Insight into Primordial Magnetic Fields from 21-cm line observation with EDGES experiment

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ABSTRACT
The recent observation of the 21-cm global absorption signal by the EDGES suggests that the intergalactic medium (IGM) gas has been cooler than the cosmic microwave background (CMB) during $15 \lesssim z \lesssim 20$. This result can provide a strong constraint on heating sources for the IGM gas during these redshifts. In this Letter we study the constraint on the primordial magnetic fields (PMFs) by the EDGES result. The PMFs can heat the IGM gas through their energy dissipation due to the magnetohydrodynamic (MHD) effects. By numerically solving the thermal evolution of the IGM gas with the PMFs, we find that the EDGES result gives a stringent limit on the PMFs as $B_{1\mathrm{Mpc}} \lesssim 10^{-10}$ Gauss.

Key words: magnetic fields – intergalactic medium – dark ages, reionization, first stars – diffuse radiation – cosmology: observations – cosmology: theory

1 INTRODUCTION
In the Universe, there exist magnetic fields on a wide range of length scales $L$, from the planets (Connerney 1993) and stars (Donati & Landstreet 2009) with $L \lesssim 1$ au, to the galaxy clusters (Carilli & Taylor 2002) and large scale structure (Vacca et al. 2018) with $L \gtrsim 1$ Mpc. Various observations have been revealing the nature of these magnetic fields. However, the origin and evolution of large scale magnetic fields have not been understood yet and they are one of the challenging problems in modern cosmology (Durrer & Neronov 2013), although the dynamo process might have played an important role in the amplification of large scale magnetic fields (Brandenburg & Subramanian 2005). The dynamo mechanism is an effective amplification mechanism even for small magnetic fields. However, this mechanism cannot create magnetic fields from nothing. Therefore, many works study the possibility that the origin of these magnetic fields could be tiny magnetic fields created in some physical phenomena in the early universe, called the “Primordial Magnetic Fields (PMFs)”. In the literature, the generation of the PMFs has been suggested in various cosmological epochs including inflationary era (Turner & Widrow 1988; Ratra 1992), preheating (Bassett et al. 2001), phase transition (Baym et al. 1996; Quashnock et al. 1989; Hogan 1983), topological defects (Sicotte 1997; Avelino & Shellard 1995), Harrison mechanism (Harrison 1970; Hutschenreuter et al. 2018), and the others scenarios (Rogachevskii et al. 2017; Sriramkumar et al. 2015) (see, e.g., Kandus et al. 2011; Subramanian 2016 for reviews).

Additionally, the possibility of the existence of the intergalactic magnetic fields with $10^{-20} - 10^{-15}$ G has been argued by some previous works in this decade (Ando & Kusenko 2010; Neronov & Vovk 2010; Takahashi et al. 2013; Tashiro et al. 2014), which motivates the study of the PMFs as the origin of magnetic fields in galaxy clusters and large scale structure.

Constraints on the PMFs can provide the information to their generation mechanism and, moreover, give a hint to the physics of the early universe. So far, many constraints on the PMFs have been obtained from different cosmological observations: big bang nucleosynthesis (BBN) (Cheng et al. 1996), cosmic microwave background (CMB) temperature anisotropy (Planck Collaboration et al. 2016b) and its polarisation (Zucca et al. 2017), the CMB spectral distortion (Jedamzik et al. 2000; Kunze & Komatsu 2013), the baryon-to-photon number constraints between the BBN and the recombination epoch (Saga et al. 2018), the Sunyaev-Zel’dovich effect (Shaw & Lewis 2012; Tashiro & Sugiyama 2011; Minoda et al. 2017), the galaxy number count (Tashiro et al. 2010), the star formation history (Marinacci & Vogelsberger 2016), the dipole anisotropy of the observed extragalactic ultra-high-energy cosmic-ray particles (Bray & Scaife 2018), and so on.

