Holographic Universe in Optoustics

J. Zhang[1]

The messages from universe, which we can observe or detect on the earth, are probably the aurora, the solar wind, the flare, and the storm.

Auroras are produced when the magnetosphere is sufficiently disturbed by the solar wind that the trajectories of charged particles in both solar wind and magnetospheric plasma, mainly in the form of electrons and protons, precipitate them into the upper atmosphere.

The solar wind is a stream of charged particles released from the upper atmosphere of the Sun. This plasmaconsists of mostly electrons and protons with energies usually between 1.5 and 10 keV; embedded in the solar-wind plasma is the interplanetary magnetic field.

Observers had seen large flares that preceded severe geomagnetic storms, shocks in the solar wind, and blasts of energetic particles in interplanetary space.

Therefore, the notion that flares cause these effects took firm root in the minds of scientists and remained the conventional wisdom for many years.

I would like to take a journey, starting from electrons and protons, navigating through magnetic field, detecting evidence for dark matter and dark energy, and finally observing the holographic universe in optoustics[2].

Electrons and Protons

The Photoelectric Effect[3]

Light is made of packets of energy called photons. Photons have no mass, but they have momentum and they have an energy given by:

$$E = hf$$

The explanation for the photoelectric effect goes like this: it takes certain energy to eject an electron from a metal surface. This energy is known as the work function, which depends on the metal. Electrons can gain energy by interacting with photons. If a photon has energy at least as big as the work

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[1] J.Zhang (Clifford), Deputy Dean, Institute of DC Science.
[2] I create a new word of optoustics by adding optics and acoustics together, meaning effect of optics and acoustics.
[3] Einstein won the Nobel Prize for Physics not for his work on relativity, but for explaining the photoelectric effect.

function, the photon energy can be transferred to the electron and the electron will have enough energy to escape from the metal.

The experimental facts connected with the photoelectric effect therefore lead to the inescapable conclusion that almost monochromatic photons cannot be split into two photons of the same frequency which carry only a fraction of the energy of the original photon, which further supported by the experimental results concerning the Compton effect.

The Compton Effect

Although photons have no mass, they do have momentum, given by:

$$P = E/c = hf/c = h/\lambda$$

Convincing evidence for the fact that photons have momentum can be seen when a photon collides with a stationary electron. Some of the energy and momentum is transferred to the electron[4], but both energy and momentum are conserved in such a collision. Applying the principles of conservation of energy and momentum to this collision, one can show that the wavelength of the outgoing photon is related to the wavelength of the incident photon by the equation:

$$\lambda' - \lambda = (h/mc)\,[1 - \cos\theta]$$

where θ is the angle between the incident and outgoing photons, and m is the mass of the electron.

Consider the rest-frame of the electron before any possible emission: in this frame the total energy is mc^2. The situation is different, however, when the electron passes through the strong electric field of a nucleus in the target. The conservation equation:

$$P_i = P_f + P_n + p$$

$$E_i + Mc^2 = E_f + E_n + \hbar\omega$$

Where M the mass of nucleus, P_i initial momentum, P_f collision momentum, P_n nucleus momentum, p photon momentum, frequency $\omega = pc/\hbar$.

Let us rewrite the conservation equations for the case when the final electron and nuclear velocities are indeed equal:

[4] This is known as the Compton effect.

$$\mathrm{p_i - p} = \frac{(M+m)v}{\sqrt{1-(v/c)^2}}$$

$$E_i + Mc^2 - cp = \frac{(M+m)v}{\sqrt{1-(v/c)^2}}$$

Multiplying the first equation by c, and subtracting the square of the result from the square of the second equation gives us

$$\hbar\omega = pc = \frac{E_i - mc^2}{1-(E_i - p_i c\cos\theta)/(Mc^2)}$$

where θ is the angle between the emerging photon and the incident electron.

The Photon-Electron System

It is convenient to classify systems according to the number and type of particles in the initial state. Thus, the photon-electron system is defined to be a system with exactly one photon and one electron in its initial state.

The application of the theory to special system is considerably simplified by the symmetry properties of the *S*-matrix under charge conjugation. The charge conjugation operator *γ* commutes with the *S*-operator, and a state vector and its charge conjugate are related by

$$\omega^c = \gamma\omega$$

It follows that the *S*-matrix element for a process, S_{fi}, is related to the S-matrix element for the charge-conjugate process, S_{fi^c}, by

$$S_{fi^c} = (\gamma\omega_f, S\gamma\omega_i) = (\omega_f, \gamma^{-1}S\gamma\omega_i) = (\omega_f, S\omega_i) = S_{fi}$$

The matrix elements of charge conjugate processes are equal.

