

For integrating Eq. (10), we use complex integration. We select a path along the ω-axis from $-\infty$ to ∞ and then along a semicircle in the upper half-plane from $+\infty$ to $-\infty$. The integral along the semicircle is equal to zero. The integral along the ω-axis is therefore equal to $2\pi i$ times the sum of the residues in the poles in the upper half-plane.

The poles in the complex plane are obtained from the denominator, which has eight complex roots. The four roots in the upper plane are:

$$\omega = \begin{cases} \mathbf{a} &=& +\sqrt{\omega_p^2 - \dfrac{\beta^2\omega_p^4}{4}} &+& i\,\dfrac{\beta\omega_p^2}{2} \\[2ex] \mathbf{b} &=& -\sqrt{\omega_p^2 - \dfrac{\beta^2\omega_p^4}{4}} &+& i\,\dfrac{\beta\omega_p^2}{2} \\[2ex] \mathbf{c} &=& &+& i\left(\dfrac{1}{\beta} + \beta\omega_p^2 - 2\beta^3\omega_p^4\right) \\[2ex] \mathbf{d} &=& +\,\omega_0 &+& i\,\dfrac{\gamma}{2} \end{cases}. \tag{12}$$

When the relations in Eqs. (11) are well fulfilled, we get

$$\omega = \begin{cases} \mathbf{a} &=& +\,\omega_p &+& i\,\dfrac{\beta\omega_p^2}{2} \\[2ex] \mathbf{b} &=& -\,\omega_p &+& i\,\dfrac{\beta\omega_p^2}{2} \\[2ex] \mathbf{c} &=& &+& i\,\dfrac{1}{\beta} \\[2ex] \mathbf{d} &=& +\,\omega_0 &+& i\,\dfrac{\gamma}{2} \end{cases}. \tag{13}$$

The results of the integration on the right side of Eq. (10) is then

$$\frac{d\hbar\,\omega_0}{dx} = -\frac{\hbar\,\omega_0\,\gamma}{2\pi\,c}\,2\pi i\,[\,\mathrm{Res}\,(\mathrm{a}) + \mathrm{Res}\,(\mathrm{b}) + \mathrm{Res}\,(\mathrm{c}) + \mathrm{Res}\,(\mathrm{d})\,]$$

$$= -\frac{\hbar\,\omega_0\,\gamma_0\,\omega_p^2}{c\,\omega_0^2}\left[\frac{\gamma}{4\,\gamma_0} + \frac{\gamma}{4\,\gamma_0} + \frac{\gamma}{2\,\gamma_0}\frac{\left(1 - 1/\left(\beta\,\omega_0\right)^2\right)}{\left(1 + 1/\left(\beta\,\omega_0\right)^2\right)^2} + \frac{1}{\left(1 + \left(\beta_0\,\omega_0\right)^2\right)}\right].$$
(14)

As above, we use the notation γ for the actual quantum mechanical width, which includes the collision width, of the incident photon, while the classical photon width is given by

$$\gamma_0 = \beta_0\omega_0^2 = \left(\frac{2}{3}\frac{e^2}{mc^3}\right)\omega_0^2 = 6.266\cdot 10^{-24}\,\omega_0^2,$$
(15)

From Eqs. (6) and (15) we get that the factor in front of the brackets in Eq. (14) is

$$\frac{\hbar\,\omega_0\,\gamma_0\,\omega_p^2}{c\,\omega_0^2} = \hbar\,\omega_0\,6.652\cdot 10^{-25}\,N_e$$
(16)

where $6.652\cdot 10^{-25}$ is the Compton scattering cross section.

The two roots a and b of Eq. (13) lead to the two first terms inside the brackets of Eq. (14); and the **root c** leads to the third term inside the brackets of Eq. (14), the plasma-redshift term. These three terms are due to the dielectric constant $\varepsilon \neq 1$. The conventional calculations have disregard these three roots. The **root d** leads to the last term inside the brackets of Eq. (14); that is, the conventional Compton scattering term.

The roots a and b in Eqs. (12) and (13), correspond to Stokes scattering or Raman scattering. The plasmas are usually in a thermodynamic equilibrium. The quantum mechanical treatment shows then that we have about equal number of negative and positive oscillators in the plasma. The energy loss and gains in the Raman scattering then averages out to a zero and we observe no Raman Effect. However, the angular spread caused by this Raman scattering, although very small, starts to be detectable in case of the most distant supernovas, as predicted by Eq. (52) of Brynjolfsson (2004a).

The imaginary root c in Eqs. (12) and (13) is a pure absorption term. It leads to the third term, the plasma-redshift term, inside the brackets of Eq. (14). Although, the real part of the frequency in the **root c** is zero, the imaginary part or the collision damping in the electron plasma results in energy losses of the incident photon. At high temperatures, the value of $1/\beta$ can be equal to or exceed the frequency ω_0. This **imaginary root c**, which is important only in a hot sparse plasma, has not been considered before.

Comment 3. When $\beta\omega_0$ is very small, the plasma-redshift term in Eq. (14) is nearly zero. But when the product $\beta\omega_0$ increases beyond one, this plasma-redshift term results in the large plasma-redshift cross section. This plasma-redshift term is proportional to the quantum mechanical photon width γ, which includes the collision broadening effects. The variation in the photon width, γ, from line to line explains the variation of the solar redshift from line to line as given by Eq. (27) below. It also explains the variation in the limb effect from line to line. The conventional literature could not explain these details.

Comment 4. Quantum mechanical calculations often disregard the dielectric constant. For example, in his excellent monograph Heitler (1954) disregards the dielectric constant when he calculates the cross section for Double and Multiple Compton scattering. When one of the outgoing photon is very soft (which results in the so-called infrared problem), the interaction never involves only one electron (not even in the most sparse plasmas of intergalactic space). This is contrary to Heitler's assumption in his calculations. Heitler's estimate of the cross section for Double and Multiple Compton scattering does not apply therefore to plasmas, although he solved the mathematical problem, based on the photon's interaction with one single electron, correctly; see (Brynjolfsson 2004a).