Now the measurement of cosmological 21-cm line signatures is also expected to be a useful tool to constrain
the PMFs. The 21-cm line signal depends on the physical states of neutral hydrogen gas such as its number density, its temperature and so on. The PMFs can provide the effect on them through the magnetohydrodynamic (MHD) effects (Sethi & Subramanian 2005). Tashiro & Sugiyama (2006) has pointed out that the measurements of 21-cm fluctuations by future radio interferometer telescopes can provide a strong constraint on the PMFs and has stimulated further detailed studies (Shiraishi et al. 2014; Kunze 2018).

Bowman et al. (2018) recently has reported the detection of the strong radio absorption signal around 78 MHz with the Experiment to Detect the Global Epoch of Reionization Signature (EDGES). This result indicates that the gas temperature in the intergalactic medium (IGM) was cooler than the CMB temperature at the corresponding redshifts, i.e., $15 \leq \nu \leq 20$. Therefore, the EDGES result can be interpreted as a constraint on a cosmological heating source. Note that the EDGES result is difficult to be explained by the standard scenario: the amplitude of the absorption trough reported by EDGES is almost two times larger than the maximal value expected in the standard cosmology. However, instead of giving the explanation on this anomaly, several authors have already applied the EDGES result to constrain some models: the Hawking evaporation of small Primordial Black Holes (PBHs) (Clark et al. 2018), the emission from the accretion disks around large PBHs (Hektor et al. 2018), the decaying (Clark et al. 2018; Mitridate & Podo 2018) or annihilating dark matter (Cheng et al. 2018; D’Amico et al. 2018), warm dark matter (Safarzadeh et al. 2018), primordial power spectrum (Yoshiura et al. 2018) and so on.

In this work, we study the implication of the EDGES result to the PMFs and derive a constraint on them. The PMFs work as extra cosmological heating sources through the so-called ambipolar diffusion, in particular, in the late universe. Several authors have already pointed out that the measurement of the global 21-cm signal can put the constraint on the PMFs (Schleicher et al. 2009; Sethi & Subramanian 2009). Following these works, we evaluate the global thermal history with the PMFs and obtain a constraint on the amplitude and the scale dependence of the PMFs.

In this Letter, we describe the time evolution of the 21-cm line signal, and discuss the impact of the PMFs on it in Sec. 3. We report a constraint on the PMFs from the result of EDGES experiment in Sec. 4, and finally summarize our findings and discuss the future outlook in Sec. 5. In the analysis of this paper, we assume a flat ΛCDM model and fix cosmological parameters as obtained by Planck 2015 (Planck Collaboration et al. 2016a): $\Omega_m = 0.308, \Omega_b = 0.048, h \equiv H_0/100 \text{ km s}^{-1} \text{ Mpc}^{-1} = 0.678$ where $\Omega_m$ and $\Omega_b$ are the density parameters for the matter and baryon, and $h$ is the reduced Hubble constant.

### 2 GLOBAL 21-CM LINE SIGNAL

In cosmological 21-cm line measurements, we observe the differential brightness temperature, which is given by

$$
\delta T_b(z) \approx 27T_{\text{HI}}(z)\left[1 - \frac{T_b(z)}{T_{\text{spin}}(z)}\right] \times \left(\frac{\Omega_b h^2}{0.02}\right) \times \left(\frac{1.15 z^{1/2} - 1 + z}{10}\right)^{1/2} \text{[mK]},
$$

where $T_{\text{HI}}$ is the neutral fraction of hydrogen, $T_b$ is the CMB temperature, and $T_{\text{spin}}$ is the spin temperature which is defined by the ratio between the populations of the two hyperfine levels of neutral hydrogen atoms (for a review, see Furlanetto et al. 2006). Since we are interested in the global 21-cm line signal, all of these quantities are treated as background ones.

