Electric charge asymmetry of the Universe and magnetic field generation

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If at an early stage of the evolution of the Universe the gauge symmetry of electromagnetism was spontaneously broken, an electric charge asymmetry would develop. After restoration of gauge invariance, the asymmetry should disappear so that the net electric charge density must vanish, the compensating charge being produced from the Higgs vacuum in the form of heavy charged particles. Energetic products of their decay would create an electric current and a local charge asymmetry. Alternatively, such an asymmetry could be created even if the electric current was always conserved but an asymmetry in another nonconserved charge existed. The primary currents which created the asymmetry as well as those damping it via plasma discharge could generate chaotic magnetic fields on astronomically interesting scales. These fields might be large enough to seed the observed magnetic fields in galaxies via a protogalactic dynamo.

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Spiral galaxies are known to possess magnetic fields with a characteristic length scale of the order of the disk scale length and amplitude of about 10^{-6} G. The origin of galactic magnetic fields remains an unsolved problem, although there are several interesting proposals of mechanisms for field generation on galactic scales (for a review, see [1]). The scenario of field amplification by the dynamo mechanism [2,3] needs for its realization a seed field which is not easy to create over galactic scales. It was shown in Ref. [4] that magnetic fields would be generated in the primeval plasma if there were primordial vorticity perturbations, but the origin of the latter is unclear. The suggestion has recently been made [5] that vorticity can be generated by cosmic strings. An alternative proposal utilizes ejecta from primordial massive stars which may carry stellar dynamo-generated fields that are amplified in scale via interaction and dispersal of supernova remnants to yield a seed field on 100 pc scales [6]. Several scenarios of magnetic field creation by inhomogeneities arising in the course of the phase transitions occurring during cooling down of the very early Universe have been considered [7–9], as well as those produced by macroscopic manifestations of parity violation in the process of black hole evaporation [10], by explicit breaking of electromagnetic gauge invariance due to coupling with gravity [11], and by nonminimal coupling of the electromagnetic field tensor to the inflaton field [12]. The challenge for all of these models is to generate seed fields whose energy density is a significant fraction of the cosmological energy density at early epochs: Galactic magnetic fields have an energy density approximately equal to that of the cosmic microwave background radiation.

Here we propose a new mechanism for creation of galactic magnetic fields which can potentially explain the observed features. Under certain conditions, an electric generator (battery) might be operating in the Universe, creating electric currents and a local charge asymmetry on cosmological or galactic scales. There are at least two possible mechanisms for this process. One explores the idea that electric charge was not conserved in the early Universe. To achieve this, electromagnetic gauge invariance should be (spontaneously) broken early in the course of the evolution of the Universe and restored later at lower temperatures. This assumption was put forward about a decade ago in Ref. [13] in order to solve the magnetic monopole problem.

During the period of broken gauge invariance, an electric charge asymmetry of the Universe could, or rather should, be generated in the same way as baryon asymmetry. Indeed, all of the essential ingredients for creation of charge asymmetry such as C and CP breaking, thermal nonequilibrium, and last but not least, charge nonconservation could be naturally realized in the phase of broken gauge invariance. It is essential for our model that the resulting electric asymmetry is generically spatially nonuniform. Consideration of the baryon asymmetry of the Universe (for a review, see Ref. [14]) shows that in many models baryon asymmetry could be spatially inhomogeneous over astronomically large scales. We believe that any type of charge asymmetry in the Universe was generated by essentially the same physical phenomena: charge nonconservation, C and CP breaking, and deviation from thermal equilibrium (though, as was understood relatively recently, none of them is obligatory). Hence one may expect that an electric or any other kind of charge asymmetry should possess the same kind of inhomogeneities which are inherent for baryon asymmetry. It is noteworthy that the scale of the inhomogeneities could be much larger than the horizon scale at an early stage, if the conditions for their generation were prepared during inflation.