1.4. Collision Frequency

The damping in the oscillations of the plasma electrons is much larger than the conventional damping caused by the photon field alone; that is, $\beta\omega^2 = \left(\alpha + \beta_p\omega^2\right) \gg \left(\beta_p\omega^2\right)$. The collision damping α in Eq. (1) is often equated with $2/\tau$, where τ is the time between collisions, see for example Eq. (15.6) to Eq. (15.10) of reference (Panofsky & Pillips 1956).

From the stopping theory of charged particles, we know that the energy absorbed per colliding electron in a small increment dp of the impact parameter p is proportional to $\{[xK_1(x)]^2 + [(x/\gamma_r)K_0(x))]^2\}(dp/p)$, where $K_0(x)$ and $K_1(x)$ are the modified Bessel functions of zero and first order, $x = p\omega_0/\gamma_r v$, and where γ_r in this case is the relativistic factor $\gamma_r = 1/\sqrt{1 - v^2/c^2}$, which is not to be confused with the notation γ for the radiation damping. (Niels Bohr deduced these relations in 1913 and 1915 (Bohr 1915)). In the following, we will assume that the relativistic factor, γ_r, is about equal to one, which eliminates any confusion about the notations. The quantity within the braces is about equal to one for $x < 1$, and for $x \geq 1$, it falls off exponentially. The cross section for energy absorption is then about equal to $\pi p^2 = \pi(v/\omega_0)^2$. If we multiply the cross section by the electron flux per second, $N_e v$, we get that the collision frequency in pertinent collisions with the electron in the plasma from all the other electrons is given in the next equation, where the kinetic energy of the electron is $(1/2)mv^2 = (3/2)kT$.

$$1/\tau = \pi \, \frac{v^2}{\omega_0^2} \, N_e \, v \approx \pi \, \frac{v^3}{\omega_0^2} \, N_e = \pi \, \frac{(3kT)^{3/2}}{m^{3/2} \, \omega_0^2} \, N_e \,, \qquad (17)$$

In a more exact evaluation, we would have to weigh the expression with the actual energy distribution of the plasma electrons. For the rough estimates, which are adequate for this illustration, such exact evaluation is not needed. For $N_e \approx 10^9$ and $T \approx 5 \cdot 10^5$ K, which correspond to the condition low in the solar corona, we get that $2/\tau = 6.81 \cdot 10^{35}/\omega_0^2 \; s^{-1}$. For 5000 Å light, the value of $\omega_0^2 \approx 1.42 \cdot 10^{31}$ and the value of $2/\tau \approx 4.8 \cdot 10^4$. The corresponding classical value of the photon width is: $\gamma_0 = \beta_0\omega_0^2 = 6.266 \cdot 10^{-24} \, \omega_0^2 \approx 8.9 \cdot 10^7 \gg 2/\tau$. The quantum mechanical width, γ, of the photon from the undisturbed atom differs from the classical width, but usually we have $0.1\gamma_0 \leq \gamma \leq 10\gamma_0$. We see thus that for the photon frequencies of main interest, the Compton scattering in the corona is not affected significantly by the collision damping $2/\tau$.

For the roots a, b, and c, the important frequencies are on the order of or less than the plasma frequency, $\omega_p = 5.64 \cdot 10^4\sqrt{N_e}$. When we insert ω_p for ω_0 in Eq. (17), we get: $2/\tau \approx 6.05 \cdot 10^8 \, T^{3/2}$, which for the hot sparse plasmas of our interest is very large compared with ω_p. For these low frequencies, the collision damping is very important, and we must consider the collective interactions. The plasma-redshift frequencies are even lower than ω_p. In the next section, we will show how we can quantify the collision damping more accurately when we treat the phenomena quantum mechanically.

1.5.   Estimates of β for Low Frequencies

When the plasma is disturbed, the forces within the plasma will result in characteristic oscillations with eigenfrequency ω_p, given by Eq. (6). We see this frequency in the denominator of Eqs. (7) and (10). For $\omega_q = 0$, the plasma frequency, ω_p, is the principal frequency for absorption. Each electron will oscillate as a classical oscillator with a restoring force proportional to the displacement r. The electron will oscillate as a classical oscillator due to the polarization with the restoring force $m\ddot{r} = -kr$ and the frequency $\omega = \sqrt{k/m}$. For each plasma electron, we have that $k = 4\pi N_e e^2$. The force $m\ddot{r} = -kr = -4\pi N_e e^2 r$ corresponds to the polarization given by Eq. (3). When we solve the classical equation $m\ddot{r} = -kr = -4\pi N_e e^2 r$, the solution is that of a classical harmonic oscillator with the frequency $\omega_p = \sqrt{k/m} = \sqrt{4\pi N_e e^2/m}$, as defined by Eq. (6). For $\omega_q = 0$, it is a characteristic eigenfrequency for each plasma electron, as Eq. (10) shows.

The forced oscillations of the electron will result in the usual radiation damping. The positive ions will also act like harmonic oscillators, but their radiation damping is much smaller, because of their larger mass.

We can treat the electron plasma quantum mechanically as in section 3 of reference (Brynjolfsson 2004a). The plasma consists of great many oscillators. The Hamiltonian for each oscillator is given by

$$H_0 = -\frac{\hbar^2}{2m}\nabla^2 + \frac{1}{2}m\omega_p^2 r\bar{r} \tag{18}$$

The energy levels of the oscillator in the plasma are given by (as in Eq. (9) of reference (Brynjolfsson 2004a))

$$E_\Lambda = E_{n,\,l} = \left(\Lambda + \frac{3}{2}\right)\hbar\left(\omega_p - i\frac{\beta\omega_p^2}{2}\right), \tag{19}$$

where $\Lambda = 2n + l$ and n is the principal quantum number, and l is the angular quantum number, and where each of the quantum numbers n and l can be any whole number and zero: $0, 1, 2, 3, ., ., ., ,$. We could also set $\Lambda = n_x + n_y + n_z$, where n_x, n_y, and n_z are the quantum numbers in the linear oscillations along the coordinate axes. The imaginary value in the energy level is included to indicate the finite lifetime of the state. (We often calculate first the eigenstates and then the stability of the states. The complex notation for the energy levels can be considered a shorthand notation for this two-step process.) For simplifying the analysis, the effect of the magnetic fields is disregarded. The modification by magnetic field can be treated separately, see section 4 of (Brynjolfsson 2004a).