When the spin temperature is the same as the CMB one, the 21-cm signal vanishes as shown in Eq. (1). In the cosmological context, there are two processes to make the spin temperature deviate from the CMB one. One is the collisional interaction and the other is the interaction with Ly-α flux field (Wouthuysen 1952; Field 1959). They can couple the hyperfine structure with the IGM gas temperature. Besides the CMB interaction (the 21-cm photon emission and absorption), these processes control the evolution of the spin temperature, depending on the redshift. As a result, the spin temperature evolves between the values of the CMB and the IGM gas temperatures.

In the standard cosmology, the spin temperature evolution can be divided into four regimes (Pritchard & Loeb 2012). First, after the decoupling of the gas temperature from the CMB one around $z \sim 20$, the spin temperature follows the gas temperature through the collisional coupling. Then at the second stage, as the universe evolves, the gas temperature and its number density decrease. Accordingly, the collisional interaction is dominated by the coupling to CMB. As a result, the spin temperature approaches to the CMB one. The resultant global 21-cm signal becomes faint in this regime. In the third regime, the first luminous objects play important roles. They can produce the strong Ly-α field and the Ly-α interaction becomes effective. Through this interaction, the spin temperature starts to couple with the gas temperature again. So far the gas temperature evolves adiabatically after the decoupling from the CMB one and the gas is cooler than the CMB. Accordingly the spin temperature also gets lower than the CMB one. Eq. (1) tells us that the global 21-cm signal becomes negative in this regime. In other words, the signals are observed as the absorption trough. After that, as the star and galaxy formation becomes active, a lot of UV photons are produced at this regime. They start to ionise and heat up the gas. Quickly the gas temperature increases and surpasses the CMB temperature. The spin temperature also becomes higher than the CMB and the signal becomes positive. This means that the signal is measured as emission. Gradually the IGM gas is ionised and, finally, the ionisation of the IGM is completed. At this point the global 21-cm signal totally vanishes again.

Recently, the EDGES experiment has reported the absorption signal of global redshifted 21-cm lines between redshifts $15 \leq \nu \leq 20$. This result would indicate that the neutral hydrogen hyperfine structure coupled well with the gas tem-
temperature through Ly-α field and the gas temperature was lower than the CMB one around these redshifts. In the next section, we show how the PMFs affect the evolution of the IGM gas temperature.

3 THE FORMALISM OF THE PRIMORDIAL MAGNETIC FIELDS

Since our final aim is to constrain the strength and scale-dependence of the PMFs, we provide their statistical properties on which we focus in this Letter. First we describe the evolution of magnetic fields as \( B(t, \mathbf{x}) = f(t)B_{\text{init}}(\mathbf{x})/a^2(t) \) where \( a(t) \) is the scale factor normalized as \( a(0) = 1 \) at the present time \( t_0 \). Here a factor \( a^{-2}(t) \) comes from the adiabatic evolution of the PMFs while \( f(t) \) represents the evolution due to the energy dissipation. We set \( f(\text{init}) = 1 \) and we will discuss more details later. Next we assume the PMFs as statistically homogeneous and isotropic random Gaussian fields. Therefore the statistical property of the PMFs is completely determined by only the power spectrum, \( P_B(k) \), as Landau & Lifshitz (1980)

\[
\langle B_x^*(k)B_y(k) \rangle = \frac{(2\pi)^3}{2} \delta_D(\mathbf{k} - \mathbf{k'}) \delta_{ij} - \hat{k}_i \hat{k}_j)P_B(k),
\]

where \( B_x^*(k) \) is a Fourier component of \( B_{\text{init}}(\mathbf{x}) \) with a mode \( \mathbf{k} \) and we assume that the PMFs are non-helical fields.