Since the gauge invariance is broken spontaneously, some remnants of it always existed. This demands in particular that the charge asymmetry in the sector of physical particles be accompanied by an asymmetry in the vacuum that is equal in amplitude but opposite in sign. The latter, however, remains unobservable while $U(1)_{em}$ is broken. When the gauge invariance is restored, electric charge which was hidden in the vacuum reappears in the form of charged (and heavy) Higgs bosons and so the net electric charge density becomes zero. In this sense, the present scenario resembles the scenario of baryogenesis with conserved baryonic charge proposed in Ref. [15] (see also [14]). However, there is more to the present story. Although the charge density became zero everywhere after the gauge invariance $U(1)_{em}$ was restored, the number density of heavy charged Higgs particles was not homogeneous. The latter should be unstable unless special care is taken to ensure their stability. If the masses of these Higgs particles are high enough so that the mean free path of their decay products is larger than the typical wavelength of the inhomogeneity, the decays would create electric currents carried by energetic light particles. This current, however, does not produce a magnetic field in a homogeneous plasma in the approximation when the current is created by a superposition of spherically symmetric noninteracting sources. Inhomogeneities in the primeval plasma would distort the high symmetry of the current and give rise to a nonzero magnetic field B. These inhomogeneities might be the usual density inhomogeneities which initiated large-scale structure formation. The scattering on inhomogeneities might be small for the currents created by the energetic particles from the charged Higgs boson decays because the scattering cross section is inversely proportional to the square of the energy. After the Higgs bosons decayed and their relativistic decay products left the regions of higher Higgs boson number density, these regions would acquire some net charge excess that can be either positive or negative. This charge excess would be neutralized by the process of electric discharge in the primeval plasma. Thus there should be two types of currents in the plasma: the original ones created by the energy of the decays and the secondary ones of usual electromagnetic origin which are generated by the electrostatic forces induced by the inhomogeneities in the electric charge distribution. The latter appear due to the original inhomogeneities in number density of the charged Higgs bosons. The secondary currents would be probably more effective in magnetic field generation since they are proportional to the plasma conductivity which is not uniform due to the density inhomogeneities in the primeval plasma. The size of these inhomogeneities as well as those in the Higgs boson number density could be comparable to or larger than galactic scales since both kinds of inhomogeneities were presumably formed during the inflationary stage. Because of the inhomogeneous conductivity, the secondary currents were not as symmetric as the original currents, possessing a nonzero curl, and could create a magnetic field. As we will argue in what follows, this might be a promising mechanism for generation of stochastically distributed primordial magnetic fields which can serve as seeds of magnetic fields in galaxies.

Spontaneous breaking of gauge U(1) symmetry of electromagnetism can be realized by the usual Higgs mechanism. Contrary to the standard scenario, however, we need the symmetry to remain unbroken at low temperatures. Usually, the temperature corrections resulting from, say, the self-coupling $\lambda\phi^4$ of the Higgs field give a positive contribution to the effective mass of $\phi, \delta m^2 \approx \lambda T^2$, since λ is assumed to be positive to ensure the stability of the theory. If, however, there exists a coupling to another scalar field ϕ' of the form $\lambda' |\phi|^2 |\phi'|^2$, the constant λ' is not necessarily positive and the corresponding temperature corrections to the mass of ϕ may be negative. This would induce symmetry breaking at high temperature [16,17]. If the zero-temperature mass of ϕ is positive, the symmetry should be restored at low temperatures.

In the realistic electroweak model, the gauge symmetry of electromagnetism U(1) is included in the larger symmetry group $SU(2)_I \otimes U(1)_Y$ and the realization of our scenario may be slightly more complicated. One interesting possibility is that the electroweak gauge symmetry was always broken both at low (as in the standard approach) and high temperatures (as proposed here). It would inhibit, in particular, electroweak baryon nonconservation and might save a preexisting baryon asymmetry from being destroyed. In this way, the conclusion of Ref. [18] that all initial $(B+L)$ asymmetry was reduced to zero by B-nonconserving processes above the electroweak phase transition can be avoided.

Let β_e be the ratio of electric charge density generated at the epoch of broken U(1) to the entropy density:

$$\beta_e = N_e / S \ . \tag{1}$$

The analogous quantity characterizing the baryon asymmetry is measured to be $\beta_B = N_B / S = 10^{-9} - 10^{-10}$. Once electric charge asymmetry is generated, it is exactly compensated by the electric charge stored in the vacuum. To show this, let us consider the field equations for the coupled charged scalar and electromagnetic fields:

$$\partial_\mu^2 \phi - ie(\partial_\mu A_\mu)\phi - 2ie A_\mu \partial_\mu \phi - e^2 A_\mu^2 + \partial U / \partial \phi^* = 0 \ , \tag{2}$$

$$\partial_\mu F_{\mu\nu} + 2e^2 |\phi|^2 A_\nu = ie(\partial_\nu \phi^* \phi - \phi^* \partial_\nu \phi) + J_\nu^{\text{ext}} \ , \tag{3}$$

where J_ν^{ext} is an external electromagnetic current created, say, by fermions and U is the potential energy of ϕ which may include contributions from other fields and, in particular, from the coupling to fermions $(\phi \bar{\psi}_1 \psi_2 + \text{H.c.})$. Here the electric charge of ϕ is equal to the difference between the charges ψ_1 and ψ_2. Such an interaction would make the electromagnetic current of fermions nonconserved in the state with nonzero vacuum expectation value of the charged field ϕ, and a charge asymmetry can then develop in the fermionic sector. Note that the total current