The photon-electron system can have two kinds of final states, those in which there is only one electron present, and those in which there are also one or more pairs present. Processes leading to the former kind of final states may be called photon-electron scattering, whereas processes in which pairs are produces can be referred to as pair production in photon-electron collision, or pair production by photons in the field of an electron.

The scattering of a photon by an electron is called Compton scattering[5]. We call the photon-electron scattering process with two photons in the final state the double Compton scattering. The matrix element for this process is higher by one order in *e* than the matrix element for single Compton scattering. An electron of momentum *p*, and a photon of momentum *k* and polarization *e*, scatter into an electron of momentum *p'*, and a photon of momentum *k'* and polarization *e'*.

$$\sigma = (2\pi)^2 \frac{\epsilon\omega}{|p\cdot k|} S_f S_i \delta(p' + k' - p - k)|(f|M|i)|^2$$

where $\epsilon = p^0$, $\omega = k^0$.

Double Compton scattering is the lowest order approximation to photon-electron scattering in which the final state consists of an electron and two photons.

The summation over the spin directions of the initial and final states can be performed in exactly the same way as in the Compton effect. The final result is considerably simplified if we assume that the incident and outgoing photons are unpolarized, so that we shall also sum over the polarizations.

$$\sigma = \frac{\alpha {r_0}^2}{(4\pi)^2} \int \frac{X}{|p\cdot k|\epsilon'\omega_1\omega_2} \delta(p' + k_1 + k_2 - p - k) d^3k_1 d^3k_2 d^3p'$$

where $\alpha = e^2/4\pi \approx 1/137$ is the fine-structure constant.

The External Field

A classical Maxwell field and current density distribution are described by two four-vector field $\varphi_\mu(x)$ and $s_\mu(x)$. These are c-numbers and are related by the equations

$$\partial^\lambda \partial_\lambda \varphi_\mu(x) = -s_\mu$$

The current density satisfies the continuity equation

$$\partial^\mu s_\mu(x) = 0$$

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[5] A. H. Compton, *Phys. Rev.* 21,715 (1923).

which is consistent with the Lorentz condition for the potentials.

$$\partial^{\mu}\varphi_{\mu}(x) = 0$$

These equations can be obtained from a quantized theory by identifying the expectation values of the field and current operation with classical quantities.

Many photon-electron processes of radiation theory are enhanced considerably when the incident photon is replaced by an external field. Of particular interest is the nuclear Coulomb field, since it is extremely strong.

For example, Compton scattering is to be compared with Coulomb scattering, double Compton scattering with bremsstrahlung [6], pair production in photon-electron collisions with pair production by an electron in a Coulomb field.

We are dealing with a double expansion, one in α which describes the radiation field of the electron, and one in αZ which corresponds to a power series expansion in the external field.

We shall refer to Coulomb scattering as the scattering of an electron in a Coulomb field to all orders in this field, but to zero order in the radiation field. The effect of the presence of photons emitted and reabsorbed by the scattered electron, briefly called radiative correction.

It appears advantageous to evaluate the Coulomb scattering cross section to lowest order in αZ first in a purely formal way. The cross section for a static field is given

$$\sigma = (2\pi)^2 \frac{1}{\beta} S_f \hat{S}_i \delta(\epsilon_f - \epsilon_i) |f|M|) i|^2$$

where β is the incident velocity.

Electron-photon Coincidence

From the experiments of electron-photon coincidence, we can find that there is a strong angular correlation with the photons being predominantly emitted on the same side relative to the primary beam as the decelerated electrons. This behavior can be understood classically by considering the radiation of the electrons along the hyperbolic orbit around the nucleus.

[6] The process of radiative transitions between bound state levels has a counterpart in the continuous spectrum, namely, the change in momentum of an electron scattering in an external field which the emission of a photon.

The measurement of an angular distribution of decelerated outgoing electron for a fixed photon emission direction was reported by Hub and Nakel[7]. The electron distribution is strongly peaked into the forward direction. This, photons with an energy that is far from the high-energy limit are radiated preferably by electrons that are deflected through relatively small angles in the nuclear Coulomb field.

On the other hand, at the high-energy limit where all the kinetic energy of the incident electron is radiated, the electron angular distributions are more isotropic.

Bremsstrahlung produced by a beam of unpolarized electrons in the field of an atomic nucleus is, in general, partially linearly polarized depending on the photon energy, the photon emission angle, the initial electron energy, and the atomic number of the target.

The theoretical curve was calculated by means of the Elwert-haug theory[8]. It includes corrections for the finite solid angles and energy width of the detectors, the plural scattering of the electrons in the target, and a contribution of electron-electron bremsstrahlung.

The radiation is almost completely polarized with the electric vector in the reaction plane, containing the momenta of the incoming and outgoing electron and the photon.