For simplicity, we assume that the power spectrum of the PMFs is given by the single power-law in \( k \) space as

\[
P_B(k) = \frac{(2\pi)^n a_n + 5}{f(\pi a_n + 2)} B_0^{n B} k^{n B + 3} \quad \text{for} \quad k < k_{\text{cut}}. \tag{3}
\]

Here, we set \( k_{\text{cut}} = 2\pi \, \text{Mpc}^{-1} \) as a reference scale, \( B_0 \) is the amplitude of the PMFs normalized at the length of 1 Mpc, and \( n_B \) gives the scale dependence. The case of \( n_B = 3 \) corresponds to the scale-invariant PMFs and we adopt \( n_B > 3 \) throughout this paper to avoid the infrared divergence of the PMFs and \( f(x) \) represents the gamma function.

On the other hand, in the ultraviolet regime, the PMFs have the cut-off scale due to the MHD effect in the early universe. According to the previous studies (Jedamzik et al. 1998; Subramanian & Barrow 1998), the radiative diffusion before the cosmic recombination epoch damps the PMFs on small scales. The damping scale increases as the universe evolves. Therefore, we assume that the power spectrum has the sharp cut-off on the damping scale at the recombination epoch,

\[
P_B(k) = 0 \quad \text{for} \quad k > k_{\text{cut}}. \tag{4}
\]

Here the cut-off scale is given by

\[
k_{\text{cut}} = \left( \frac{V_A}{\sigma_T} \right)^{1/2} \left( \frac{\pi^{1/2} \Gamma(n_B/2)}{\Gamma(n_B)} \right)^{1/2} t_\text{rec} / \Omega_m h^3 a^3(\mathbf{x}(t)) \approx 3.32 \times 10^{-3} \left( \frac{B_0}{1 \, \mu\text{G}} \right)^2 \left( \frac{\Omega_m h^2}{0.02} \right)^{-1/2} \left( \frac{\sigma_T}{0.15} \right)^{1/2}. \tag{5}
\]

where \( V_A \), \( \sigma_T \), \( n_e \) and \( t_{\text{rec}} \) are the Alfvén velocity, the cross section for Thomson scattering, the electron number density, and the cosmic time at the recombination epoch, respectively. In constraining the PMFs, it is useful to introduce \( B_A \), which represents the magnetic field strength smoothed at any spatial length \( \lambda \). One can easily find that \( B_A \) is related to \( B_n \) as

\[
B_A^2 = \int_0^{\infty} e^{-k^2 x^2} P_B(k) \frac{d^3 k}{(2\pi)^3} = B_n^2 \left( \frac{k_{\text{cut}}}{k_n} \right)^{n_B + 3}, \tag{6}
\]

where \( k_A = 2\pi / \lambda \). In Eq. (6), we choose the Gaussian window function in Fourier space.

After the recombination epoch, the magnetic fields dissipate their energy and heat the IGM gas through two processes (Sethi & Subramanian 2005). One is the decaying turbulence and the other is the ambipolar diffusion. Sethi & Subramanian (2005) has shown that, while the first one is active around the recombination epoch, the latter becomes effective at the late universe \((z < 500)\). The evolution of the IGM gas temperature \( T_{\text{gas}} \) with the heating from the PMFs is given by

\[
\frac{dT_{\text{gas}}}{dt} = -2HT_{\text{gas}} + \frac{x_e}{\eta_\gamma} \frac{8\rho_c \Omega_c}{m_e c^2} (T_r - T_{\text{gas}}) + \frac{2}{3k_B n_h} (\dot{Q}_{\text{AD}} + \dot{Q}_{\text{DT}}), \tag{7}
\]

where \( H \), \( x_e \), \( \rho_c \), \( m_e \), \( c \), \( k_B \) and \( n_h \) is the Hubble parameter, the ionisation fraction of the baryon gas, the energy density of CMB photons, the rest mass of an electron, the speed of light, the Boltzmann constant, and the number density of the baryon gas, respectively. \( \dot{Q}_{\text{AD}} \) and \( \dot{Q}_{\text{DT}} \) in the right hand side represent the global heating rate due to the ambipolar diffusion (AD) and the decaying turbulence (DT), respectively. These two heating rates are written as Sethi & Subramanian (2005)