When the magnetic field is zero, the states are degenerate; that is, several states can have the same energy for $\Lambda > 1$. For example, for the states $\Lambda = 4$, we have $(n, l) = (2, 0)$, $(1, 2)$, or $(0, 4)$. These three states in turn have the multiplicity of $2l + 1$, or 1, 5, and 9, respectively, for a total of 15 states. More generally, for each quantum number Λ, the total number of states is $(\Lambda + 1)(\Lambda + 2)/2$, all with the same energy.

The damping of the electrons is derived from the transition rate of the electrons from one state to another in the plasma (see Eq. (10) of reference [1]). The Einstein A coefficients for an oscillator are

$$A_{\Lambda \to \Lambda - 1} = \frac{4e^2 \omega_p^2}{2mc^3} \left[\frac{n_x + n_y + n_z}{3} \right] = \frac{2e^2 \omega_p^2}{3mc^3} \Lambda = \Lambda \beta_0 \omega_p^2 = \beta \omega_p^2 \qquad (20)$$

where $\Lambda = n_x + n_y + n_z$.

The radiation damping constant is thus found to be $\beta = \Lambda \beta_0$ for each electron, where Λ is the quantum number for the energy level of the electron (see Eq. (10) of reference (Brynjolfsson 2004a)). In hot plasma, the energy levels will be highly excited and have approximately thermal distribution. When a classical electron oscillates in the field of one field harmonic, $(A/\varepsilon) \exp(i\omega_q t)$, the radiation losses will be proportional to $(2e^2/3mc^3)\omega_q^2$. But in the collisions field from the surrounding electrons, the radiation losses can be very different. Eq. (19) shows that the electrons in the plasma are vibrating with high frequencies resulting in radiation losses that are Λ times higher or $\beta \omega_p^2 = \Lambda \beta_0 \omega_p^2$, where Λ is on the order of $\Lambda = 3kT/\hbar\omega_p$.

1.6. The Distribution for the Damping Factor β

The estimates of the collision damping of electrons in a sparse hot fully ionized hydrogen plasma is relatively simple. All the atoms are fully ionized and the electrons relatively free. The quantum mechanical estimates of the energy levels and the transition rates between the levels, which are equal to the damping in the plasma, are known and given by Eq. (20).

We have that

$$\beta = \Lambda \beta_0 \approx \frac{3kT}{\hbar \omega_p} \beta_0 , \qquad (21)$$

where the total energy of the state is $3kT$. For $T = 5 \cdot 10^5$ K and $N_e = 10^9 \, cm^{-3}$ the ratio $3kT/(\hbar \omega_p) = 2.07 \cdot 10^{-10}/1.88 \cdot 10^{-18} = 1.1 \cdot 10^8 = \Lambda$. We have therefore,

that in this case, which corresponds to the middle of the transition zone to the solar corona, that $\beta = 1.1 \cdot 10^8 \beta_0 = 6.89 \cdot 10^{-16}$. For 5000 Å light, the value of $\beta \omega_0 = 2.6$, and the value of $(1 - 1/(\beta\omega_0)^2)/(1 + 1/(\beta\omega_0)^2)^2 = 0.65$. The plasma redshift term in Eq. (14) is therefore significant for $\lambda \approx 5000$ Å photons in the middle of the transition zone to the solar corona. A better estimate takes into account the statistical distribution of the oscillators. At these high temperatures, the Boltzmann, Fermi-Dirac, and Bose-Einstein statistics all give the same thermal distribution.

When we weigh the plasma-redshift term in Eq. (14) by the thermal distribution, we get that the oscillator strength function $F_1(a)$ as a function of $a = \hbar\omega_p/(\beta_0\omega_0 kT) = 3.6509 \cdot 10^5 \lambda_0 \sqrt{N_e}/T$, is given by Table 1 below; see subsection 3.2 of (Brynjolfsson 2004a) for the mathematical deduction.

Table 1. $F_1(a)$ as a function of $a = \hbar\omega_p/(\beta_0\omega_0 kT)$

a	$F_1(a)$	a	$F_1(a)$	a	$F_1(a)$	a	$F_1(a)$
0.0	1.000	1.0	0.571	2.0	0.228	6.0	-0.070
0.1	0.990	1.1	0.527	2.2	0.183	7.0	-0.073
0.2	0.962	**1.163**	**0.500**	2.4	0.144	8.0	-0.071
0.3	0.921	1.2	0.485	2.6	0.111	9.0	-0.067
0.344	**0.900**	1.3	0.445	**2.671**	**0.100**	10.0	-0.061
0.4	0.872	1.4	0.407	2.8	0.082	20.0	-0.024
0.5	0.821	1.5	0.372	3.0	0.057	40.0	-0.0071
0.6	0.769	1.6	0.339	3.5	0.010	50.0	-0.0047
0.7	0.717	1.7	0.309	3.633	0.000	100.0	-0.0012
0.8	0.667	1.8	0.280	4.0	-0.022	200.0	-0.0008
0.9	0.618	1.9	0.253	5.0	-0.057	∞	-0.0000

The redshift is then proportional to $\Delta N_2 = N_e F_1(a)$, which is thus a measure of the oscillator strength of the redshift term, the third term inside the brackets in Eq. (14). Table 1 shows that $F_1(a)$ is close to 1 for small values of a, or for high frequencies ω_0, or for small wavelengths λ_0 in cm of the incident radiation. From the definition $a = \hbar\omega_p/(\beta_0\omega_0 kT) = 3.651 \cdot 10^5 \lambda_0 \sqrt{N_e}/T$ above, we have that the photon's cut-off wavelength is

$$\lambda_0 = \frac{2\pi c}{\omega_0} = 2.739 \cdot 10^{-6} \, a \, \frac{T}{\sqrt{N_e}}. \tag{22}$$

From Table 1, we can find the oscillator strength for a given value of a. For example, for $a \leq 1.163$, we have that $F_1(a) \geq 0.5$ or that the oscillator strength is $\geq 50\,\%$, where the 50% cut-off wavelength, $\lambda_{0.5}$, for the redshift is determined by inserting 1.163 for a into Eq. (22). We get

$$\lambda_{0.5} = 2.739 \cdot 10^{-6} \cdot 1.163 \frac{T}{\sqrt{N_e}} = 3.185 \cdot 10^{-6} \frac{T}{\sqrt{N_e}} \, cm = 318.5 \frac{T}{\sqrt{N_e}} \, \text{Å}. \tag{23}$$

The 90\,\% and 10\,\% oscillator strengths are obtained for a equal to 0.344 and 2.671, respectively. The corresponding 90\,\% and 10\,\% cut-off wavelengths are obtained by inserting the corresponding values for a into Eq. (22).