Magnetic Field

An Electromagnetic Wave

Maxwell's description of the electromagnetic field was essentially complete. We have arrived by different routes at various pieces of it, which we shall now assemble in the form traditionally called Maxwell's equations:

$$\text{curl } E = - \frac{\partial B}{\partial t}$$

$$\text{curl } B = - \epsilon_0 \mu_0 \frac{\partial E}{\partial t} + \mu_0 J$$

$$\text{div } E = \frac{\rho}{\epsilon_0}$$

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7 Hub, R. and Pankau, E., Phys. Lett. 44A (1967) 601
8 Elwert, G. and Haug, E., Phys. Rev. 93 (1969) 90

$$\text{div } B = 0$$

These are written for the fields in vacuum, in the presence of electric charge of density ρ and electric current, that is, charge in motion of density J.

Notice that the lack of symmetry in these equations, with respect to B and E, is entirely due to the presence of electric charge and electric conduction current. In empty space, the term with ρ and J are zero, and Maxwell's equations become

$$\text{curl } E = -\frac{\partial B}{\partial t} \qquad \text{div } E = 0$$

$$\text{curl } B = \epsilon_0 \mu_0 \frac{\partial E}{\partial t} \qquad \text{div } B = 0$$

We are going to construct a rather simple electromagnetic field that will satisfy Maxell's equation for empty space. Suppose there is an electric field E, everywhere parallel to the z axis, whose intensity depends only on the space coordinate y and the time t. Let the dependence have this particular form:

$$E = \hat{z} E_0 \sin(y - vt)$$

in which E_0 and v are simply constants. This field fills all space. We will need a magnetic field, too. We shall assume it has an x component only, with a dependence on y and t similar to that of E_z:

$$B = \hat{x} B_0 \sin(y - vt)$$

where B_0 is another constant.

We will show now that this electromagnetic field satisfies Maxwell's equations if certain conditions are met. It is easy to see that div E and div B are both zero for this field. The other derivatives involved are

$$\text{curl } E = \hat{x} \frac{\partial E_z}{\partial y} = \hat{x} E_0 \cos(y - vt)$$

$$\frac{\partial E}{\partial t} = -v\hat{z} E_0 \cos(y - vt)$$

$$\text{curl } B = -\hat{z} \frac{\partial B_x}{\partial y} = -\hat{z} B_0 \cos(y - vt)$$

$$\frac{\partial B}{\partial t} = -v\hat{x}B_0\cos(y - vt)$$

The Magnetic Moment of the Electron

It is well known that the electron which satisfies the wave equation of Dirac has a magnetic dipole moment of magnitude $\mu_0 = e/2m$. In our field theoretical formulation of quantum electrodynamics this property of the electron is implicitly contained in the interaction operator

$$H(\tau) = ie\int_{\sigma(\tau)} \bar{\psi}\gamma_\mu\psi\,\varphi^\mu d\sigma$$

For the matter field ψ which interacts with an external electromagnetic field φ^μ.

Interactions with Magnetic Fields

Many important results of the theory of synchrotron radiation are obtained within the framework of classical electrodynamics. The classical treatment of the synchrotron radiation is limited by the condition

$$\frac{E_e}{m_e c^2}\frac{B}{B_{cr}} << 1$$

where $B_{cr} = \frac{m_e^2 c^3}{e\hbar} \approx 4.4 \times 10^{13}$ G is the so-called critical value of the magnetic field relevant to quantum effects.

General, the energy of synchrotron photons are much less than the energy of parent electrons. For interaction of electrons and photons with magnetic field, it is convenient to introduce interaction probabilities instead of standard total cross sections. But in the literature this parameter is still formally called a cross-section. These probabilities, normalized to the strength of the magnetic field.

Relativistic Electron-Photon Cascades

Relativistic electrons-directly accelerated, or secondary products of various hadronic process – may result in copious γ-ray production caused by interactions with ambient targets in forms of gas, radiation and magnetic field. In different astrophysical environments γ-ray production may proceed with high efficiency through bremsstrahlung, inverse Compton scattering and synchrotron.

Generally, γ-ray production in a given process is effective when the relevant radiative cooling time does not significantly exceed the source age, the time of non-radiative losses caused by adiabatic expansion or by particle escape, and the cooling time of competing radiation mechanisms resulting in low-energy photons outside the γ-ray domain.

As long as the charged particles are effectively confined to the γ-ray production region, in some circumstances these conditions could be fulfilled even in environments with relatively low gas and photon densities or a weak magnetic field.

Magnons in Ferromagnets

Spin waves represent a set of oscillators with frequencies ω_k in *k* – space that describe transverse deviations of magnetization from the magnetically ordered vacuum state. By convention, the quantum of spin wave is a quasi particle, which is called a magnon.

