Ultrahigh energy neutrinos from population III stars: Concept and constraints

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In this paper we reconsider the model of neutrino production during the "bright phase," first suggested in 1977, in the light of modern understanding of both the role of Pop III stars and of acceleration of particles in supernova shocks. We concentrate on the production of cosmogenic ultrahigh energy neutrinos in supernova (SN) explosions that accompany the death of massive Population III stars. Protons are assumed to be accelerated at such SN shocks and produce neutrinos in collisions with CMB photons. In the calculations we deliberately use simplified assumptions which make the physical results transparent. Pop III stars, either directly or through their SN explosions, are assumed to be responsible for the reionization of the universe as observed by WMAP. Since the evolution of massive Pop III stars occurs on time scales much shorter than the Hubble time H^{-1}, we consider the burst of UHE proton production to occur at fixed redshift ($z_b = 10$ and $z_b = 20$), though more realistic models can easily be built. We discuss in some detail the problems involved in the formation of collisionless shocks in the early universe as well as the acceleration of charged particles in a medium that has potentially very low preexisting magnetization, if any at all. The composition of the accelerated particles in Pop III stars explosions is expected to be proton dominated, based upon the predictions of BBN and the Hydrogen-enhanced stellar-wind from primary Pop III stars. A simple calculation is presented to illustrate the fact that the diffuse neutrino flux from the bright phase burst is concentrated in a relatively narrow energy interval, centered at $E_\nu^c = 7.5 \times 10^{15}(20/z_b)^2$ eV. The ν_μ flux may be detectable by IceCube without violating the cascade upper limit and without exceeding the expected energetics of SNe associated with Pop III stars. A possible signature of the neutrino production from Pop III stars may be the detection of resonant neutrino events ($\bar{\nu}_e + e^- \to W^- \to$ hadrons) at energy $E_0 = 6.3 \times 10^{15}$ eV. For the burst at $z_b = 20$ and $\bar{\nu}_e$-flux at the cascade upper limit, the number of resonant events in IceCube may be as high as 10 events in 5 years of observations. These events have equal energies, $E = 6.3 \times 10^{15}$ eV, in the form of e-m cascades. Taking into account the large uncertainties in the existing predictions of cosmogenic neutrino fluxes at $E > 10^{15}$ eV, we argue that UHE neutrinos from the first stars might become one of the most reliable hopes for UHE neutrino astronomy.

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I. INTRODUCTION

Ultrahigh Energy (UHE) neutrino astronomy at energies above 10^{16}–10^{17} eV may open a new observational window on the universe. The most reliable prediction for the diffuse fluxes of these neutrinos can be achieved for the cosmogenic neutrinos [1], produced in $p\gamma$ collisions of UHE protons with CMB photons (see [2,3] for recent calculations and reviews). Such reliability mainly follows from three factors: (i) the flux of parent protons at redshift z of neutrino production can be normalized to the observed flux of protons at present (redshift $z = 0$); (ii) the number density and energy spectrum of the target (CMB) photons are both known with high accuracy at any cosmological epoch z; (iii) the interactions of UHE protons with CMB photons occur at center of mass (cm) energy that does not exceed ~GeV, well accessible to laboratory experiments and hence well known.

On the other hand, the numerical predictions of neutrino fluxes at Earth are affected by huge uncertainties, mainly due to the unknown evolution with redshift of the source space density and luminosity, the generation spectrum ($\propto E^{-\gamma_g}$), and the maximum energy ($E_{\max}^{\rm acc}$) of accelerated protons. Because of these uncertainties, any upper bound on the flux of cosmogenic neutrinos becomes of the highest importance. For cosmogenic neutrinos this limit is not provided by the Waxman-Bahcall bound [4], because the latter is imposed by the parent proton flux observed at $z = 0$ and fixed in the calculations. Meanwhile the calculated flux can differ by orders of magnitude due to uncertainties in γ_g, $E_{\max}$, and the evolution of sources [5].

A more general upper limit, valid, in particular, for the cosmogenic neutrino flux is the *cascade upper limit* [6]. The production of UHE neutrinos from pion decays is accompanied by the production of UHE electrons, positrons, and photons, which start an e-m cascade due to the collisions with CMB photons. The HE photons, remnants of these cascades, contribute to the diffuse gamma-radiation with gamma rays of energy reaching several hundreds GeV. Recent measurements of the diffuse gamma-ray flux by the Fermi-LAT telescope [7] was used in Ref. [8] to put a stronger upper limit on the diffuse neutrino flux, which excluded many models with detectable cosmogenic neutrino flux.

The detectability of cosmogenic neutrinos is further disfavored by the recent Auger measurements of UHECR mass composition [9]: both $X_{\max}$ and RMS data strongly suggest that at energies larger than 2×10^{18} eV the primaries are nuclei with steadily increasing mass number A. In contrast with these data the HiRes [10] and Telescope Array [11] mass composition agrees well with a pure proton composition. The cosmogenic neutrino flux produced by primary nuclei is considerably lower than in the case of primary protons [12]. The Auger results on mass composition, together with the Fermi-based neutrino upper limit suggest a rather grim possibility of detection of cosmogenic neutrinos.

Here we revise a model of UHE neutrino production related to the so-called Population III (Pop III) stars ([13–16]). Now this population of stars attracted much attention as the most plausible sources of ionizing photons in the early universe, at redshift $z \sim 6$–20, as observed by WMAP [17].