$$J_\mu^{\text{tot}} = ie(\partial_\mu \phi^* \phi - \phi^* \partial_\mu \phi) + J_\mu^{\text{ext}} - 2e^2 |\phi|^2 A_\mu \tag{4}$$

is always conserved because $\partial_\mu J_\mu^{\text{tot}} = \partial_\mu \partial_\nu F_{\mu\nu} = 0$, and thus the total charge density remains zero even after a charge asymmetry in the fermionic sector has developed. It is interesting that in the static case the total electric charge should vanish in the broken-symmetry phase with nonzero $\langle \phi \rangle$. The easiest way to see this is to use the unitary gauge in which $\mathrm{Im}\phi=0$, and correspondingly the equations of motion take the form

$$\partial_\mu^2 \phi - e^2 A_\mu^2 \phi + \mathrm{Re}\frac{\partial U}{\partial \phi^*}=0 \ , \tag{5}$$

$$\partial_\mu F_{\mu\nu} + 2e^2|\phi|^2 A_\nu = J_\nu^{\mathrm{ext}} \ . \tag{6}$$

If there is a nonzero external electric charge density

$$J_\nu^{\mathrm{ext}} = \delta_{0\nu}\rho_e(r) \tag{7}$$

and $\langle \phi \rangle$ is spatially constant, the solution of Eq. (6) for the Coulomb-like potential is

$$A_0(\mathbf{r}) = \frac{1}{4\pi}\int \frac{d^3r'}{|\mathbf{r}'-\mathbf{r}|} e^{-m_\gamma|\mathbf{r}-\mathbf{r}'|}\rho_e(\mathbf{r}') \ , \tag{8}$$

where $m_\gamma = \sqrt{2}e\,|\phi|$ is the photon mass in this phase. Now it can be easily verified that

$$Q_{\mathrm{tot}} = \int d^3r [J_0^{\mathrm{ext}}(\mathbf{r}) - 2e^2|\phi|^2 A_0(\mathbf{r})] = 0 \ . \tag{9}$$

The medium behaves like a (super)conductor with an excess of charge driven to the surface.

When the symmetry is restored and the potential of ϕ effectively takes the form $U(\phi)=m_\phi^2|\phi|^2$, the equations of motion in the homogeneous case have the solution

$$A_\mu = \delta_{0\mu} m_\phi / e \ , \tag{10}$$

$$|\phi|^2 = J_0^{\mathrm{ext}}/2e^2 A_0 = J_0^{\mathrm{ext}}/2em_\phi \ . \tag{11}$$

This solution would be more transparent if the gauge transformation $A_0 \to A_0 - m_\phi/e = 0$ and $\phi \to |\phi|\exp(im_\phi t)$ is made. It shows that the mean charge density is zero (as expected) and that the electric charge stored in fermions is compensated by the charge produced from the vacuum in the form of real ϕ particles. Thus, when the gauge symmetry was restored, the heavy ϕ particles should be created, although the temperature might be well below their mass. A natural expectation is that these particles would almost instantly decay into energetic charged fermions, although one cannot exclude that there is a mechanism inhibiting the decay so that their lifetime is very large. In what follows, we will consider both possibilities.

It is possible and perhaps even more promising to make a cosmological electric battery without any electric charge nonconservation. Assume that there exists a new nonconserved charge Y corresponding to a global symmetry group G_Y. We assume for definiteness that the particles possessing this charge are rather heavy and not yet observed. One would expect that Y asymmetry of the Universe was generated at an early stage as happened with the baryon asymmetry. We assume also that the Y asymmetry manifests itself in a nonzero density of heavy and long-lived X particles ($Y_X \neq 0$). Since the symmetry G_Y is assumed to be only approximately valid, X particles should be generically unstable and their lifetime can be evaluated by rescaling the proton lifetime in grand unified models. The lifetime scales inversely proportionally to the fifth power of mass of the decaying particle. Thus, with the proton lifetime around 10^{40} sec, the lifetime of X would be about $10^{40}(m_X/m_p)^5$ sec. This estimate is clearly very approximate since the decay mechanisms and the masses of the intermediate bosons are most probably different in these two cases. Moreover, X might be a boson in contrast to a proton, but the hint is clear that the lifetime of X's possessing approximately conserved quantum number Y may be sufficiently large to satisfy the constraints presented below.