A magnon's energy is equal to $\hbar\omega_k$. According to quantum mechanics, the magnetic moment of the system is equal to the negative partial derivative of the energy in the magnetic field. This, a magnon carries the magnetic moment

$$\mu_k = - \frac{\partial \hbar\omega_k}{\partial H}$$

where H is the external magnetic field. Perturbation of the magnetic system can be regarded as an evolution of a magnon gas and this approach helps to explain many thermodynamic, dynamic, and kinetic phenomena in magnetics.

Magnetoelastic Waves

The elementary excitations of the ordered magnetic system are magnons, and those of the elastic solid are phonons. Magnetoelastic interactions give rise to a coupling of magnetic and elastic oscillations.

The resulting normal modes contain both the magnetic and elastic components. This means that such quasi particles can be excited, in principle, either by an alternating magnetic field or by elastic vibrations.

Here, we consider quasi phonons in the easy plane antiferromagnet, where the Magnetoelastic phenomena exhibit so-called exchange enhancement. The Hamiltonian of the two-sublattie antiferromagnet can be represented as

$$H = H_{afm} + H_e + H_{me}$$

where

$$H_e + H_{me} = \int (\mu_e + \mu_{me}) dr$$

are the elastic and magnetoelastic energies of the sample.

Indirect Elastic-Elastic Interactions

In the region of acoustic frequencies which is far below the frequencies of quasi ferromagnetic magnons, the magnetoelastic interactions can be simplified by introducing effective elastic-elastic interaction.

We shall assume conditions of quasi equilibrium of the form

$$\frac{\partial (H^{(2)} + H_{me})}{\partial \psi_k} = 0$$

where

$$\psi_k \equiv \hat{C}_k^-$$

is the generalized coordinate.

Relaxation of Bose Quasi Particles

Consider now the thermal bath Hamiltonian as a set of harmonic oscillators

$$H_b = \hbar \sum_q \Omega_q d_q^\dagger d_q$$

With the spectrum Ω_q. These oscillators could be magnons, phonons and other elementary Bose excitations.

Now we shall consider the magnon interaction with conduction electrons. It is convenient to describe small oscillations of the magnetization in terms of creation $\alpha^\dagger$ and annihilation α Bose operators introduced by a linearized Holstein-Primakoff transformation. For the transverse magnetization components one can write

$$M_x \approx -\sqrt{\frac{\hbar\gamma M_s}{2V}}(\alpha^\dagger + \alpha)$$

$$M_\gamma \approx -\frac{1}{i}\sqrt{\frac{\hbar\gamma M_s}{2V}}(\alpha^\dagger - \alpha)$$

The interaction of the uniform magnetic precession with electrons is assumed to be of the form:

$$\mathrm{V} = \frac{D^\dagger\alpha + D\alpha^\dagger}{\sqrt{N}}$$

where

$$\mathrm{D} = \frac{1}{N}\sum_{q.q'} \sum_j f_{q.q'}(r_j) d^\dagger_{q'} d_q$$

The amplitude

$$f_{q.q'}(r_j) = |f_{q.q'}| \exp[i\emptyset(r_j)]$$

Describes the scattering process in the vicinity of crystal defect at the point r_j;

$\emptyset(r_j)$ is the phase. The term $d^\dagger_q$ and d_q are the fermion creation and annihilation operators.

Relaxation of Coupled Oscillations

We analyze the relation in the system of coupled oscillation. We have already considered coupled oscillations in magnetic systems: nuclear magnons and quasi phonons which appear as the result of interaction of quasi ferromagnons with nuclear spin deviations and elastic vibrations of the lattice.

The frequencies of normal modes of the coupled oscillations can differ noticeably from the frequencies of original pure modes. The change in spectra also leads to the change of relaxation.

The shape of the resonance line $f_1(\Omega)$ is the normalized weight function, according to which the statistical average parameters of the oscillation, the moments M_n are determined:

$$M_n = \int_{-\infty}^{\infty} (\Omega - \omega_1)^n f_1(\Omega) d\Omega$$

Where ω_1 is the eigenfrequency of the oscillator and n = 1,2,…For example, for a harmonic oscillator with no damping

$$F(\Omega) = \sigma(\Omega\text{-}\omega),\ M_n = 0$$

If an interaction is turned on between two oscillators with known eigenfrequencies ω_1, and ω_2, and line shapes $f_1(\Omega)$ and $f_2(\Omega)$, the line shape of any new normal mode of this coupled system can be represented in the form

$$\tilde{f}(\Omega) = \iint \delta[\Omega - \widetilde{\omega}(\Omega_1, \Omega_2)]\, f_1(\Omega_1) f_2(\Omega_2) \mathrm{d}\Omega_1 \mathrm{d}\Omega_2$$

Where $\widetilde{\omega}(\omega_1, \omega_2)$ is the frequency of the corresponding normal mode as obtained from the characteristic equation. We can easily obtain the formula for the moments of the new normal mode:

$$\widetilde{M_n} = \iint \left[\widetilde{\omega}(\Omega_1, \Omega_2) - \widetilde{\omega}(\omega_1 \omega_2)\right]^n f_1(\Omega_1) f_2(\Omega_2) d\Omega_1 d\Omega_2$$

Expressions give a general solution of the problem of finding the linewidths of the normal modes of two coupled oscillation.