Originally, one of the main motivations that led to propose the existence of Pop III stars was the apparent gap between the chemical composition of the universe as predicted by BBN, mainly light elements up to Lithium, and the presence of metals needed for the formation of Population II stars. Pop III stars bridge this gap by synthesizing the first metals in the universe and spreading them through the interstellar and intergalactic medium, where they are later instrumental for the formation of Pop II stars. For this scenario all Pop III stars or a large fraction of them must finish their evolution by SN explosions, and therefore these stars should be very massive, heavier than $50M_\odot$.

Pop III stars can also play an important role in tackling one of the unsolved fundamental problems in astrophysics, namely the origin of magnetic fields in the universe. Bisnovatyi-Kogan, Ruzmaikin, and Sunyaev first proposed in 1973 [18] that the earliest stars could be a site for magnetic field generation. The basic idea was that the field could be amplified in massive stars due the dynamo effect and then expelled into the circumstellar medium through stellar winds. The authors realized that the weak seeds necessary for the mechanism to be effective could be produced through the Biermann battery process [19], an idea that is currently used in many scenarios for the origin of cosmological magnetic fields. The authors acknowledge that the first-star scenario they proposed is similar to the one put forward by Hoyle [20] for AGN. In Sec. II the creation of magnetic field is considered in more detail.

After recombination at $z_{\rm rec} \approx 1100$, the baryons in the universe are in the form of cold, neutral, atomic hydrogen and light atoms. However, we know from WMAP observations [17] that the universe must have been reionized at a later epoch. In the simple assumption of a burst-like ionization, the redshift of this event that best fits the WMAP data is $z_{\rm reion} = 11.0 \pm 1.4$. More realistic scenarios including continuous and multiburst reionization have been developed in the literature (e.g. [21,22]). The reionization occurred due to photons produced by astrophysical sources. The most plausible sources of such photons are indeed hot Pop III stars, which produce ionizing photons either directly [23] or as a result of Pop III supernova explosions [24]. However, some other sources may also be required. For instance for reionization of remote regions and voids, photons with larger path-length are needed. In Ref. [25] X-ray photons from binaries with a black hole as accreting component have been proposed.

On the other hand, severe constraints on the sources of reionization have been obtained in Ref. [26].

Pop III stars arise from baryons gravitationally pulled into potential wells formed by dark matter (DM) dominated mini-halos. These mini-halos collapse at $z \sim 20$–30 and eventually capture baryons which can form stellar objects. Numerical simulations of this collapse show no fragmentation of baryonic matter and the first stars form with large masses, spread over a wide range, with a typical value $M \sim 100M_\odot$ (see [27] and references therein).

At masses $140M_\odot \le M_* \le 260M_\odot$, Pop III stars are fully disrupted in SN explosion due to e^+e^- pair-instability. These SNe are called Pair-Instability SNe (PISN). The energy output of PISN in the form of ejecta is in the range 10^{51}–10^{53} erg, and could be much higher than for Pop II and Pop I SNe. In the mass range $M < 140M_\odot$ and $M > 260M_\odot$, a massive black hole and disc are produced, with ejection of gas in the form of jets. These stars are good candidates for GRBs, and the production of UHE neutrinos from Pop III GRBs has been recently studied in [28]. For further information on Pop III stars and SNe see the reviews in [29].

In the following we will refer to this epoch of Pop III star formation as the *bright phase*, an expression often used also to refer to a period of enhanced quasar activity (about $\sim$20 quasars are observed at $5.7 \lesssim z \lesssim 6.4$ [30]).

The paper is organized as follows: in Sec. II we discuss some subtle issues related to the formation of shocks in Pop III supernova explosions and in the particle acceleration at such shocks. In Sec. III the basic features of the bright-phase model are described. In Sec. IV we present the calculations of the neutrino fluxes and the constraints on such fluxes. The conclusions are presented in Sec. V.

II. MAGNETIC FIELD, SHOCK FORMATION AND PARTICLE ACCELERATION

The environment in which Pop III stars explode as SNe is quite unlike that of SNe in the present universe and one may wonder whether particle acceleration may take place at all. One may argue that the region where the Pop III star explode as a SN may be polluted by the magnetic field expelled through the pre-supernova wind [18] possibly scaled as $\sim 1/r$ from the star's surface (or from the

Alfvenic surface if this is outside the star's surface). The magnetic field at the star's surface may be produced through the dynamo mechanism [31,32]. At large distances from the star, where particle acceleration occurs the field is therefore expected to be much weaker than at the star, but still strong enough to be a useful seed for processes to occur at later times, when the SN shock crosses those regions. One may also think of situations in which the shock propagates in very weakly magnetized regions, appropriate to the intergalactic medium at redshift $z \sim 10$–20. This leads to two questions: (1) does a shock develop in the supernova explosions, since such shocks are collisionless, namely mediated by electromagnetic instabilities? (2) if the shock forms, is there enough magnetic turbulence to lead to diffusive particle acceleration up to energies $\sim 10^{19}$–10^{20} eV, needed for neutrino production?