It is not essential for what follows if X is electrically neutral or charged. In the latter case, the electric asymmetry created by the X's should be compensated by, say, electrons or quarks. If the X's are electrically charged, the scenario of current generation is essentially the same as in the previous case. As we see in what follows, the effectiveness of the battery is substantially higher if X decays into heavy charged particles (the requirement for the mass of the latter is presented below).

If the X's are electrically neutral, they could still produce the battery in the case when they decay, say, into the channel $X^0 \to X^+ + l^-$, where X^+ is a heavy and (quasi)stable new charged particle and l^- is a light one such as a charged lepton or quark (we have chosen here a particular sign convention that is $X^0 \to X^+ + l^-$, while $\bar{X}^0 \to X^- + l^+$). In this case, the current is created due to different mean free paths of heavy and light fermions in the plasma. The mean free path of a charged particle in the primeval plasma is determined by Compton scattering or by e^+e^--pair production. In the ultrarelativistic limit, the cross section for Compton scattering on particles with mass m and charge e ($e^2 = 4\pi\alpha = 4\pi/137$) is given by the expression

$$\sigma_C = \frac{2\pi\alpha^2}{s}\ln\frac{s}{m^2} \ , \tag{12}$$

where $s=(p+k)^2 \approx 6ET$, with E being the energy of the charged particle and $\omega = 3T$ being the thermally averaged photon energy. The cross section is dominated by backward scattering when all the energy of the relativistic charged particle is transferred to the photon. The mean free path of the charged particles in these conditions is

$$l_{\mathrm{free}} \approx (\sigma_C N_\gamma)^{-1} \approx E/(\alpha^2 T^2 \ln s/m^2) \ .$$

It is larger than or comparable to the horizon size $l_h \approx t$ if $E \geq (10^{-3}-10^{-4})m_{\mathrm{Pl}}$; i.e., the mass of the decaying particle should be in the grand unification range. However, the energy of the decay products is redshifted down, and this invalidates the condition $l_{\mathrm{free}} > l_h$ in a few Hubble times. Hence, to create the electric currents on scales of interest for galactic magnetic field formation, the lifetime of the ϕ or X particles should be sufficiently high, $t_d > 10^6$ sec, and correspondingly the temperature at the moment of the decay should be sufficiently low, $T_d < 10^{-3}$ MeV. Note that for an early decay the galactic scale is much larger than the horizon at the moment of the decay. This makes it difficult to generate magnetic fields at the galactic scale.

If one demands that the energy density contained in the particles creating the electric current be smaller than, say, 10% of the energy density in the radiation, then this condition together with the lower bound on E mentioned above implies a strong upper limit on the charge asymmetry,

$$\beta_e < T/E \approx 10^{-23} . \tag{13}$$

Otherwise, there might be excessive distortions in the spectrum of the background radiation.

The limit (13) on β_e could be nonvalid if the mass of the charged decay products is sufficiently high, $m^2 > s - m^2 = 6ET$, although the particles in the comoving reference frame were still relativistic, $E \gg m$. In this case, the cross section of the photon elastic scattering is given by the Thomson limit $\sigma_T = 8\pi\alpha^2/3m^2$, and moreover the stopping power is further reduced by the factor $(p/\omega) \gg 1$ where p is the momentum of the charged particle and ω is the energy of the background photons in the comoving frame.

Another dangerous process is that of e^+e^--pair creation through the reaction $X^+ + \gamma \to X^+ + e^+ + e^-$ with a cross section which does not decrease with energy, $\sigma_{\text{pair}} \approx \alpha^3 \ln^3(s/m_e^2)/m_e^2$. However, for sufficiently heavy X^+, this process is not effective since the threshold for pair production is $E_X = m_X(m_e/3T)$ and is not reached for m_X and T in the range of interest. To summarize, this range is determined by the conditions

$$6m_1 T < m_2^2 , \tag{14a}$$

$$\beta m_1 < T , \tag{14b}$$

$$T < 30 m_2 (m_1/m_{\text{Pl}})^{1/2} , \tag{14c}$$

where m_1 is the mass of the decaying particle and m_2 is the mass of the heavy decay product. The first condition implies that the cross section of the interaction with the plasma is given by the Thomson limit. The second condition means that the energy density of the decaying particles is below the energy density of the radiation, and the third inequality ensures the large mean free path of the decay products. In particular, with m_1 and m_2 in the range 10^5-10^4 GeV, it is possible to satisfy all these restrictions with $\beta = 10^{-11}-10^{-12}$.