Microwave Pumping of Magnons

Magnons can be excited by the external microwave magnetic field applied to the magneto-ordered system. Depending on the geometry, frequency and power, the growth of magnon population can be caused by linear or nonlinear processes.

The simplest example of a linear process is the ferromagnetic resonance when the uniform magnetization precession is excited by the transverse alternating magnetic field with frequency ω_0. Nonlinear processes open the possibility of exciting magnons with $k \neq 0$ via parametric resonance conditions.

One of the most frequently used and powerful techniques is the method of so-called parallel pumping when the microwave field is parallel to the DC magnetic field. In this case the decay parametric process occurs: a microwave photon creates a pair of magnons with the oppositely directed wave vectors and the energies $\hbar\omega_k = \hbar\omega_{-k} = \hbar\omega_p/2$. The photon wave vector in this case is negligibly small and can be considered as zero. This process has a pumping field threshold above which the number of magnon pairs grows exponentially.

The emission of electromagnetic waves are parametrically excited by the quasi phonons. The emission can be observed after the end of the microwaves pump pulse. In these cases the microwaves signal passed through the cavity can be detected by the crystal detector and displayed on an oscilloscope.

As the frequencies ω_P and ω_R converge, the increasingly stronger repulsion of the branches of the mixed photon-quasi phonon modes can be observed, we conclude that the new normal modes appear due to nonlinear interaction between the microwave resonator and the sample.

It is well know that the collective oscillations of parametric pairs represent the oscillation of their number and phase in the vicinity of equilibrium values. However, the efficiency of nonlinear photon-quasi phonon interaction depends on the number of parametric pairs.

For this reason both the efficiency of photon-quasi phonon coupling and value of splitting oscillate with the same frequency and so the side band amplitudes contain the information not only about parametrically excited magnetoelastic waves but also on the efficiency of nonlinear photon-quasi phonon interaction.

A direct confirmation of existence of nonlinear coupled photon-quasi phonon modes can be obtained by the observation of nonmonotonic electromagnetic emission from the resonator and sample system after the switching off of the pumping field[9]. The electromagnetic emission after the end of the microwave pump pulse gives the time evolution of this nonequilibrium system in the absence of the external driving force.

Bose-Einstein Condensation of Quasi Equilibrium Magnons

Let us consider a general form of the boson populations:

$$n_k = [\exp(\frac{\hbar\omega_k - \mu}{k_B T})]^{-1}$$

where the chemical potential μ is always less than the energy $\hbar\omega_k$ in order to have positive and finite n_k. The chemical potential is defined as a solution of the integral equation

$$N = \sum_k n_k = \sum_k [\exp\left(\frac{\hbar\omega_k - \mu}{k_B T}\right) - 1]^{-1}$$

where N = const. is the total number of bosons. From this formula follows the condition of Bose-Einstein condensation μ = 0 for real particles with mass m and energy $\hbar\omega_k = (\hbar k)^2/2m$

$$\frac{N}{V_s} = \int_0^\infty \frac{4\pi k^2\,dk}{(2\pi)^3}[\exp\left(\frac{(\hbar k)^2}{2mk_B T} - \frac{\mu}{k_B T}\right) - 1]^{-1}$$

$$= \frac{(2mk_B T)^{3/2}}{8\hbar^3 \pi^{3/2}} Li_{3/2}[\exp\,(\frac{\mu}{k_B T})]$$

Here V_s is the volume of the system and

$$Li_v(\mathrm{z}) \equiv \sum_{n=1}^{\infty} \frac{z^n}{n^v}$$

is the PolyLog function.

There is a principal difference between magnons as quasi particles and real

[9] Andrienko, A.V. and Safonov, V.L.(1995) observation of coupled photon-phonon oscillation with parametric excitation of magnetoelastic waves in an antiferromagnet.

Bose particles. The number of real particles is not changed when the temperature decrease. But the number of magnons decrease with decreasing temperature. There are no spin excitation. This is why the chemical potential of magnons in the thermodynamic equilibrium is always equal to zero and the total magnon number is defined by

$$N = \sum_k [\exp\left(\frac{\hbar\omega_k}{k_B T}\right) - 1]^{-1}$$

This formula obviously has no conditions for the Bose-Einstein condensation of magnons.