Numerical simulations of the development of the ion Weibel instability [33] in the absence of a preexisting magnetic field show that a shock front does in fact form [34]. The simulations are carried out for relativistic electron-ion plasmas, while the motion of the plasma ejected from a Pop III star explosion might be Newtonian or relativistic depending upon details of the explosion, but the general physical principles that lead to the formation of the shock front should be left unchanged, so we may be confident that a shock wave does indeed form even in an unmagnetized medium such as the one that might host Pop III stars. The fate of the magnetic field behind the shock is not well established: the magnetic filaments which are formed through the Weibel instability might merge to form larger scale magnetic fields, but it is not clear whether this happens before the field may be damped.

More important than the presence of preexisting magnetic field is the reionization of the surrounding gas prior to the SN explosion: a collisionless shock would hardly develop in a tenuous neutral medium, such as in the universe soon after recombination. Pop III stars can naturally fulfill the ionization requirements since, being very hot, they are expected to be responsible for the production of intense UV radiation that causes reionization (see e.g.[23] and references therein).

The issue of particle acceleration is more delicate in that diffusive acceleration at a shock front requires the existence of a turbulent magnetic field upstream and downstream of the shock. The only way that a magnetic field can be generated upstream of the shock is through streaming instability, an intrinsically nonlinear process (particle acceleration occurs if accelerated particles are able to excite the instability but the instability is excited by accelerated particles). This process is also necessary in SNe in the ISM, and in fact the instability must be driven to its nonlinear regime if to reach cosmic ray energies as high as the knee [35]. It has however been proposed that in the presence of quasi-perpendicular shocks, the acceleration time may be shorter and lead to higher energies of accelerated particles [36]. This mechanism may be at work in a fraction of supernova remnants: in particular, when the shock propagates in the wind of the pre-supernova star, the geometry is likely to be best suited for perpendicular shock drift acceleration. In Ref. [37] the authors estimated that such stars may accelerate CRs up to the knee and speculate that even higher energies may be reached in binary systems. The assessment of the relative importance of particle acceleration at parallel shocks with magnetic field amplification versus perpendicular shocks is still a subject of much investigation.

The cosmic ray driven streaming instability can proceed in a resonant [38] or nonresonant [39] way, but these are just two sides of the same physical phenomenon [40]. The nonresonant branch grows faster for higher shock velocities [40].

Once the instability turns nonlinear it is very difficult to predict its saturation, which might be dominated by intrinsic dynamical scales (such as the advection time of a fluid element through the shock) or by damping. The issue of damping might be of particular importance in the clouds where Pop III stars originate since the neutral fraction might be appreciable and ion-neutral damping could well shut the instability off. In the absence of damping and assuming a naive extrapolation of quasi-linear theory to the nonlinear regime one can estimate the saturation levels of the instability as

$$\frac{\delta B_{\rm res}^2}{8\pi} = \frac{1}{M_A} \rho V_s^2 \xi_{\rm CR} \tag{1}$$

in the resonant case, and

$$\frac{\delta B_{\rm nr}^2}{4\pi} = \frac{1}{2} \rho V_s^2 \frac{V_s}{c} \xi_{\rm CR} \tag{2}$$

in the nonresonant case. Here $M_A = V_s/v_a$ is the Alfvenic Mach number, and $v_a = B/(4\pi\rho)^{1/2}$ is the Alfven speed. Clearly neither the Alfvenic Mach number nor the Alfven speed are well defined in the nonlinear regime, but this ambiguity is the price to pay to write down a simple estimate of the magnetic field generated through streaming of accelerated particles upstream of the shock. In all these formulae we assumed that the shock is nonrelativistic. Our estimate for the amplified magnetic field is

$$\delta B_{\rm res} \approx 96 B_\mu^{1/2} n_1^{1/4} V_9^{1/2} \xi_{\rm CR}^{1/2} \mu G \tag{3}$$

in the resonant case, and

$$\delta B_{\rm nr} \approx 590 n_1^{1/2} V_9^{3/2} \xi_{\rm CR}^{1/2} \mu G \tag{4}$$

in the nonresonant case. Here $V_9 = V_s/10^9$ cm/s and $n_1 = n/1\ {\rm cm}^{-3}$. $\xi_{\rm CR}$ is the fraction of the inflow ram pressure $m_p n V_s^2$ that gets converted to accelerated particles.

Assuming Bohm diffusion, the acceleration time in the two cases reads

$$\tau_{\rm res}(E) \approx 3.5 \times 10^{-7} E({\rm eV}) B_\mu^{-1/2} n_1^{-1/4} V_9^{-5/2} \xi_{\rm CR}^{-1/2}\ {\rm s} \quad (5)$$

in the resonant case, and

$$\tau_{\rm nr}(E) \approx 5 \times 10^{-8} E({\rm eV}) V_9^{-7/2} n_1^{-1/2} \xi_{\rm CR}^{-1/2}\ {\rm s} \quad (6)$$

in the nonresonant case. A more careful analysis of the saturation of the instability in the resonant and nonresonant cases and a discussion of the Bohm diffusion regime were presented in [41,42].

It is important to keep in mind that in the nonresonant case the self-generated turbulence is on spatial scales that are much smaller than the gyration radius of the particles, therefore these modes are not very effective in scattering the particles, despite the large growth rate. In this sense, the assumption of Bohm diffusion for nonresonant modes is poorly justified, and would require an efficient cascade of the modes to much larger scales, an inverse cascade.