Once the redshift is initiated in the cut-off region, the redshift heating causes relatively rapid temperature increase and density decrease. For example, in the transition zone to the solar corona, the temperature increases steeply. Below 50% cut-off for a given wavelength in the transition zone, the oscillator strength function $F_1(a)$ is less than 50% and above this cut-off it is more. By averaging, we can often for each wavelength set the oscillator strength function $F_1(a) = 1$ above the 50% cut-off and equal to zero below the 50% cut-off.

In the middle of the transition zone to the solar corona, we have that the electron density is about $N_e = 10^9$ cm^{-3}, the temperature is equal to 500,000 K and the value of $TN_e \approx 5 \cdot 10^{14}$ K cm^{-3}, see Vernazza, Averett, & Loeser (1981). From Eq. (23) we get for these values that the 50% cut-off wavelength is $\lambda_0 = 5000$ Å, which is about in the middle of the solar spectrum. Photons with wavelength shorter than 5000 Å will be redshifted more than 50% of the maximum redshift.

Above the cut-off limit the temperature increases and the density decreases until nearly the entire spectrum is redshifted. For a quiescent corona, detailed analysis shows that above the cut-off limit for the main portion of the spectrum, the redshift heating exceeds the X-ray and recombination cooling, causing the temperature to increase to about two million degrees (see sections 5.1 and 5.2 of (Brynjolfsson 2004a)). At and above this maximum temperature the gravitational cooling by the solar wind causes the temperature to fall. Below the maximum coronal temperature, the heat conduction moves the excess heating into the lower transition zone and helps compensate the cooling. At higher frequencies, the cut-off zone penetrates deeper into the transition zone. For example, for the same pressure and for $\lambda_{0.5} = 1215$ Å the cut-off is at about 200,000 K. Any excess heating from the transition zone leaks into the chromosphere. As shown in Brynjolfsson (2004b) the plasma redshift explains many other solar phenomena.

The temperatures and densities in the transition zone to the corona are experimentally determined and well documented in Vernazza, Averett, & Loeser (1981). So have many other phenomena explained in Brynjolfsson (2004a). Clearly, we did not make any adjustment when deriving Eqs. (14) and (24), which are based only on conventional physics. The onset of the plasma redshift in the transition zone with the concurrent heating, as estimated quantitatively in reference Brynjolfsson (2004a), confirm the plasma redshift.

1.7. The Plasma Redshift Term in Eq. (14)

With help of the oscillator strength function $F_1(a)$ in Table 1, where $a = \hbar\omega_p/(\beta_0\omega_0 kT)$, we can now determine the plasma redshift cross section. When we in Eq. (14) focus on the redshift term (and disregard first, second, and the fourth term within the brackets), we get that the plasma redshift is

$$-\frac{d\hbar\omega}{\hbar\omega} = \frac{4\pi}{3} r_e^2 N_e \frac{\gamma}{\gamma_0} F_1(a)\,dx = 3.326 \cdot 10^{-25} N_e \frac{\gamma}{\gamma_0} F_1(a)\,dx. \qquad (24)$$

For $a = \hbar\omega_p/(\beta_0\omega_0 kT) \approx 0$, the oscillator strength, $F_1(a)$ is about equal to 1, as seen from Table 1. But as a increases the $F_1(a)$ decreases.

When we then integrate each side of Eq. (24) and set $\lambda - \lambda_0 = \Delta\lambda$, we get that the plasma redshift $z = \Delta\lambda/\lambda_0$ is given by

$$-\int_{\omega_0}^{\omega} \frac{d\omega}{\omega} = \ln\left(\frac{\omega_0}{\omega}\right)$$

$$= \ln\left(1 + \frac{\Delta\lambda}{\lambda_0}\right)$$

$$= \ln\left(1 + z\right)$$

$$= 3.326 \cdot 10^{-25} \int_0^R F_1(a) \frac{\gamma}{\gamma_0} N_e dx. \tag{25}$$

The approximation $ln\left(1 + \Delta\lambda/\lambda\right) \approx \Delta\lambda/\lambda = z$ is valid for small redshifts z, such as those in the Sun and most stars, while for large cosmological redshifts and quasars' redshifts, we must use the logarithmic expression.

1.8. The Importance of the Photon Width γ

We see from Eq. (25) that the redshift is proportional to the photon width γ. But when a photon penetrates and interacts with the electron plasma, the photon width, γ, must approach the classical photon width, γ_0, which is equal to the quantum mechanical width of photons interacting with an electron plasma. We assume that a small incremental change in the width is proportional to the difference, $\gamma - \gamma_0$, and proportional to the plasma redshift. The incremental change in the photon width is then

$$d\gamma = -\xi \frac{(\gamma - \gamma_0)\,\omega}{\gamma_0} \frac{4\pi r_e^2}{3} N_e dx, \tag{26}$$

where $\gamma_0 = \beta_0 \omega^2 = 6.266 \cdot 10^{-24} \omega^2$ is the classical width as well as the quantum mechanical width of photons penetrating and interacting with the electron plasma. In Eq. (26), ξ is an adjustment factor. A rough estimate from observations of the resonance line of Na-I in the Sun indicates that $\xi \approx 0.25$. At this stage, we cannot be sure that ξ is a constant; but it applies generally to the great many lines investigated. When we integrate Eq. (26), we find that transition from γ_i to γ_0 takes place for small redshifts of $z \leq 5 \cdot 10^{-7}$ of the first term in Eq. (27); that is, it takes place in the transition zone to the corona. When we have a better theory for the forces within the photon, we may be able to determine ξ theoretically, but at this stage we must determine it experimentally. From Eq. (26) we determine γ as a function of x, and insert that value into Eq. (25). For oscillator strength function $F_1(a)$ equal to 1, Eq. (25) takes the form