Bose-Einstein[10] can occur in quasi equilibrium thermodynamic conditions which result from strong parametric pumping of magnons. Magnons, as conventional elementary excitations of magnetically ordered electronic spins, have relatively long lifetimes.

The principal nonlinearities in the system conserve the number of magnons and do not change the Bose distribution function. Magnons are still good bosons at high densities.

In particular, it is shown that a magnetic system under strong microwave pumping there are conditions when the gas of elementary magnetic excitation can be described thermodynamically with some effective temperature and chemical potential. Bose-Einstein condensation of quasi equilibrium magnons occurs when the chemical potential approaches the bottom of the magnon spectrum.

Let us consider magnons in a magnetic sample as an isolated ideal Bose gas which is in a thermodynamic quasi equilibrium with population n_k. The relation

$$N(\mu,\ T_{eff}) = \sum_k [\exp\left(\frac{\hbar\omega_k - \mu}{k_B T_{eff}}\right) - 1]^{-1}$$

binds together the total number of quasi equilibrium magnons *N*, the chemical potential μ and the effective temperature T_{eff}.

Taking into account the condition of Bose-Einstein condensation $\mu = \mu_c$, we obtain the critical density of magnons

[10] The macroscopic quantum phenomenon in magnon gas

$$\frac{N_c}{V_s} = \frac{N(\mu_c, T_{eff})}{V_s}$$

This critical number contains both the initial density of thermal magnons $N(0,T)/V_s$ and the density N_p/V_s of magnons pumped, for example, by the external microwave field. Thus, the critical density of pumped magnons is

$$\frac{N_{p,c}}{V_s} = \frac{N(\mu_c, T_{eff})}{V_s} - \frac{N(0,T)}{V_s}$$

Let us consider the coherent emission of photons by the condensate with k = 0. The magnetodipole emission intensity is given by the formula:

$$I = \frac{2}{3c^3}\left(\frac{d^2 MV_s}{dt^2}\right)^2$$

where MV_s is the magnetic moment of the sample and c is the speed of light.

Ionized Gas

The electrons and ions in a thermal plasma scatter off one another. In classical electromagnetism, an accelerating charge radiates electromagnetic energy; therefore, we expect these scattering events to produce electromagnetic radiation. The electrons, with their much smaller mass, undergo much larger accelerations and therefore dominate the radiation.

The emission is a continuum extending from very low frequencies, up to frequencies where the emitted photon energies are comparable to the thermal energy *kT*.

A classical analysis of the power radiated by electrons scattered by ions with charge Z_ie leads one to write the free-free emissivity[11] as

$$j_{ff,v} = \frac{8}{3}\left(\frac{2\pi}{3}\right)^{1/2} g_{ff,i}\frac{e^6}{m_e^2 c^3}\left(\frac{m_e}{kT}\right)^{1/2} e^{-hv/kT} n_e Z_i^2 n_i$$

[11] Power radiated per unit frequency, per unit volume, per steradian.

where the dimensionless factor $g_{ff,i}(v,T)$ is called the Gaunt factor for free-free transitions.

The nonrelativistic quantum mechanics of hydrogen and one-electron ions is simple enough that the ground state photoelectric cross section for photons with energy $\hbar v > Z^2 I_H$ is given by an analytic expression:

$$\sigma_{pe}(v) = \sigma_0 (\frac{Z^2 I_H}{hv})^4 \frac{e^{4-4arctan(x)/x}}{1-e^{-2\pi/x}}$$

where Z is the atomic number of the nucleus, and σ_0, the cross section at threshold, is given by

$$\sigma_0 \equiv \frac{2^9 \pi}{3e^4} Z^{-2} \alpha \pi a_0^2$$

The high pressure ionized gas is eventually able to break out of the blister. Once it has broken out of the confining molecular gas, the ionized gas is able to exhaust into the lower density interstellar medium, in what is termed a "champagne flow"[12]. The actual geometry is best visualized by maps of the free-free continuum emission at radio frequencies.

Collisions of Galaxies with an Intracluster Gas

Galaxies are comparatively fragile objects, having lower escape velocities than stars and being composed of regions of dramatically different densities. The gas is loosely bound to the disk or even halo and can be removed if the Galaxy moves through a background gas within a cluster.

The typical density of the interstellar medium (ISM) within a Galaxy has a temperature of about 100K, so that the intracluster medium (ICM) progagates through the Galaxy provided $P_0 < n_{ICM} V_g^2$. This forms a wall that moves through the ISZM of the Galaxy in the frame of the Galaxy.

To compute the velocity of this snowplowed material, we have to assume that the piston moves at a velocity V_g. Then as it moves into the gas at rest, we see that

[12] Tenorio-Tagle, G., 1979, "The gas dynamics of HII regions. I-The champagne model." *Astr.As.trophys.*, 71

$$v_g^2 = 2(p_2 - p_1)(\frac{1}{\rho_1} - \frac{1}{\rho_2})$$

which comes from the shock adiabat[13].