\[
\dot{Q}_{\text{AD}} = \frac{|(\nabla \times \mathbf{B}) \times \mathbf{B}|^2}{16\pi^2 c^3} x_e \tag{8}
\]

\[
\dot{Q}_{\text{DT}} = \frac{3w_B}{2} \frac{H}{8\pi} |\mathbf{B}|^2 \left[ \frac{\log (1 + t_d / t_{\text{rec}})}{\log (1 + t_d / t_{\text{rec}})} \right]^{w_B} \tag{9}
\]

with the drag coefficient \( \xi = 1.9 \left( T_{\text{gas}}/1 \, K \right)^{0.375} \times 10^{14} \, \text{cm}^3 \, \text{s}^{-1} \) as referred in Schleicher et al. (2008), the mass density of the baryon gas \( \rho_B \), the time-dependence of the decaying turbulence \( w_B \equiv (n_B + 3)/(n_B + 5) \), and the physical time-scale for the decaying turbulence \( t_d \). Here, since we assume a blue spectrum of the PMFs, \( n_B > -3 \), we can take \( t_d = (k_{\text{cut}} V_A)^{-1} \) as the Alfvén time scale at the cut-off scale of the PMFs, as referred in Sethi & Subramanian (2005). Also, we can calculate the absolute value of the Lorentz force and the magnetic energy in Eqs. (8) and (9) at a time \( t \) by substituting Eq. (3) with \( f(t) = |(\nabla \times \mathbf{B}) \times \mathbf{B}| \) and \( |\mathbf{B}| \) as

\[
\mathbf{B}(k) = \frac{4}{3} \int_0^{t} \frac{d^3 k}{(2\pi)^3} f(k) \mathbf{B}(k) P_B(k) \mathbf{B}(k), \tag{10}
\]

and \( |\mathbf{B}|^2 = \int (dk / 2\pi)^3 f(k) P_B(k) \mathbf{B}(k) \mathbf{B}(k) f^4(\xi + c)^2 \). In this study, we neglect following two effects on the thermal evolution of the IGM gas included in the previous work (Schleicher et al. 2005).

\footnote{We have used this definition of the smoothed amplitude of the PMFs from Planck Collaboration et al. (2016b). Although this definition is slightly different from some other works (such as Fedeli & Moscardini 2012; Saga et al. 2018), the value of \( B_A \) is almost unchanged among different definitions.}

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First, we do not take into account any radiative cooling effects, such as collisional excitation and ionisation, recombination, and bremsstrahlung. We have confirmed these terms are negligible for the PMF model parameters of our interest. The other assumption is that there are no astrophysical objects. We mention this point in Sec. 5.

In order to solve Eq. (7), we also need to follow the evolution of the ionisation fraction,

\[
\frac{dx_e}{dt} = \frac{1 + K_\alpha \Lambda n_e (1 - x_e)}{1 + K_\alpha (\Lambda + \beta_e) n_e (1 - x_e)} 
\times \left[ -\alpha_e n_e x_e^2 + \beta_e (1 - x_e) \exp \left( -\frac{3E_{\text{ion}}}{4k_B T_g} \right) \right] + \gamma_e n_e x_e ,
\]

(10)

where \( K_\alpha, \Lambda, \alpha_e, \beta_e \) and \( \gamma_e \) are the parameters for the ionisation and the recombination processes. For these parameters, we adopt the functions in Seager et al. (1999) and Seager et al. (2000) with the modifications suggested in Chluba et al. (2015). In Eqs. (7) and (10), we neglect the presence of the primordial helium and the heavier elements for simplicity.

Finally, we comment on the time evolution of the PMFs. In this study, by introducing \( f(t) \), we consider the back-reaction from the PMF dissipation into the heating efficiency in Eqs. (8) and (9) as well. We can determine the time evolutionary function \( f(t) \) by solving the equation of the magnetic energy conservation,

\[
\frac{d}{dt} \left( \frac{|B|^2}{8\pi} \right) = -4H \frac{|B|^2}{8\pi} - \dot{Q}_{\text{AD}} - \dot{Q}_{\text{DT}} .
\]

Solving Eqs. (7), (10) and (11) simultaneously, we can uniquely obtain the evolution of \( T_{\text{gas}}, x_e \), and \( f(t) \) for a given PMF model with \((n_B, B_n)\).