There are at least three different ways to define the maximum energy of accelerated particles: (1) Time limitation: $\tau(E_{\rm max}^{\rm acc}) \sim T_{\rm ST}$, where $T_{\rm ST}$ is the time when the remnant enters the Sedov-Taylor phase; (2) Space limitation of the precursor: $D(E_{\rm max}^{\rm acc})/V_s \sim \eta R_{\rm ST}$, where $R_{\rm ST}$ is the radius of the supernova remnant at the beginning of the Sedov phase, and $\eta \lesssim 1$; (3) Space limitation on the gyration radius: $r_L(E_{\rm max}^{\rm acc}) \sim R_{\rm ST}$. The three criteria lead to different estimates for the value of $E_{\rm max}^{\rm acc}$.

The Sedov time and radius are related to the ejected mass and to the total energetics through:

$$T_{\rm ST} = 70 \frac{M_{\rm ej,\odot}^{5/6}}{\epsilon_{51}^{1/2} n_1^{1/3}}\ {\rm years} \quad (7)$$

and

$$R_{\rm ST} = 6.6 \times 10^{18} \left(\frac{M_{\rm ej,\odot}}{n_1}\right)^{1/3}\ {\rm cm}, \quad (8)$$

where ϵ_{51} is the total energetics in units of 10^{51} erg and $M_{\rm ej,\odot}$ is the mass of the ejecta in units of solar masses.

Among the three criteria for $E_{\rm max}^{\rm acc}$ listed above, the last one leads to the highest estimated value

$$E_{\rm max}^{\rm acc} \approx 2 \times 10^{17} B_\mu^{1/2} n_1^{-1/12} V_9^{1/2} \xi_{CR}^{1/2} M_{\rm ej,\odot}^{1/3}\ {\rm eV} \quad (9)$$

for the resonant case, and

$$E_{\rm max}^{\rm acc} \approx 10^{18} n_1^{1/6} V_9^{3/2} \xi_{CR}^{1/2} M_{\rm ej,\odot}^{1/3}\ {\rm eV} \quad (10)$$

for the nonresonant case. One can see that maximum energies as high as $(2\text{–}5) \times 10^{19}$ eV might be reached if the values of the parameters are pushed to their extremes.

As mentioned above, the presence of perpendicular magnetic field may alleviate the requirements of extreme parameters in that the acceleration time is shorter in such a configuration, as discussed in Ref. [36].

In principle larger values of $E_{\rm max}^{\rm acc}$ can also be reached if the Pop III stars explode in a dense stellar region, so that accelerated particles may feel the repeated action of multiple shocks. Such situation occurs at high redshifts where the physical volume which hosts the fixed number N of pregalactic Pop III stars is $(1+z)^3$ times smaller than at present. Even a stronger case might be realized in mini-cluster models, when the Pop III stars are produced in the relatively small volume of a mini-cluster. In this case the spatial scale to be used to determine the maximum energy is the size of the region where the shocks emerge, provided the acceleration time remains smaller than the loss time scale as plotted in Fig. 1. In such models the acceleration time may be shortened by the presence of collective magnetic field expulsion in the form of stellar winds. Acceleration at multiple shocks has been first discussed in [38] and is known to result in a hard spectrum: in the asymptotic case of a very large number of shocks, the spectrum of accelerated particles tends to $\propto E^{-3/2}$.

Larger values of $E_{\rm max}^{\rm acc}$, up to 10^{20} eV, were inferred in previous literature for the case of acceleration at relativistic shocks with large Lorentz factors $\Gamma \gg 1$, see e.g. [43]. Usually this case is considered for jets, and, in particular, for GRBs, which are often assumed to be potential sources of particles with maximum energies $\sim 10^{20}$ eV [44,45]. These energies are the consequence of the relativistic

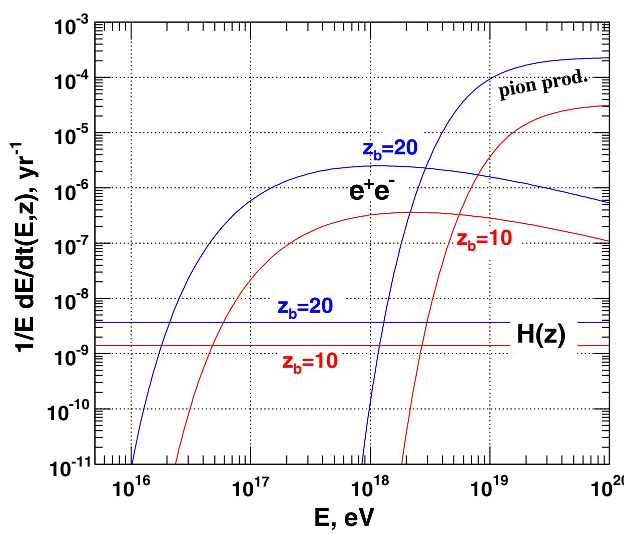

FIG. 1 (color online). Energy losses of protons at cosmological epochs $z_b = 10$ (shown in red) and $z_b = 20$ (shown in blue). Adiabatic energy losses are given by $H(z)$. Also shown are losses due to pair production on CMB, $p + \gamma_{\rm cmb} \to p + e^- + e^+$, and pion photo-production, $p + \gamma_{\rm cmb} \to N +$ pions. Photo-production responsible for the generation of neutrinos dominates at energy above the crossing of the pair-production and pion production curves ($E_c^b = 3 \times 10^{18}$ eV at $z_b = 20$ and $E_c^b = 5.5 \times 10^{18}$ eV at $z_b = 10$). The intersection of the pair-production curve with $H(z)$ (adiabatic losses) determines the energy $\varepsilon_{\rm pair}$, above which the pair-production energy losses dominate ($\varepsilon_{\rm pair} = 2.1 \times 10^{16}$ eV at $z_b = 20$ and $\varepsilon_{\rm pair} = 5 \times 10^{16}$ eV at $z_b = 10$).

motion of the outflow in jets, which in principle can shorten the acceleration time by a factor $1/\Gamma$.