Stellar Winds

Stellar winds inject mass, momentum, and energy into the ISM, and can carve out conspicuous structures in the ISM in the vicinity of the star. The most important stellar winds come from hot, young, massive stars; from red giants and supergiants, and from the progenitors of planetary nebulae.

Stellar winds normally have terminal speeds V_ω[14] that are of order the escape velocity from the stellar surface. The ionizing radiation from the star will have already created an HII region into which the wind will blow. The bubble is spherical.

Because greatly exceeds the sound speed in the HII, the wind will initially drive a shock into the HII region – the outer shock – with radius R_s and velocity V_s.

At very early times, the wind undergoes essentially free expansion, with until the wind mass and swept-up mass are comparable: $Mt \approx (4\pi/3\, \rho_0 (V_\omega t)^3)$. The free expansion phase ends at a time

$$t_0 \approx (\frac{3\dot{M}}{4\pi\rho_0 V_\omega^3})^{1/2} = 2.54 n_3^{-1/2} M_{-6}^{1/2} V_{\omega 8}^{-3/2} yr$$

The stellar wind is the energy source, hence the total energy $E(t) = (1/2)\dot{M}V_\omega^2 t$.

Now what if the momentum input is from a continuous source? For example, if there is a stellar wind inside a bubble, which is pushing on the medium and causing it to expand, what is the solution? It is very easy to show that it is $R \sim t^{1/2}$. The dimensionless approximation is that, for continuous mass loss, $MV = \dot{M}V_t$.

We can specify the source as we wish. It may be a stellar wind driven by

[13] Assume that momentum and mass are conserved across the shock interface.
[14] The asymptotic speed of the wind after it has traveled many stellar radii for the star.

radiation pressure. It may be due to wind which is driven by turbulence as in the Sun. Whatever the mechanism for generating the outflow, it can be incorporated. The same is true for the case of an adiabatic outflow.

Dark Matter and Dark Energy

Radiative Cooling

In a cold plasma with electron density n_e, electromagnetic waves propagating with $\mathrm{E} \propto e^{ikx-i\omega t}$ must satisfy the dispersion relation[15]

$$k^2c^2 = \omega^2 - \omega_p^2$$

where

$$\omega_p \equiv (\frac{4\pi n_e e^2}{m_e})^{1/2}$$

is known as the plasma frequency. From this equation, it is evident that there are no propagating mode with frequencies below the plasma frequency ω_p. Radio astronomical observation are in general conducted at frequencies $v \geq 10^8$Hz, far above the plasma frequency $v_p = \omega_p/2\pi = 8.979 \times 10^3 (n_e/cm^{-3})^{1/2}$ Hz, so we can in general make the approximation $\omega_p/\omega \ll 1$ when discussing propagation of electromagnetic waves through the ISM

An example of the "cooling function" for predominantly neutral gas, as a function of temperature is shown for abundances appropriate to diffuse HI in the Milky Way, and for two different fractional ionizations: $x_{e=0.017}$. For $10 \leq T \leq 10^4 K$, the [CII] 158 μm fine structure line is a major coolant. The [OI]63 μm fine structure line is important for $T \geq 100\mathrm{K}$. Lyman α cooling dominates only at $T \geq 1 \times 10^4 K$.

Collisional excitation of ions in low-density plasma results in radiative cooling. The emitted power depends on the ionization state, and the plasma is often assumed to be in collisional ionization equilibrium, or CIE.

CIE assumes that the plasma is in a steady state, and that collisional ionization,

[15] Kulsrud, R. M., 2005, Plasma Astrophysics (Princeton: Princeton Univ. Press)

charge exchange, radiative recombination, and dielectronic recombination are the only processes altering the ionization balance, in which case the ionization fractions for each element depend only on the gas temperature, with no dependence on the gas density.

Winds from Cool Stars

The winds from cool stars can also produce bubbles. These winds are molecular and dusty, and typically have low outflow velocities $V_\omega \approx 15 - 30\mathrm{km}s^{-1}$. Mass loss rates can vary from $\sim 10^{-7} M_\odot yr^{-1}$ for M giant stars to $\sim 10^{-4} M_\odot yr^{-1}$ for AGB stars. The free expansion phase will end at

$$t_0 \approx (\frac{3\dot{M}}{4\pi\rho_0 V_\omega^3})^{1/2}$$

For $t \geq t_0$, the bubble expansion will decelerate. The gas is relatively dense, and molecules are effective coolants. As a result, these bubbles are generally strongly radiative, consisting of a supersonic stellar wind impacting a cold dense shell.