4 RESULTS

Assuming the power spectrum of the PMFs, we solve Eqs. (7), (10) and (11) numerically to follow the gas temperature, the ionisation fraction, and the magnetic energy with setting the initial time \( t_{\text{init}} \) to the recombination epoch.

We show the evolution of the gas temperature for different \( B_n \) with the nearly scale-invariant PMFs \((n_B = -2.9)\) in Fig. 1. Even with the sub-nano Gauss PMFs, the IGM gas temperature is strongly affected by the PMFs. In this case, the gas temperature starts to be heated around \( z \sim 200 \), and is well deviated from the one in the case without the PMFs \((B_n = 0 \text{ nG})\).

We have confirmed that the ambipolar diffusion contributes to the heating of the IGM more than the decaying turbulence at the redshifts of our interest. Also the redshift dependence of the heating term in Eq. (8) follows as, \( 2\dot{Q}_{\text{AD}}/3k_B n_e \propto (1 + z) \). On the other hand, the dominant cooling term in the right hand side of Eq. (7) for \( z \ll 200 \) is the first one, and \( -2HT_{\text{gas}} \propto (1 + z)^{3/2} \), which decreases faster than the heating term due to the ambipolar diffusion. Therefore as we take larger values for \( B_n \), the heating due to the ambipolar diffusion gets effective and dominates the cooling terms at higher redshifts. After the IGM is heated well and the magnetic energy is considerably dissipated, the heating source terms in Eq. (8) become decreased. As a result, the gas temperature starts to deviate from the adiabatic evolution \((B_n = 0 \text{ nG})\) at first, and the IGM is gradually getting cool after the saturation. We find the larger value of \( B_n \) makes the earlier and strong heating, and the value of \( n_B \) determines the duration of the PMF heating.

The EDGES experiment has reported the detection of the global 21-cm absorption signals in the redshifts range \( 15 \leq z \leq 20 \) (Bowman et al. 2018). In Fig. 1 we highlight the redshift region of the EDGES absorption signals with the blue shade.

To create the absorption signal, the IGM gas temperature should be lower than the CMB temperature as shown in Eq. (1). Therefore, when the absorption signal is detected at the redshift \( z_{\text{abs}} \), we can exclude a PMF parameter region \((n_B, B_n)\) which gives \( T_{\text{gas}} > T_e \) at \( z_{\text{abs}} \). According to the EDGES result, we set \( z_{\text{abs}} = 17 \) which is the central redshift of the absorption signal in the EDGES experiment. Calculating the evolution of the gas temperature for different \( B_n \) with fixing \( n_B \), we get a novel constraint on the PMFs from the condition \( T_{\text{gas}} < T_e \) at \( z_{\text{abs}} = 17 \). The obtained constrained are \( B_n \leq 1.2 \times 10^{-7} \text{ nG} \) for \( n_B = -2.9, B_n \leq 6.3 \times 10^{-3} \text{ nG} \) for \( n_B = -2.0, B_n \leq 2.0 \times 10^{-4} \text{ nG} \) for \( n_B = -1.0 \). We can fit our new constraint in a linear relation between \( B_n \) and \( n_B \) as \( \log(B_n/1 \text{ nG}) \leq -(3n_B + 10)/2 \) for \( -3.0 < n_B < -1.0 \). We plot our PMF constraint on the \((n_B, B_n)\) plane in Fig. 2 with a solid line. For comparison, we also show the constraint from the Planck collaboration (Planck Collaboration et al. 2016b) and the one from the magnetic reheating before the recombination (Saga et al. 2018) in the dotted and dashed lines, respectively. In the range of \( -3.0 < n_B < -2.0 \), the EDGES experiment could constrain the PMF amplitude most tightly.