Particle acceleration at relativistic shocks requires the presence of strong turbulence generated in the downstream region, in order to avoid particle trapping, typical of relativistic shocks [46]: particle acceleration is possible only for quasi-parallel shocks, namely for magnetic fields oriented within an angle $\sim 1/\Gamma$ from the normal to the shock surface. At perpendicular shocks the return probability from downstream tends to vanish thereby making the spectra steeper. In fact, even the compression of the large scales in the upstream turbulent fields at the shock surface may lead to spectra steeper than the canonical $E^{-2.23}$ [47], that is expected for parallel relativistic shocks in the regime of small pitch angle scattering. The presence of strong turbulence may alleviate this problem, so to make GRBs from Pop III stars potential sources of UHECRs with energy $E_{\rm max}^{\rm acc} \sim 10^{20}$ eV, although many physical effects may appreciably reduce the maximum energy, and a case-by-case investigation should be made.

The issue of whether particle acceleration can occur in the first stars is also of relevance for the origin of cosmic magnetic fields: in a recent paper [48] the authors proposed that if supernovae arising from the death of primeval stars accelerate CRs effectively (though not necessarily to very high energies), the instability induced by the escape of these cosmic rays into the intergalactic medium may lead to the formation of magnetic seeds that can possibly be reprocessed and amplified at later cosmic epochs.

III. THE MODEL

Here we consider a model in which the bright phase is powered by SN explosions of Pop III stars.

In order to enrich space with metals, the rate of SN explosions must be high and thus Pop III stars have to be massive, typically with $M > 50 M_\odot$. Stars in the interval of masses (140–260)$M_\odot$ undergo e^+e^- instability and end their evolution with full destruction in SN explosions. Outside this interval, and most notably at $M > 50 M_\odot$, a SN explosion leaves behind a massive black hole. In both cases a SN explosion results in the formation of a shock front where particles from the interstellar medium can be accelerated (see Sec. II for a detailed discussion). The formation of such (collisionless) shocks is expected to be made possible because of the ionization of the circumstellar region caused by photons produced by either the Pop III star itself or its supernova. The chemical composition of the accelerated particles should reflect that of the circumstellar medium, which is quite different from that in our Galaxy. The primordial gas after BBN is characterized by 75% of Hydrogen and by the absence of the heavy elements (metals). Some heavy elements pollution is expected to be caused by multiple SN explosions in the same spatial regions. On the other hand, the metallicity of the gas in Pop II star formation regions is expected to be relatively low, therefore it appears reasonable to assume that the medium around Pop III stars is dominated by light elements. We assume here that most of the accelerated particles are protons.

As discussed above, the SN explosion of Pop III stars can result in the formation of either nonrelativistic or relativistic shock fronts, the latter usually associated to the formation of jets. The spectrum of accelerated particles in the two cases is usually taken to be $\propto E^{-\gamma_g}$ with $\gamma_g = 2.0$ for nonrelativistic shocks and $\gamma_g \approx 2.2$–2.3 for relativistic shocks, although numerous physical effects can change this naive expectation. In both cases we assume the maximum energy of accelerated particles to be $E_{\rm max}^{\rm acc} \sim 10^{20}$ eV and the minimum energy to be as low as $E_{\rm min} \sim m_p c^2 \sim 10^9$ eV. Increasing $E_{\rm min}$ results in larger neutrino flux in the calculations below.

As far as $E_{\rm max}$ is concerned, the energy 10^{20} eV may be problematic in nonrelativistic shocks, although particle reacceleration at multiple SN shocks may hopefully provide this energy (see Sec. II). The possibility of reaching $E_{\rm max} \sim 10^{20}$ eV might be more promising in the case of relativistic shocks and is often invoked especially in GRB models of UHECRs, but one should admit that this possibility is poorly understood.