In both cases considered above, the initial value of the electric charge density was identically zero immediately after symmetry restoration and prior to the decays of ϕ's or X's. However, the number density of the heavy decaying particles might be different at different space locations due to possible inhomogeneities in $\beta_e = \beta_e(\mathbf{r})$ or $\beta_Y(\mathbf{r})$ (in what follows, we will not distinguish between β_e and β_Y and between ϕ and X since both mechanisms operate essentially along similar lines at this stage). When ϕ decayed and the decay products streamed away from the region of higher ϕ density, an inhomogeneous distribution of electric charge would have been created. The charge distribution is determined by two factors: by the current of particles from the ϕ decays (it is not of electromagnetic origin) and by the electric discharge tending to restore the neutrality of the plasma. The characteristic time and space scales of the discharge are determined by the plasma conductivity κ. The local (in space and time) relation between the electric field and the current density, $\mathbf{J} = \kappa \mathbf{E}$, can be used here since the characteristic frequency of the field variation is small in comparison with the plasma frequency, and the mean free path of the charged particles in the plasma is small in comparison with the distance over which the field changes considerably. The Maxwell equations governing these processes can be written as

$$\nabla \times \mathbf{B} = 4\pi \mathbf{J}^{\text{ext}} + 4\pi\kappa(x)\mathbf{E} + \partial_t \mathbf{E} , \tag{15}$$

$$\partial_t \mathbf{B} = -\nabla \times \mathbf{E} , \tag{16}$$

$$\nabla \cdot \mathbf{E} = 4\pi\rho , \tag{17}$$

$$\nabla \cdot \mathbf{B} = 0 , \tag{18}$$

where $\mathbf{J}^{\text{ext}}$ is the current induced by the ϕ decays. This current is transported by particles moving practically with the speed of light and is equal to

$$J^{\text{ext}} = e\,\delta\beta_e(R) S , \tag{19}$$

where $\delta\beta_e(R)$ is the inhomogeneity of the charge asymmetry on scale R and $S \approx T^3$ is the entropy density. This expression is valid on the boundary of this region; at a larger distance l, the current decays as $(R/l)^2$. The time duration of this pulse of current is of the order of R (we use the natural system of units with $c = h/2\pi = k = 1$). After the decays of ϕ mesons, their energetic decay products started to stream from the regions with an excess of ϕ's. This created a deficit of charge in this region with electric charge density ρ_e. In turn, it gave rise to the electric field E forcing the restoration of the plasma neutrality. The charge density ρ_e satisfies the equation

$$\partial_t \rho_e + 4\pi\kappa\rho_e = -\partial_i J_i^{\text{ext}} - E_i \partial_i \kappa . \tag{20}$$

In a realistic cosmological plasma, the conductivity κ is rather large, while its variation $\delta\kappa/\kappa$ is small. If they are due to the usual density inhomogeneities which gave rise to the large-scale structure formation, one would expect $\delta\kappa/\kappa \approx 10^{-4}$. This allows neglect of the second term on the right-hand side (RHS) of Eq. (20), and we obtain the solution

$$\rho_e(t) \approx \exp(-4\pi\kappa t) \int_{t_0}^{t} dt' \partial_i J_i^{\text{ext}}(t') \exp(4\pi\kappa t)$$
$$\approx \frac{\partial_i J_i^{\text{ext}}(t)}{4\pi\kappa} . \tag{21}$$

The induced charge density is inversely proportional to the conductivity, and the time necessary to reach the stationary state is small in comparison with R, $\Delta t = (4\pi\kappa)^{-1}$. This expression determines in particular the limiting value of the charge asymmetry that could be developed by the ϕ decays. The electrostatic potential induced by the nonzero charge density in the plasma should not exceed the energy of the decay products:

$$\frac{1}{R} \int d^3 r\, \rho = \frac{1}{4\pi\kappa R} \int d^3 r\, \partial_i J_i^{\text{ext}} = e\,\delta\beta_e T^3 R/\kappa < m_\phi . \tag{22}$$

If all the other quantities also reached their stationary values (although we see in what follows that this is not the case), the magnetic field would be determined by Eqs. (15)–(18) with $\partial_t \mathbf{E} = \partial_t \mathbf{B} = 0$, which reduce to

$$\partial_i(\kappa E_i) = -\partial_i J_i^{\text{ext}} , \tag{23}$$

$$\nabla \times \mathbf{E} = 0 . \tag{24}$$

The solution may be conveniently presented in the form

$$\mathbf{E} = (\mathbf{J}^{\text{ext}}/\kappa) + \mathbf{E}^{(1)} , \tag{25}$$

with

$$\partial_i(\kappa E_i^{(1)}) = 0 \tag{26}$$

and

$$(\nabla \times \mathbf{E}^{(1)})_i = -\left[\nabla \times \frac{\mathbf{J}^{\text{ext}}}{\kappa}\right]_i = \varepsilon_{iln}\frac{J_l^{\text{ext}}}{\kappa}\frac{\partial_n \kappa}{\kappa} \neq 0 . \tag{27}$$