Finally, by varying \( z_{\text{abs}} \) from 15 to 20, we have investigated the dependence of the PMF constraint on \( z_{\text{abs}} \). We have found out that the resultant upper limit of \( B_n \) for any
fixed $n_B$ differs less than 10 per cent for different $z_{abs}$ between 15 and 20. We can conclude that the dependence on $z_{abs}$ is very weak.

5 CONCLUSION

In this Letter, we obtained a novel constraint on the PMFs from the result of the EDGES 21-cm global signal. The EDGES has reported the detection of the 21-cm absorption signals in the redshifts between $15 \leq z \leq 20$. This result suggests that the IGM gas is cooler than the CMB during this epoch. The PMFs can heat up the IGM gas by the dissipation of their energy through the ambipolar diffusion and the decaying turbulence. Therefore, the EDGES result can provide the constraint on the PMFs.

We have numerically evaluated the thermal evolution of the IGM gas with the PMFs. By requiring $T_{gas} < T_{T}$ at $z_{abs} = 17.0$ at which the EDGES absorption profile is centred, we have obtained the stringent upper bound on the PMFs roughly about $B_n \lesssim 0.1 \, \text{nG}$. We also find that this PMF constraint is not changed very much if the absorption redshift is deviated from $z_{abs} = 17.0$.

To end this Letter, we briefly discuss the impact of the astrophysical objects on the PMF constraint, which is not considered in our calculation. Of course they can significantly heat up and ionise the IGM gas, and the constraint from the 21-cm absorption signal might become tighter if we include such astrophysical processes (e.g., the star formation, the AGN activity, and so on). On the other hand, some previous works have pointed out that the condition of the astrophysical object formation is also affected by the PMFs (Wasserman 1978; Shibusawa et al. 2014). In order to calculate the global 21-cm signal history with these effects, we should solve the fully non-linear MHD equations and the realistic astrophysical processes (Minoda et al. 2017; Kunze 2018). We leave these points to future work.

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REFERENCES

Ando S., Kusenko A., 2010, ApJ, 722, L39
Avelino P. P., Shellard E. P. S., 1995, Phys. Rev. D, 51, 5946
Bassett B. A., Pollifrone G., Tsujikawa S., Viniegra F., 2001, Phys. Rev. D, 63, 103515
Baym G., Bödeker D., McLerran L., 1996, Phys. Rev. D, 53, 662
Bowman J. D., Rogers A. E. E., Monsalve R. A., Mozdzen T. J., Mahesh N., 2018, Nature, 555, 67
Brandenburg A., Subramanian K., 2005, Phys. Rep., 417, 1
Bray J. D., Scaife A. M. M., 2018, ApJ, 861, 3
Carilli C. L., Taylor G. B., 2002, ARA&A, 40, 319
Cheng B., Olinto A. V., Schramm D. N., Truran J. W., 1996, Phys. Rev. D, 54, 4714
Cheng H.-C., Li Z., Zheng R., 2018, Journal of High Energy Physics, 9, 98
Chiba J., Paoletti D., Finelli F., Rubiño-Martín J. A., 2015, MNRAS, 451, 2244
Clark S. J., Dutta B., Gao Y., Ma Y.-Z., Strigari L. E., 2018, Phys. Rev. D, 98, 043006
Connerney J. E. P., 1993, J. Geophys. Res., 98, 18
D’Amico G., Panci P., Strumia A., 2018, Phys. Rev. Lett., 121, 011103
Donati J.-F., Landstreet J. D., 2009, ARA&A, 47, 333
Durrer R., Neronov A., 2013, A&ARv, 21, 62
Fedeli C., Moscardini L., 2012, J. Cosmology Astropart. Phys., 11, 055
Field G. B., 1959, ApJ, 129, 536
Furlanetto S. R., Oh S. P., Briggs F. H., 2006, Phys. Rep., 433, 181
Harrison E. R., 1983, Phys. Rev. Lett., 51, 1488
Hutschenreuter S., Dorn S., Jasche J., Vazza F., Paoletti D., Lavaux G., Enßlin T. A., 2018, Classical and Quantum Grav- ity, 35, 154001
Jedamzik K., Katalinić V., Olinto A. V., 1998, Phys. Rev. D, 57, 3264
Jedamzik K., Katalinić V., Olinto A. V., 2000, Phys. Rev. Lett., 85, 706
Kandus A., Kunze K. E., Tsagas C. G., 2011, Phys. Rep., 505, 1
Kunze K. E., 2018, preprint, (arXiv:1805.10943)
Kunze K. E., Komatsu E., 2014, J. Cosmology Astropart. Phys., 1, 009
Landau L. D., Lifshitz E. M., 1980, Statistical physics. Pt.1, Pt.2
Mariotti M., Padoan P., 2018, J. Cosmology Astropart. Phys., 5, 069
Neronov A., Vovk I., 2010, Science, 328, 73
Planck Collaboration et al., 2016a, A&A, 594, A13
Planck Collaboration et al., 2016b, A&A, 594, A19
Privitera J. R., Loeb A., 2012, Reports on Progress in Physics, 75, 086901
Quashnock J. M., Loeb A., Spergel D. N., 1989, ApJ, 344, L49
Ratra B., 1992, ApJ, 391, L1