$$\ln\left(1 + z\right) = 3.326 \cdot 10^{-25} \int_0^R N_e dx + \frac{\gamma_i - \gamma_0}{\xi\,\omega} = 3.326 \cdot 10^{-25} \int_0^R N_e dx + \frac{\delta\lambda_i - \delta\lambda_0}{\xi\,\lambda}, \tag{27}$$

where γ_i is the initial photon width and γ_0 is the final photon width, and where the electron integral is greater or equal to $1.5 \cdot 10^{18}$. The latter form of the equation applies when the photon width $\delta\lambda$ is in wavelength units. The second term on the right side of Eq. (27) is important for small redshift, such as the redshifts in the solar corona. It is also significant in collapsar, such as the white dwarfs, because of the large pressure broadenings (including Stark broadening). Interstellar space contains enough electrons to produce this redshift even if the star has a thin or no corona (a cold dwarf). In galactic coronas, intergalactic space, and in supernovas, the second term on the right side in Eq. (27) is small and can be disregarded, because the first term on the right side of Eq. (27) is large and the line broadenings rather small.

Eqs. (26) and (27) take into account that when emitted from the Sun, the photon widths of strong and weak lines vary greatly; but when measured on Earth, the photon widths should all be about equal to $\beta_0\omega^2 = 6.266 \cdot 10^{-24}\omega^2$, corresponding to $\delta\lambda \approx 0.118$ mÅ. The width, $\beta_0\omega_0^2$, of the solar photons arriving on Earth have not been measured so far, but they could be measured.

1.9. The Magnetic Field Affects the Cut-off Wavelength

The effect of the magnetic field, B, is significant and is described in section 4 of Brynjolfsson (2004a). When a magnetic field penetrates the plasma, the electrons will encircle the field lines. The effect of the magnetic field couples to the incident photon field, and the field of the Fourier components of the fast moving charged particles. This doubles the number of roots in the dielectric constant without affecting the total cross section. However, the radiation losses, which are caused by the acceleration of the electrons as they encircle the field lines, are significant. This increases the damping of the plasma electrons and must be added to the collision damping $\beta\omega^2$. The main result (see section 4 of Brynjolfsson (2004a)) is that Eq. (23) must be replaced by

$$\lambda_{0.5} = 3.185 \cdot 10^{-6} \left(1 + 1.3 \cdot 10^5 \frac{B^2}{N_e} \right) \frac{T}{\sqrt{N_e}}, \qquad (28)$$

At extremely high temperature, we must include a factor that takes into account relativistic effects. However, for the redshift in the solar corona and in most astrophysical plasmas, this factor, which is about $[1 + (1 - v^2/c^2)^{-1/2}]/2$, is not important. The effect of the magnetic field in Eq. (28) is to bring the plasma redshift cut-off zone deeper into the solar atmosphere. Sometimes, the cut-off zone reaches down into the chromosphere and for large fields even into the reversing layers of the Sun. An initiation of a plasma redshift in these layers results in formation of hot plasma "bubbles" followed by spicules, flares, prominences, and arches in the Sun, as described in Brynjolfsson (2004a). In these cases, the plasma redshift initiates and facilitates conversion of magnetic field energy to heat, as described in section 5.5 and Appendix B of Brynjolfsson (2004a). For the cosmological redshifts, the role of the magnetic field is usually insignificant. But in stars, and active galactic nuclei, including quasars, the magnetic field plays an important role.

2.   Comparison with Experiments

2.1.   The Temperatures and Densities in the Solar Corona

The plasma redshift predicts well the temperatures and densities in the solar corona; see sections 5.1 to 5.5 of Brynjolfsson (2004a). The many observations by many researchers have given a good determination of the temperature and the densities. However, it has never been possible to explain them theoretically. But in section 5.2 of Brynjolfsson (2004a), (see Table 2 and Figs. 1 and 2 of that source), we show the results of extensive computer calculations based on plasma redshift heating, X-ray cooling, ionization cooling, and the thermal conduction give a good description of the temperature and density distribution.

It has long been a mystery what determines the observed solar wind. In section 5.3 to 5.3.5 of Brynjolfsson (2004a), we show how the plasma redshift heating and the repulsion of the diamagnetic moments by the magnetic field are important for explaining the solar wind.

2.2.   Solar Redshift Experiments

The plasma redshift theory predicts well the observed solar redshifts, without the gravitational redshift; see sections 5.6 to 5.6.3 of Brynjolfsson (2004a) and Table 3 and Fig. 4 of that source; and see also Brynjolfsson (2004d), which gives a detailed theoretical explanation of this remarkable discovery.

The photons and all atomic and nuclear frequencies are gravitationally redshifted while in the Sun, but during their travel to the Earth the gravitational redshift is reversed. It is generally accepted that the gravitationally redshift of frequencies of atoms and nuclei in the Sun are reversed when the atoms and the nuclei move from the Sun to the Earth. But the frequencies of the photons are generally assumed to stay constant as the photons move from the Sun to the Earth. However, the comparison of the prediction of the plasma redshift theory with observations shows that, contrary to general beliefs, the frequencies of photons, like the frequencies of the atoms and the nuclei, reverse their gravitational redshift when they move from the Sun to the Earth.

Einstein (based on his classical physics thinking that equally many waves should arrive on Earth as left the Sun) believed that the photons retain their redshifted frequencies when they travel from the Sun to the Earth. Therefore, when the solar photons arrived on Earth, they would be observed gravitationally redshifted. It is generally believed that great many experiments have proven this assumption to be correct. However, an analysis of all these experiments, see Brynjolfsson (2004d), shows that they were incorrectly designed and incorrectly interpreted, because according to the uncertainty principle, the photons in these experiments had no chance to change their frequency during the short time between the emission and absorption of the photons. For example, in Pound et al. experiments the photons' travel time was $7.5 \cdot 10^{-8}$ seconds, while a minimum of $1.9 \cdot 10^{-5}$ seconds is required for the photons to change their frequency; see details in Brynjolfsson (2004d). In contrast, the solar photons have ample time, 8.3 minutes, to reverse their frequencies during their travel to Earth.