Cosmic Microwave Background

The cosmic microwave background (CMB) is very close to blackbody radiation with a temperature $T_{CMB} = 2.7255 \pm 0.0006\mathrm{K}$ [16]. The radiation is essentially isotropic, with the primary departure from isotropy consisting of a dipole perturbation due to motion of the Sun relative to the CMB rest frame with a velocity *v* = 372±1kms^{-1} toward *l* = (264.31±0.15)deg, b = (48.05±0.10)deg.

The position of the CMB acoustic peaks lie on the expand history from the decoupling epoch to the present epoch, and contains the information of dark energy.

Two distance ratios are often used to constrain dark energy. The first distance ratio is the so-called "acoustic scale" l_A, which represents the CMB multipole corresponds to the location of the acoustic peak. It can be calculated as

[16] Fixsen, D. J., 2009, "The Temperature of the Cosmic Microwave Background." Ap. J., 707, 916

$$l_A = (1+z_*)\frac{\pi D_A(z_*)}{r_s(z_*)}$$

Here z_* denotes the redshift of the photon decoupling epoch, $D_A(z)$ is the proper angular diameter distance, and r_s is the commoving sound horizon size.

The second distance ratio is the so-called “shift parameter” R which takes the form

$$R\,(z_*) = \sqrt{\Omega_{m0}}H_0(1+z_*)D_A(z_*)$$

Baryon Acoustic Oscillations

Baryon acoustic oscillation (BAO)[17], which happens on some typical scales, provides a “standard ruler” for length scale in cosmology to explore the expansion history of the universe.

From measurements of the galaxy clusters, the BAO scales in both transverse and line of sight directions are obtained; they correspond to the quantities $r(z)/r_s(z)$ and $r_s(z)/H(z)$.

In addition, three characteristic quantities of BAO, including the *A* parameter, $D_v(0.35)/D_v(0.2)$, and $r_s(z_d)/D_V(z)$, are often used to constrain dark energy parameters.

The A parameter is defined as

$$A_{th} = \frac{\sqrt{\Omega_{m0}}}{(H(z_b)/H_0)^{\frac{1}{3}}}[\frac{1}{z_b\sqrt{|\Omega_{k0}|}}f_k(H_0\sqrt{|\Omega_{k0}|}\int_0^{z_b}\frac{dz'}{H(z')})]^{\frac{2}{3}}$$

Let us turn to the quantity.

$$D_V(z) = [(1+z)^2 D_A^2 \frac{z}{H(z)}]^{1/3}$$

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[17] Refer to an overdensity or clustering of baryonic matter at certain length scales due to acoustic waves which propagated in the early universe.

We introduce the quantity.

$$z_d = \frac{1291(\Omega_{m0}h^2)^{0.251}}{1+0.659(\Omega_{m0}h^2)^{0.828}}[1+b_1(\Omega_{b0}h^2)^{b_2}]$$

Holographic Universe

The Holographic Principle

Holographic principle asserts that the world can be understood as a hologram. In other words a theory of gravity is dual to a boundary field theory without dynamic gravity in one less dimensions.

‘t Hooft[18] and Susskind[19] realized that the black hole entropy can be understood as a dimensional reduction, or holographic principle. The entropy is viewed as degrees of freedom measured in Planck units, which lives on the surface of the strongly gravitating system

$$S_{matter} \leq \frac{A}{4G}$$

Holographic Dark Energy

The reasoning of last section can be applied to the vacuum, which leads to a holographic model of dark energy. Li[20] pointed out that if one take $L = 1/R_h$, where R_h is the future event horizon defined as

$$R_h = \alpha \int_t^\infty \frac{dt}{\alpha} = \alpha \int_\alpha^\infty \frac{d\alpha}{H\alpha^2}$$

the energy density becomes

$$\rho_{de} = 3c^2 M_P^2 R_h^2$$

which does behave as dark energy.

18 G. ‘t Hooft, arXiv: gr-qc/9310026.
19 L. Susskind, J. Math. Phys. 36 (1995) 6377
20 M. Li, Phys. Lett. B 603 (2004) 1.

Summary

How matter and energy shape our universe, and more fundamental, what this stuff that makes up everything really is. We can now consider the following model: the universe operates holographically with regard to optics and acoustics.

Energy interacts with constructive and destructive interference, to form holograms which we perceive as matter. Just as the optical holograms we make give us appearances of nonexistent 3D images, energy operating on a much more basic level of density, form holograms which we perceive as actual objects.

To review briefly, emissions and winds take on a stationary appearance, while energy continues to pass through the system, each successive emission and wind takes the position of the one before. The emissions and winds are generated in hologram reconstruction since, as the hologram continues to be illuminated or flowed over a period of time, the same emission and wind continues to be formed.

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