Note that $\nabla \times \mathbf{J}^{\text{ext}} = 0$ and a nonzero RHS of Eq. (27) emerges from inhomogeneities in the conductivity κ. We can evaluate $E^{(1)}$ in order of magnitude as $E^{(1)} \approx (J^{\text{ext}}/\kappa)(\delta\kappa/\kappa)$, where $\delta\kappa(R)$ is the variation of the conductivity on scale R. Now, from Eq. (15), it follows that

$$B \approx 4\pi(\delta\kappa/\kappa)J^{\text{ext}}R . \tag{28}$$

With J^{ext} given by Eq. (19), we obtain

$$B/T^2 \approx 4\pi e(\delta\kappa/\kappa)\delta\beta_e(R)(RT) . \tag{29}$$

Note the presence of the enormous enhancement factor RT, which permits generation of magnetic field $B \approx T^2$ and energy density $\rho_B \approx \rho_\gamma$ even with $\delta\beta_e \le 10^{-20}$ on galactic scales. Unfortunately, there is not sufficient time to reach the stationary limit of B [Eq. (29)], as can be seen from Eq. (16), and hence the maximum possible value of B should be much smaller. Nevertheless, since the seed magnetic field can be amplified by the galactic dynamo mechanism by a factor of up to 10^{10} (or even higher) and the value of β_e could be as large as 10^{-11}, there is ample room to retreat below the estimate (29).

Let us turn now to the nonstationary case. The time evolution of B is governed by Eq. (16) where E is determined by the equation

$$\partial^2 E_i + 4\pi\kappa\partial_t E_i = -4\pi\partial_i\rho_e - 4\pi\partial_t J_i^{\text{ext}} . \tag{30}$$

Together with Eq. (20), this forms the complete set of equations determining E and ρ_e. These equations can be solved perturbatively in the cosmologically interesting case of large, slowly varying $\kappa(x)$. It is evident that the time scale for variation of E is the same as that of the charge density and is equal to $\Delta t = (4\pi\kappa)^{-1}$. The solution can again be written in the form (25) with $E^{(1)}$ satisfying the equation

$$\partial_t E_i^{(1)} = -\frac{1}{4\pi\kappa}\left[\partial_i\left[\frac{\partial_l J_l^{\text{ext}}}{\kappa}\right] + (\partial_t^2 - \partial_j^2)\left[\frac{J_i^{\text{ext}}}{\kappa}\right]\right] ; \tag{31}$$

thus, $E^{(1)}$ can be evaluated as

$$E^{(1)} \approx \delta\beta_e T^2 (T/\kappa)(t/R)(R\kappa)^{-1}$$

and its curl is smaller than that of (J^{ext}/κ) at least by a factor $(t/R)(\kappa R)^{-1} \ll 1$. This means that we can neglect $E^{(1)}$ in estimating the time evolution of the magnetic field,

$$\partial_t B_i = \epsilon_{ilm} J_l^{\text{ext}} \partial_m \kappa^{-1} , \tag{32}$$

and hence during the time interval $t \approx R$, B can rise at most up to

$$B \approx \delta\beta_e T^2 \frac{t}{R}\frac{T}{\kappa}\frac{\delta\kappa}{\kappa} . \tag{33}$$

This is much smaller than the estimate (29) of B if the stationary state were reached. Since the magnetic field is much smaller than the electric field, we can neglect the influence of the former on the conductivity at this stage.

This result strongly depends upon the plasma conductivity, and in what follows we evaluate it for different values of the temperature T. The electron-positron conductivity is controlled by scattering on the electromagnetic radiation with the Thomson cross section $\sigma_{\text{Th}} = 8\pi\alpha^2/3m_e^2$. It is assumed that the temperature is below the electron mass. At higher temperatures, the cross section of $e\gamma$ interactions is $\sim\alpha^2/T^2$ and the conductivity is very large, $\kappa \approx T/\alpha$. The mean free path of electrons in the radiation bath is equal to $l_f^{(e)} = (\sigma_{\text{Th}} N_\gamma)^{-1}$. Under the influence of the electric field, electrons are accelerated in accordance with the equation

$$m_e \dot{v}_E = eE . \tag{34}$$

The time during which this acceleration is efficient is approximately $t_f^{(e)} = l_f^{(e)}/\langle v_T \rangle$, where $v_T = (3T/m_e)^{1/2}$ is the thermally averaged electron velocity. Hence the drift velocity under the influence of the electric field E is $v_E = eEt_f^{(e)}/m_e$, and the conductivity is

$$\kappa_e = \frac{J_e}{E} = \frac{eN_e v_E}{E} = \frac{\sqrt{3}}{8\pi\alpha}\frac{N_e}{N_\gamma}\left[\frac{m_e}{T}\right]^{1/2} m_e . \tag{35}$$

This expression is valid for sufficiently weak electric fields such that the energy gained by an electron on its mean free path due to the acceleration in the electric field is smaller than its thermal energy (see, e.g., Ref. [19], pp. 139–140) or, in other words, that there are no runaway particles. We shall check the validity of this assumption on the final result.