There is, however, a fundamental difference between the photons and the particles. In case of particles, we can say that we transfer energy to them as we

lift them from the Sun to the Earth. This transfer of energy to the particles is the cause of the reversal of their gravitational redshift. But in case of photons, they move to Earth without us lifting them. Their increase in frequency must therefore be due to repulsive force acting on the photons as seen by a distant observer (on Earth). A local observer (an observer at the position of the photon) sees the photons as weightless. Although the photons, in a local system of reference, therefore have no gravitational mass, $m_g = 0$, their inertial mass is as before $m_i = h\nu/c^2$. In a reference system of a distant observer, the photons are being repelled by the gravitational field. As we show later, this discovery leads to the conclusion that there are no black holes, and that the world can renew itself forever. This modification of the equivalence principle does not destroy the General Theory of Relativity. It only modifies it. Presently, the repulsion applies only to photons (see Brynjolfsson (2004d)).

2.3. Galactic Corona

Experiments have shown that our Galaxy has an extensive and dense corona; see section 5.7 of Brynjolfsson (2004a). The stars, the quasars, and the galaxies must, like the Sun, be surrounded by extensive coronas, which lead to intrinsic redshifts of all these objects. Lyman Spitzer Jr. & Ostriker (1997) was one of the first to realize that our Galaxy must have extensive corona. Not knowing the plasma redshift, nor the hot or relatively dense intergalactic plasma, he had difficulties explaining the required heating; see section 5.7 to 5.7.7 of Brynjolfsson (2004a). The experiments by Pettini et al. (1989), using the bright light from SN1987A in the Large Magellanic Cloud (LMC), made it clear that the Galactic corona contains hot plasma with electronic column density in excess of $1.6 \cdot 10^{21}$ cm^{-2}. This column density requires not only the plasma redshift heating from the galactic light, but also the X-ray heating and conduction heating from the hot intergalactic plasma, which is heated by the cosmological plasma redshift. This corona and the HII regions within the Galaxy lead to intrinsic redshifts of our Galaxy. Other galaxies and galaxy clusters must also have similar intrinsic redshifts. Some of the hot and dense coronal plasma around galaxies has energy in excess of the gravitational potential and must therefore diffuse into intergalactic space and fill it with plasma, which is heated by the cosmological redshift. The intrinsic redshifts of galaxies and of galaxy clusters increase the value of the Hubble constant, $H_0 = 3.076 \cdot 10^5 \cdot (N_e)_{avg}$, for nearby galaxies as mentioned in subsection 2.5 below.

2.4. Cosmological Redshift

The plasma redshift explains fully the cosmological redshift; see section 5.8 and 5.9 and Fig, 5 and 6 of Brynjolfsson (2004a); also Brynjolfsson (2004b) and Brynjolfsson (2006a) (see, especially, Fig. 2 of that source), which show that the plasma redshift explains well the observed magnitude-redshift relation of SNe Ia. There is no need for artificial parameters, such as Dark Energy or Dark Matter. Proper evaluation of the experimental data shows also that there is no Cosmic Time Dilation, and therefore no expansion. The expansion in the Big Bang requires Cosmic Time Dilation. Many authors contend to have experimentally proven the time dilation. In all cases that I have analyzed, I find this conclusion is based on impermissible disregard for the Malmquist bias.

In fact, when the Malmquist bias is taken properly into account, their data indicate to me that there is no time dilation. The average electron density and temperature in intergalactic space is about $N_e = 2 \cdot 10^{-4}$ cm^{-3}, and $T_e = 2.7 \cdot 10^6$ K; see Eqs. (57) and (62) of Brynjolfsson (2004a).

2.5.   Hubble Constant

The Hubble constant is given by $H_0 = 3.076 \cdot 10^5 \cdot (N_e)_{avg} \ km \, s^{-1} \, Mpc^{-1}$, where $(N_e)_{avg}$ is the average electron density in cm^{-3} along the line of sight; see Eq. (49) of Brynjolfsson (2004a). When we correct for the intrinsic redshifts of the galaxies, the Hubble constant is about $H_0 = 60 \ km \, s^{-1} \, Mpc^{-1}$; see Brynjolfsson (2004c). The uncertainty is similar to that reported by the supernova researchers, who use the Dark Energy and the Dark Matter parameters to reduce the variance. The intrinsic redshifts of galaxies increase the uncorrected average value of H_0. When uncorrected for intrinsic redshifts, the value of H_0 is, therefore, larger for nearby galaxies than for distant galaxies. Big Bang cosmologists, who deny intrinsic redshifts, conclude therefore incorrectly that the expansion of the Universe is accelerating.

2.6.   Distances

In plasma redshift cosmology, the distance, R, is given by $R = (c/H_0) \cdot \ln(1 + z)$**, or** $R = [0.9746/(N_e)_{avg}] \cdot \ln(1 + z)$ Mpc, where c is the light velocity in $km \, s^{-1}$, and $(N_e)_{avg}$ in cm^{-3} is the average electron density along the line of sight; see Eqs. (49) and (50) of Brynjolfsson (2004a). By differentiating, we see that dR/dz is proportional to $1/(1 + z)$, which decreases with z. Also this causes the Big Bang cosmologists to think that the expansion of the Universe is speeding up.

2.7.   Compton Scattering

In plasma redshift cosmology the redshift and the Compton scattering reduce the light intensity by a factor of $1/(1+z)^3$, while in Big Bang the redshift and the time dilation reduce the light intensity by a factor of $1/(1+z)^2$. The greater reduction in plasma redshift cosmology is partially compensated for by the difference in distance-redshift relation. This also explains why it is so important for the supernova researchers to retain the reduction caused by the cosmic time dilation. In plasma redshift cosmology, we explain the measured magnitude-redshift relation without any parameters and without any time dilation, while the Dark Energy and Dark Matter parameters in addition to Cosmic Time Dilation are essential for the Big Bang explanation.