Rather than achieving a quantitatively accurate prediction of the neutrino flux from Pop III stars, our aim here is to develop a simple calculation that catches the main physical ingredients and illustrates in a clear, transparent way, the relevant aspects of the problem. For this reason we consider a model in which the redshift of the bright phase is fixed and the duration of the phase is short. All physical processes involved in the bright phase and the neutrino production are much shorter than the Hubble time $H^{-1}(z_b)$ at the redshift of the burst z_b. The latter can be calculated as $H(z_b) = H_0\sqrt{\Omega_m(1+z_b)^3 + \Omega_\Lambda}$ with $H_0 = 72$ km/s Mpc, $\Omega_m = 0.27$, and $\Omega_\Lambda = 0.73$. At $z_b = 10$ the Hubble time is $H(z_b)^{-1} \approx 8 \times 10^8$ yr, to be compared with the lifetime $\sim 10^6$ yr of a Pop III star with $M \sim 100 M_\odot$ on the Main Sequence (the longest time scale for the evolution of such stars), with the time of particle acceleration, $t_{\rm acc} \sim 10^4$–10^5 yr, and with the characteristic time of photopion energy losses, $\sim 3 \times 10^5$ yr, responsible for neutrino production (see Fig. 1 for comparison of these times at various energies). These estimates justify the assumption of a bursting bright phase: all processes are assumed to occur at a fixed redshift z_b, for which further on we consider $z_b = 10$ and $z_b = 20$ as our benchmark cases (see Fig. 1). This allows us to obtain analytical expressions for the neutrino fluxes. A more realistic model of a bright phase extended in time may be obtained by integration over different z_b. However, it is interesting to note that models with two bursts of reionization can be formulated [22], as based on taking into account the important role of neutral H_2 molecules in star formation. The redshifts of the predicted bursts are $z_1 \approx 11$–15 and $z_2 \approx 6$, close to the value of z_b that we consider.

Having in mind the same goal of simplicity of formulae and transparency of the calculations, we assume power law spectra with slope $\gamma_g = 2.0$, which does not lead to dramatic differences compared with similar spectra with $\gamma_g \approx 2.2$–2.3.

IV. ASSUMPTIONS AND ESTIMATES

The burst of accelerated protons generated during the bright phase at epoch z_b produces a spectrum that can be expressed as a function of the energy density $\omega_p(z_b)$ of the same protons at that epoch as:

$$n_p(E_b, z_b) = \frac{\omega_p(z_b)}{\ln(E_b^{\max}/E_b^{\min})} E_b^{-2}. \quad (11)$$

Henceforth $\omega_p(z_b)$ will be used at the main normalization parameter of our calculations.

The cosmogenic neutrinos are produced in collisions of UHE protons with CMB photons through generation of pions, $p + \gamma_{\rm cmb} \to \pi^\pm +$ anything. On average a proton transfers 20% of its energy to the leading pion, and this energy is approximately equally divided among the four leptons (three neutrinos and one electron), so that $E_\nu = (1/20)E_p$. Photopion production dominates starting at the proton energy E_c^b, where energy losses due to photopion production exceed pair-production losses ($p + \gamma_{\rm cmb} \to p + e^- + e^+$). At $z = 0$ this condition occurs at proton energy $E_c \approx 6 \times 10^{19}$ eV, therefore at epoch z_b this energy is $E_c^b(z_b) = E_c/(1 + z_b)$ (see Fig. 1), since the energy of CMB photon is $(1 + z_b)$ times higher. Hence the characteristic neutrino energy in the observed spectrum is

$$E_\nu^c = \frac{1}{20}\frac{E_c}{(1+z_b)^2} = 7.5 \times 10^{15}\left(\frac{20}{1+z_b}\right)^2 \text{ eV}, \quad (12)$$

where an extra factor $(1 + z_b)$ appears because of the redshifted neutrino energy. Below E_ν^c, the neutrino spectrum becomes flatter, being produced by pions moving backwards in the cm-system. Above E_ν^c, the neutrino spectrum follows the parent proton spectrum, $\propto E^{-2}$. Equation (12) thus gives the neutrino energy where detectability reaches its maximum. The main contribution to the number of detected neutrinos is given by some energy interval centered at E_ν^c.

For the calculation of the neutrino spectrum it is convenient, following [1,49,50], to introduce the so-called "unmodified" proton spectrum, calculated taking into account only adiabatic energy losses due to redshift. This proton spectrum at $z = 0$ is

$$n_p^{\rm unm}(E) = \frac{\omega_p(z_b)}{(1+z_b)^4}\frac{1}{\ln(E_{\max}/E_{\min})} E^{-2}. \quad (13)$$

Equation (13) is easy to understand because $\omega_p(z_b)/(1 + z_b)^4$ is the unmodified energy density of protons at $z = 0$: $(1 + z_b)^3$ is due to the expansion of the universe and the extra factor $(1 + z_b)$ accounts for adiabatic energy losses.

The spatial density of UHE neutrinos at $z = 0$ can now be readily calculated using the unmodified proton spectrum and taking into account that each proton in one photonuclear $p\gamma$ collision produces three neutrinos and the energy of a neutrino E_ν is connected with the energy E of the "unmodified" proton as $E_\nu = (1/20)E$, independent of the redshift z when the collision occurs, because adiabatic energy losses of neutrinos and UHE protons are identical.

Thus, one obtains at $E_\nu \geq E_\nu^c$:

$$n_\nu(E_\nu)dE_\nu = \frac{3f_{\pi^\pm}}{20}\frac{\omega_p(z_b)}{(1+z_b)^4}\frac{1}{\ln(E_{\max}/E_{\min})}\frac{dE_\nu}{E_\nu^2}, \quad (14)$$

where $f_{\pi^\pm} \approx 2/3$ is the fraction of charged pions.