If one literally takes the same expression for the proton conductivity with the substitution of m_p instead of m_e, one finds that the latter is $\kappa_p \approx 10^5 \kappa_e (N_p/N_e)$. This is not true, however, because of the Coulomb scattering of protons on electrons (and positrons). Note that the number density of electrons and positrons N_e could be much larger than $N_p \approx 3\times 10^{-10} N_\gamma$ and still $N_e \ll N_\gamma$ in the temperature range near $m_e/15$.

The proton conductivity can be found essentially along the same lines as above with some complications connected with the long-range character of the Coulomb force. The resulting expression can be found, e.g., in Refs. [19,20]:

$$\kappa_p = \left[\frac{2}{\pi}\right]^{3/2}\frac{N_p}{N_e}\left[\frac{T}{m_p}\right]^{1/2}\frac{T}{\alpha \ln\Lambda} , \tag{36}$$

where $\ln\Lambda$ is the so-called Coulomb logarithm which comes from the integration of the Rutherford cross section. Λ is equal to the product of the average momentum transfer $\Delta p = T$ and the Debye screening length $\lambda_D = (T/4\pi\alpha N_e)^{1/2}$. Thus

$$\ln\Lambda \approx 0.5\ln(N_\gamma/\alpha N_e) = 10-20 \ .$$

The minimum value of the total plasma conductivity, $\kappa \approx 5\times 10^{-5}T$, is achieved at $T_m \approx m_e/15$ and correspondingly $N_e/N_\gamma \approx 1.7\times 10^{-7}$ and $N_p/N_e \approx 1.7\times 10^{-3}$ (for $N_p/N_\gamma = 3\times 10^{-10}$). It means that the most favorable conditions for magnetic field generation are realized for wavelength R_m equal to the horizon l_h at $T = T_m$. In terms of the comoving scale, it is equal approximately to 3 kpc. The amplitude of the magnetic field, as given by Eq. (33), is $B/T^2 \approx 10^4(\delta\kappa/\kappa)\delta\beta_e$. This does not mean, of course, that the magnetic field is largest on this scale because the result depends also on the fluctuation spectrum $\delta\beta_e(R)$ and $\delta\kappa(R)$ and upon the details of the subsequent field amplification by the adiabatic compression and/or by the galactic dynamo. The scale R_m is rather low because the galaxies themselves do not follow the average Hubble expansion, and to get the present-day galactic size we need a scale two to three orders of magnitude larger. A magnetic field of the appropriate wavelength should be generated at $T \approx 10^{-3}m_e$ when the horizon was about $l_h = 4\times 10^6$ sec. The plasma conductivity at that time was dominated by proton conductivity, and to obtain the right amplitude of the magnetic field at the galactic scale with the subsequent dynamo amplification by, say, ten orders of magnitude, one needs

$$\delta\beta_e(\delta\kappa/\kappa) > 3\times 10^{-16} \ . \tag{37}$$

Now we test the validity of the assumption of the absence of runaway particles. The amplitude of the discharge current in the plasma is close to that of the original one from ϕ decays, $J^{\rm dis} \approx J^{\rm ext} \approx e\delta\beta_e N_\gamma$. This means that the drift velocity of the particles participating in $J^{\rm dis}$ is $v_E = \delta\beta_e(N_\gamma/N_B) = \delta\beta_e/\beta_B$ (here $\beta_B = 3\times 10^{-10}$ is the baryon asymmetry of the Universe). The kinetic energy of the drift in the electric field should be smaller than the thermal energy, $mv_E^2/2 < 3T/2$. It is satisfied if

$$\frac{3T}{m_p} > \left[\frac{\delta\beta_e}{\beta_B}\right]^2 . \tag{38}$$

Comparing this with the condition (37), we find that for successful magnetic field generation the inhomogeneities in the plasma conductivity should be rather high: $(\delta\kappa/\kappa) > 10^{-3}$ for $T = 10^{-3}m_e$ and $(\delta\kappa/\kappa) > 3\times 10^{-5}$ for $T = m_e/15$. Note, however, that the above mechanism might be operative even if the condition (38) is not valid.