2.8.   Surface Brightness

In plasma redshift cosmology, the surface brightness: $i_{SB} = (i_{SB})/[R_{pl}^2 \cdot (1+z)^3]$**, while in Big Bang cosmology it is:** $i_{SB} = (i_{SB})/[R_{bb}^2 \cdot (1+z)^4]$, where R_{pl} and R_{bb} are the distances in plasma-redshift and Big-Bang cosmology, respectively; see Eq. (A13) and Eq. (A14) of Brynjolfsson (2006b). We see thus that in plasma-redshift cosmology the surface brightness varies with z like the light intensity I, while in Big-Bang cosmology the surface brightness reduces much faster with z (a factor of $1/(1+z)^2$ faster) than the light intensity.

As shown in Brynjolfsson (2006b), the measurements by Sandage and Lubin [Sandage & Lubin (2001); Lubin & Sandage (2001a); Lubin & Sandage (2001b); Lubin & Sandage (2001b)] are in agreement with plasma-redshift cosmology while conflicting with Big-Bang cosmology. Lubin and Sandage explain this discrepancy with the Big Bang, as due to some kind of evolution. But I have failed to find any corroborating evidence for such an evolution (Brynjolfsson 2006b).

2.9. Cosmic Microwave Background

The cosmic microwave background (CMB) is emitted by the hot intergalactic plasma; see section 5.10 and Appendix C of Brynjolfsson (2004a). The energy density of the emitted microwaves is given by $a\,T_{CMB}^4 = 3NkT_e\,erg\,cm^{-3}$, where $a = 7.566 \cdot 10^{-15}\,dyne\,cm^{-2}\,\mathrm{K}^{-4}$ is the Stefan-Boltzmann constant; T_{CMB} = 2.728 K is the blackbody temperature of CMB; $N \approx N_p + N_{He} + N_e = (2.3/1.2)\,N_e \approx 1.917\,N_e\ cm^{-3}$ and $(N_e)_{avg} = 2 \cdot 10^{-4}\ cm^{-3}$; k=$1.38 \cdot 10^{-16}$ the Boltzmann constant; and $T_e = 2.7 \cdot 10^6$ K is the average particle temperature; see Eq. (61) of Brynjolfsson (2004a). The temperatures and the densities of the particles vary significantly, but their pressure $p = N\,k\,T_e$ (which determines the energy density) is nearly constant, when averaged over huge dimensions (about 5000 Mpc) of space. The Compton scattering on the electrons reduces the angular fluctuations. For the densities of intergalactic plasma and in the frequency range of the CMB, the plasma-redshift cross section is more than million times greater than the cross section in the free-free absorption and emission. Therefore, the blackbody temperature T_{CMB} is well defined and isotropic.

However, in the galaxy clusters, the plasma pressure should increase due to the higher gravitational attraction. Therefore, the T_{CMB} should increase especially in direction of nearby galaxy clusters. I believe this is the main cause for the observed dipole anisotropy in the CMB. This dipole indicates that the local group is moving with velocity of about 612 $km\,s^{-1}$ in direction of $l \approx 247$ and $b \approx 37$. This is roughly in the direction of Centaurus super cluster, the Great Attractor, and the Virgo cluster. This explanation is corroborated by Fig. 8 of Mould et al. (1991), which shows that redshift distances to several galaxies in the Centaurus supercluster exceed significantly the Tully-Fisher distances. In plasma redshift cosmology, these excessive redshift distances indicate that the average plasma densities to these galaxies are relatively high, which is consistent with higher CMB temperatures in these directions.

2.10. Cosmic X-ray Background

The cosmic X-ray background is emitted by the hot intergalactic plasma; see section 5.11, and sections C2 and C2.1 and C3 of the Appendix C in Brynjolfsson (2004a). After more than 30 years of intense studies, the origin of the X-ray background remained a mystery. However, the values of $(N_e)_{avg} = 2 \cdot 10^{-4}\ cm^{-3}$ and $(T_e)_{avg} = 2.7 \cdot 10^6$ K, which are needed for explaining the redshift-magnitude relation for SN Ia, and the blackbody spectrum of CMB, match exactly those that are needed for explaining the CXB; see discussion of Eq. (C21) in section C3 of Brynjolfsson (2004a).

The imbalance between the plasma redshift heating, which is a first order process, and the X-ray cooling, which is a second order process, causes significant temperature variation, with huge hot bubbles separated by colder regions at the

surfaces of the bubbles. These colder region are usually close to the galaxies, and when viewing bright objects through the colder regions, we usually see line absorptions. The centers of the hot bubbles are usually farther away from the galaxies, and when viewing bright objects through these hot regions, they are free of line absorption, except far into the X-ray region, due to the high temperatures. This is all consistent with observations.

2.11.   The High Average Densities in Intergalactic Space

In plasma redshift cosmology the average baryonic density is $\rho_{pl} = 3.806 \cdot 10^{-28} \cdot (H_0/60)\ g\,cm^{-3}$, and the Hubble constant is $H_0 \approx 60$. In Big-Bang cosmology, the baryonic density is $\rho_{BB} = 1.88 \cdot 10^{-29} \cdot \Omega_b\,h^2\ g\,cm^{-3}$. Often used values are $\Omega_b = 0.04$ and $h = (H_0/100) \approx 0.7$, which lead to $\rho_{BB} \approx 3.68 \cdot 10^{-31}$ $g\,cm^{-3}$. We see thus that $\rho_{pl} \approx 1000 \cdot \rho_{BB}$.

The Big Bang cosmologists often add Dark Matter and Dark Energy. In plasma-redshift cosmology, there is no need for Dark Matter nor Dark Energy; and the average density corresponds to electron density of $N_e = (1.95 \cdot 10^{-4}) \cdot (H_0/60)\ cm^{-3}$, and an average pressure of $p = 1.917k \cdot (1.95 \cdot 10^{-4}) \cdot (2.7 \cdot 10^6) = 1.4 \cdot 10^{-13}\ dyne\,cm^{-2}$, which is equal to the pressure of the CMB radiation.

In the Big-Bang cosmology, the need for Dark Energy and accelerated expansion stems mainly from the disregard for intrinsic redshifts. The need for Dark Matter stems mainly from the disregard for intrinsic redshifts, and from failure to realize that the coronas of galaxies are much denser than that assumed in the Big-Bang cosmology, because there were no means of heating the plasma.