The maximum neutrino energy at $z = 0$ can be estimated as

$$E_\nu^{\max} = \frac{1}{20}E_{\max} = \frac{1}{20}\frac{E_b^{\max}}{1+z_b}. \quad (15)$$

The neutrino flux,

$$E_\nu^2 J_\nu(E_\nu) = 0.1\frac{c}{4\pi}\frac{\omega_p(z_b)}{(1+z_b)^4}\frac{1}{\ln(E_{\max}/E_{\min})}, \quad (16)$$

is fully determined by the value of the basic parameter $\omega_p(z_b)/(1 + z_b)^4$.

We estimate first the range of basic parameters resulting in detectable neutrino fluxes and then discuss whether these parameters meet the restrictions imposed by the *e-m* cascade and energy release by Pop III SN explosions.

The value of the basic parameter $\omega_p(z_b)/(1 + z_b)^4$ for detectable neutrino flux can be found from the future IceCube sensitivity [51] $E^2 J_{\nu_\mu}(E) = 3\times10^{-9}\,\text{GeV/cm}^2\,\text{ssr}$. (The present upper limit of IceCube is a factor of 3 higher). From this condition the basic parameter must be

$$\omega_p(z_b)/(1 + z_b)^4 \geq 9.5 \times 10^{-7} \text{ eV cm}^{-3}. \quad (17)$$

It is also necessary to estimate the maximum energy of acceleration $E_{\max}^{\rm acc} = E_b^{\max}$ discussed in Sec. II. The rigorous limit follows from the condition that pion production dominates upon pair production (see Fig. 1):

$$E_b^{\max} = E_c(z_b) = \frac{E_c}{1+z_b} = 3 \times 10^{18}\left(\frac{20}{1+z_b}\right) \text{eV}. \quad (18)$$

The same limit follows from the condition which provides an appreciable neutrino flux, $E_\nu^{\max} \geq E_\nu^c$. In the present estimates we assume that $E_b^{\max} \sim 10^{20}$ eV. The case $E_b^{\max} \gtrsim 1 \times 10^{19}$ eV needs more detailed calculation of the neutrino spectra and will be presented in a forthcoming paper.

In the following we discuss the constraints on the cosmogenic neutrino flux from the bright phase.

The production of UHE neutrinos from pion decays is accompanied by *e-m* cascade radiation, which provides us with an upper limit on the neutrino flux [6]. The neutrino flux given by Eq. (16) with energy density (17) must respect this limit as given by the recent Fermi observations [7]. In [8] this limit is expressed in terms of an upper limit on the energy density of the cascade radiation $\omega_{\rm cas}^{\rm max} \leq 5.8 \times 10^{-7}$ eV/cm^3. We can estimate the cascade energy density in our calculations from $\omega_p(z_b)$ evaluated in Eq. (17).

At energies $E_b \geq \varepsilon_{\rm pair}$ (see Fig. 1), protons lose energy on time scales shorter than the Hubble time $H^{-1}(z_b)$, producing e^+e^- pairs and pions. The energy density $\omega_{\rm cas}^b$ of the cascade radiation at epoch z_b can be written as

$$\omega_{\rm cas}^b \lesssim \int_{\varepsilon_{\rm pair}}^{E_b^{\rm max}} n_p(E_b) E_b dE_b, \tag{19}$$

and for the cascade energy density at $z = 0$ we have

$$\omega_{\rm cas} \lesssim \frac{\omega_p(z_b)}{(1+z_b)^4}\left[\ln\frac{E_b^{\rm max}}{\varepsilon_{\rm pair}}\right]\left[\ln\frac{E_b^{\rm max}}{E_b^{\rm min}}\right]^{-1} = 2.9 \times 10^{-7}\ {\rm eV/cm^3}. \tag{20}$$

One can see that the cascade energy density (20), which follows from neutrino production in the bright phase does not exceed the upper limit $\omega_{\rm cas}^{\rm max} = 5.8 \times 10^{-7}$ eV/cm^3, obtained in [8] from Fermi observations. In fact, this upper limit is derived for the cascade spectrum produced at redshift $z = 0$, while in our case the cascade is produced at large z_b and its spectrum is redshifted. This makes the upper limit higher than $\omega_{\rm cas}^{\rm max}$ from [8].

Finally, we address the issue of whether the energy density of protons in Eq. (17) which is needed to warrant detectability of the neutrino flux in IceCube can be provided by Pop III SN explosions. For this, we estimate a fraction ξ of cosmological baryon mass processed by Pop III pre-supernovae.

From the energy density of UHE protons $\omega_p(z_b) = W_{\rm SN}^p n_*(z_b)$, provided by the SN energy release in the form of UHE protons $W_{\rm SN}^p$ and from the space density of Pop III pre-supernovae n_*, one obtains

$$\xi = \frac{\omega_p(z_b)}{(1+z_b)^4}\frac{M_*(1+z_b)}{W_{\rm SN}^p \Omega_b \rho_{\rm cr}}. \tag{21}$$

Here we used the equality $M_* n_*(z_b) = \xi \rho_b(z_b)$, where $M_* \sim 100 M_\odot$ is a typical mass of Pop III pre-supernova and $\rho_b(z_b) = \Omega_b \rho_{\rm cr}(1+z_b)^3$ is the total baryonic mass density at epoch z_b, where for the baryonic gas density at $z = 0$ we use the WMAP [17] value $\Omega_b \rho_{\rm cr} = 4.27 \times 10^{-31}$ g/cm^3. For $M_* = 100 M_\odot$ the SN energy release can be conservatively estimated as $W_{\rm SN}^p \sim 10^{51}$ erg. As a result we obtain for the case of Eq. (17) the value $\xi = 1.4 \times 10^{-2}$, i.e. about 1% of the total baryonic mass processed by Pop III pre-supernovae.