An interesting possibility is the generation of magnetic field after hydrogen recombination at $T \leq T_{\rm rec} = 0.35$ eV ≈ 4000 K. Though κ_e [Eq. (35)] rises as $T^{-1/2}$, the ratio N_e/N_γ drops by 10^5 at recombination so that

$$\kappa_e \approx 5\times 10^{-5}(T_{\rm rec}/T)^{1/2}T_{\rm rec} \ .$$

The horizon volume during recombination includes 10^5-10^6 potential galaxies, and so the scale of the magnetic field is sufficiently large. At this stage one should expect larger variations in the plasma conductivity due to rising density inhomogeneities (up to $\delta\kappa/\kappa \sim 1$). Since the galactic size $l_{\rm gal}$ is much smaller than $R = l_h$, the time duration of the electric current pulse corresponding to the inhomogeneity of β_e on the scale R is also much longer than $l_{\rm gal}$ and this permits the enhancement of the magnetic field amplitude by the factor $t/l_{\rm gal} \approx 10^2$ [see Eq. (33)]. In the most optimistic case, one can get $(B/T^2) \approx 10^6\delta\beta_e(R)$.

These results suggest that the proposed mechanism can create large enough seed fields for the generation of the observed galactic magnetic fields. It is noteworthy that magnetic fields generated by the present mechanism are correlated with the inhomogeneities in the energy density of the primeval plasma so that the field amplitude should be larger in regions with higher matter density. Magnetic fields in the early universe could create a nonzero local vorticity because of the different interactions of protons and electrons with the radiation. Particles infalling into the region of higher matter density due to the gravitational attraction would acquire a nonzero angular momentum in the magnetic field (if it exists). The angular momentum gained by protons would be equal in magnitude and opposite in sign to that gained by the electrons. So, for a neutral flux of noninteracting particles, the net angular momentum should be zero. However, the much smaller friction force acting on the protons from the background radiation would destroy the balance. This process is in fact inverse to the one considered in Ref. [4] where it was assumed that magnetic fields were generated by vorticity in the early Universe. Note also that the decays of ϕ bosons could create energetic cosmic rays with characteristic energy

$$m_\phi/(z_d+1) \approx 3\times 10^4 \text{ GeV } (100 \text{ GeV}/T_d)(m_\phi/m_{\rm Pl}) \ ;$$

here, $(z_d+1) = T_d/T_0$ and $T_0 = 3$ K is the present-day temperature of the relic background radiation.

It is of interest to note that observations of global magnetic fields in spiral galaxies cannot yet distinguish between two rival hypotheses for field origin that appeal either to a dynamo that leads to a closed axisymmetric field [21,22] or to an open bisymmetric spiral configuration [23]. Both types of field configuration are observed, and since with reconnection even a primordial field can produce closed magnetic structures, it is not possible to make any definitive assertions about the required degree of amplification of primordial fields [24]. Without a dynamo, differential rotation accounts for ~ 100 windings of the primordial field, suggesting that the primordial field was about $\sim 10^{-8}$ G, while a protogalactic dynamo may succeed even if the pregalactic field is $10-15$ orders of magnitude lower. A more direct observation of early galactic fields comes from observations of the extragalactic component of Faraday rotation toward quasars, which appears to be associated with intervening metal-line absorption systems [25], generally presumed to be metal-poor clouds in galaxy halos. Equipartition field strengths are inferred at a redshift $z = 1-2$, implying that a relatively ordered field already was present in galaxy halos.

In conclusion, we have shown that an electric battery might be operating at an early stage of the evolution of the Universe. It should lead to creation of chaotic cosmological magnetic fields with length scale and amplitude appropriate for the explanation of seeds for the observed magnetic fields in galaxies. The spectral distribution of the magnetic field depends on the details of the electrogenesis and can be practically arbitrary starting from a scale-free form to one having a prominent peak at some particular wavelength. The spectrum should be considerably modified in the course of cosmological evolution, and in particular the fields in galaxies should be enhanced both due to operation of a protogalactic dynamo and because of the anticipated correlation with the matter inhomogeneities. The model is most easily realized if there are new quasistable neutral or charged particles with mass around 100 TeV as well as charged long-lived or even stable particles with mass an order of magnitude smaller. The latter, if they are stable, were proposed [26] as bearers of the hidden mass of the Universe. The search for charged dark matter [27] allows their existence in the desired region. The decays of X into Y and the subsequent decay of Y (if it is unstable) could give rise to observable features in the microwave background radiation.

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