In plasma redshift cosmology, 5 independent relations: 1) the magnitude-cosmological redshift relation, 2) the CMB relations for blackbody temperature as a function of density and particle temperature, 3) the intensity of X-ray background (XRB) as a function of density and particle temperature, 4) the redshift-distance relation for SN Ia, and 5) the relation between cosmological redshifts and surface-brightness, are all consistent with the an average density $N_e = (1.95 \cdot 10^{-4}) \cdot (H_0/60)\ cm^{-3}$, and an average electron temperature $T_e = 2.7 \cdot 10^6$ K in intergalactic space.

2.12.   Eternal Renewal of Matter

In plasma redshift cosmology old star matter is converted to primordial like matter in black hole candidates (BHCs) and supermassive black hole candidates (SMBHCs). When a large burned out star collapses to form BHC, the gravitational energy gained by a particle approaching the BH limit exceeds the particle's rest-mass energy.

When these hot particles collide, they will fission into protons, neutrons, quark-gluon plasma, and weightless photons. These weightless photons collect at the centers of the BHCs. The exchange forces acting on the fermions at the surface of the photon bubbles at the center exceed the gravitational attraction at the surface of the bubble. These exchange forces together with the pressure of the photon bubble counterbalance the gravitational attraction of the layers above the surface of the photon bubble. This tremendous pressure pushes the weightless photons towards the center of the BHC and SMBHC and prevent the particles from ever reaching the BH limit. This is all self-regulating. When matter falls on the BHC or SMBHC, more matter is transformed into photons. This large

reservoir of photons and primordial matter may then be released in form of gamma-ray bursts and primordial matter, which then lead to primordial-like nucleosynthesis. This explains also the starforming regions around the SMBHC at the center of our Galaxy.

3.   Conclusions

We have shown in sections 2.1 to 2.12 that the observations of great many cosmological phenomena verify the predictions of the plasma redshift theory, which is deduced in sections 1.1 to 1.9 from conventional laws of physics. We have failed to find any experiment that contradicts it.

The failure to deduce plasma redshift from the basic physical laws has been due to misleading approximations in the conventionally used equations. The conventional approximations were usually permissible for their intended usage in laboratory medium. But in the hot sparse plasmas of the coronas of stars, galaxies, and in intergalactic space, some of these approximations become impermissible.

Plasma redshift is given by Eq. (27) and Eq. (28). In the deduction, we did not change or surmise any new physical laws. We used only conventional, long-established, and basic axioms of physics. Plasma redshift cosmology should replace, therefore, the Big-Bang cosmology, which requires many physically strange assumptions, such as, the Big-Bang, Inflation, Dark Energy, Dark Matter, and Black Holes.

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Appendix A

A1   Electromagnetic Waves in Dielectrics

We consider a homogeneous and isotropic medium with a dielectric constant, ε, and a permeability, μ. Initially, these material constants do not vary with the coordinates nor with time. When using Gaussian (cgs) system of units, we get from Maxwell's electrodynamic theory that

$$\nabla \times \mathbf{E} = -\frac{1}{c} \cdot \frac{\partial \mathbf{B}}{\partial t} = -\frac{\mu}{c} \cdot \frac{\partial \mathbf{H}}{\partial t} \tag{29}$$

$$\nabla \cdot \mathbf{D} = \nabla \cdot \varepsilon \cdot \mathbf{E} = 0 \tag{30}$$

$$\nabla \times \mathbf{H} = \frac{1}{c} \cdot \frac{\partial \mathbf{D}}{\partial t} = \frac{\varepsilon}{c} \cdot \frac{\partial \mathbf{E}}{\partial t} \tag{31}$$

$$\nabla \cdot \mathbf{B} = \nabla \cdot \mu \cdot \mathbf{H} = 0 \tag{32}$$

where the dielectric constant, ε, and the permeability, μ, depend only on the frequency and where c is the velocity of light.

Comment A1. We can obtain the corresponding equations in the mks (rationalized) system of units by replacing ε with $\varepsilon \cdot (\varepsilon_0 \cdot c)$, and μ with $\mu \cdot (\mu_0 \cdot c)$, where ε_0 and μ_0 are the dielectric constant and permeability in vacuum, and where $\varepsilon_0 \cdot \mu_0 = 1/c^2$.

We assume that the field varies sinusoidally as the real part of $\exp(i\,\omega\,t)$. For facilitating the calculations, we use complex notations for the different quantities; and in the usual manner, we use their modulus for comparison with experiments.

We write the general solutions to these equations on the form:

$$E_y = \frac{A(\omega)}{\varepsilon \cdot \sqrt{\mu}} \cdot \exp[i\,\omega(t - x\,\sqrt{\varepsilon\,\mu}/c)], \tag{33}$$

$$E_x = E_z = 0 \tag{34}$$

$$H_z = \frac{A(\omega)}{\mu \cdot \sqrt{\varepsilon}} \cdot \exp[i\,\omega(t - x\,\sqrt{\varepsilon\,\mu}/c)], \tag{35}$$

$$H_x = H_y = 0 \tag{36}$$

Comment A2. It is possible to set $B(\omega) = A(\omega)/\sqrt{\varepsilon\,\mu}$. The coefficients in front of the exponential factors in Eqs.(33) and (35) would then be $B(\omega)/\sqrt{\varepsilon}$ and $B(\omega)/\sqrt{\mu}$, respectively. Use of these coefficients leads to the conventionally used solutions, and the misleading assumption about the variations of the fields with ε and μ. However, the forms of the coefficients in Eqs. (33) and (35) are both mathematically correct and physically simpler to interpret, because $A(\omega)$ is independent of ε and μ. For example, for $\mu = 1$, we have that $\nabla \cdot \mathbf{D} = 4\pi\rho$, where $\varepsilon \cdot \mathbf{E} = \mathbf{D}$. Therefore, when we put a dielectric material around a charge, the vector $\mathbf{D}$ is unchanged; that is, the quantity $A(\omega)$ in Eq. (33) is a constant when ε changes.