Figure 2. The upper limit of the PMFs parameters obtained by the recent observation of 21-cm line absorption by EDGES experiment (solid). The gray shaded region is the excluded region by previous works; Planck Collaboration et al. (2016b) in the dotted line and Saga et al. (2018) in the dashed line.
Rogachevskii I., Rachayskiy O., Boyarsky A., Fröhlich J., Kleerorin N., Brandenburg A., Schober J., 2017, ApJ, 846, 153
Safarzadeh M., Scannapieco E., Babul A., 2018, ApJ, 859, L18
Saga S., Tashiro H., Yokoyama S., 2018, MNRAS, 474, L52
Schleicher D. R. G., Banerjee R., Klessen R. S., 2008, Phys. Rev. D, 78, 083005
Schleicher D. R. G., Banerjee R., Klessen R. S., 2009, ApJ, 692, 230
Seager S., Sasselov D. D., Scott D., 1999, ApJ, 523, L1
Seager S., Sasselov D. D., Scott D., 2000, ApJS, 128, 407
Sethi S. K., Subramanian K., 2005, MNRAS, 356, 778
Sethi S. K., Subramanian K., 2009, J. Cosmology Astropart. Phys., 11, 021
Shaw J. R., Lewis A., 2012, Phys. Rev. D, 86, 043510
Shibusawa Y., Ichiki K., Kadota K., 2014, J. Cosmology Astropart. Phys., 8, 017
Shiraishi M., Tashiro H., Ichiki K., 2014, Phys. Rev. D, 89, 103522
Sicotte H., 1997, MNRAS, 287, 1
Sriramkumar L., Atmjeet K., Jain R. K., 2015, J. Cosmology Astropart. Phys., 9, 010
Takahashi K., Mori M., Ichiki K., Inoue S., Takami H., 2013, ApJ, 771, L42
Tashiro H., Sugiyama N., 2006, MNRAS, 372, 1060
Tashiro H., Sugiyama N., 2011, MNRAS, 411, 1284
Tashiro H., Takahashi K., Ichiki K., 2010, preprint, (arXiv:1010.4407)
Wasserman I., 1978, ApJ, 224, 337
Wouthuysen S. A., 1952, AJ, 57, 31
Yoshiura S., Takahashi K., Takahashi T., 2018, Phys. Rev. D, 98, 063529
Zucca A., Li Y., Pogosian L., 2017, Phys. Rev. D, 95, 063506

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