In such a simplified model with a fixed z_b it is not easy to discuss how well the obtained value of ξ fits the observed reionization (see [52–54]): in the absence of strong feedback from the first stars that form in the parent cloud, the values of ξ can be relatively large, but it can be appreciably reduced due to the effect of radiative feedback, that suppresses the efficiency of formation of new stars after the few first massive stars.

If to assume that each SN explosion leaves behind a black hole with mass $M_{\rm bh} \sim 3 M_\odot$, the fraction of baryonic mass in the form of these black holes is only $\eta = \xi M_{\rm bh}/M_* \sim 4 \times 10^{-4}$.

The full-scale model describing production and cooling of the first gas clouds, evolution of the first stars with subsequent SN explosions and consequent reionization are beyond our simplified model of burst-like bright phase. As a next step of our investigation (currently in progress) we plan to include the possibility of a bright phase extended in redshift, a detailed numerical calculation of the neutrino fluxes and a more careful study of the reionization induced by SN explosions.

A. $\bar{\nu}_e + e^- \to W^- \to$ hadrons resonance.

The resonant production of hadrons by UHE $\bar{\nu}_e$-neutrinos is a remarkable feature and signature of the Pop III bright phase model.

The resonant production of W-bosons was first proposed by S. L. Glashow in 1960 [55] for the generation of high-energy muons in the reaction $\bar{\nu}_e + e^- \to W^- \to \mu^- + \bar{\nu}_\mu$. In 1977 Berezinsky and Gazizov [56] suggested another channel of the Glashow resonance $\bar{\nu}_e + e^- \to W^- \to$ hadrons for the detection of UHE neutrinos. The resonant neutrino energy is given by

$$E_0 = m_W^2/(2m_e) = 6.3 \times 10^6\ {\rm GeV}, \tag{22}$$

and the rate of resonant events $\nu_{\rm res}$ in a detector with total number of electrons N_e is determined by the exact formula derived in [56]:

$$\nu_{\rm res} = 2\pi N_e \sigma_{\rm eff} J_{\bar{\nu}_e}(E_0) E_0, \tag{23}$$

where $J_{\bar{\nu}_e}(E_0)$ is the diffuse flux at the resonant energy E_0, $N_e = (5/9)M/m_H$ is the number of electrons in the underground detector with mass M, 2π is the solid angle at which a deep-underground detector is open for resonant neutrinos (for the detailed calculations see [57]), $\sigma_{\rm eff} = (3\pi/\sqrt{2})G_F = 3.0 \times 10^{-32}$ cm^2 is the effective cross section obtained by integration of the Breit-Wigner cross section over energy.

The flux of $\bar{\nu}_e$-neutrinos produced in $p\gamma_{\rm cmb}$ collisions is strongly suppressed and appears mostly due to oscillations as has been already anticipated in [56]. According to recent calculations, $J_{\bar{\nu}_e}(E)$ is approximately the same for all flavors, namely it is about $1/6$ of the all-flavor neutrino flux.

For the application of the Glashow resonance to the Pop III bright phase, we have to first determine z_b. It can be found from the condition that the energy of the parent proton for the resonant neutrino at z_b must be above the critical energy $E_c/(1+z_b)$, at which the energy losses for pion and e^+e^--pair production are equal. Using $E_c = 6\times 10^{19}$ eV one has

$$1 + z_b = \left(\frac{E_c}{20E_0}\right)^{1/2} = 21.8, \tag{24}$$

which coincides with our assumption $z_b = 20$.

The attractive feature of this model is that it imposes milder constraints on $E_{\rm max}^{\rm acc}$. Indeed, the energy of the parent proton for resonant neutrino at z_b is $E_p(z_b) = 20E_0(1+z_b) = 2.8\times 10^{18}$ eV, which is considerably lower than $E_{\rm max}^{\rm acc}$ discussed in Sec. II. Below we assume $E_{\rm max}^{\rm acc} = 1\times 10^{19}$ eV.

We calculate now the rate of the resonant events in IceCube with effective mass $M \sim 1\times 10^9$ t. According to Eq. (23), there is only one unknown quantity for the calculations, the flux $J_{\bar\nu_e}(E_0)$, which we take as maximal, namely corresponding to the cascade upper limit $\omega_{\rm cas}^{\rm max} = 5.8\times 10^{-7}$ eV/cm^3. Using Eq. (16) for the neutrino flux with $\omega_p(z_b)/(1+z_b)^4$ from Eq. (20), and assuming equipartition of neutrino flavors with a fraction $1/6$ for $\bar\nu_e$, we obtain the rate of the resonant events in IceCube as

$$\nu_{\rm res} = \frac{0.1c}{12}\,\frac{N_e\sigma_{\rm eff}}{\ln(E_b^{\rm max}/\varepsilon_{\rm pair})}\,\frac{\omega_{\rm cas}^{\rm max}}{E_0} \tag{25}$$

In more accurate calculations, where one takes into account that decaying pions transfer only half of their energy into *e-m* cascade, $E_b^{\rm max}$ in Eq. (25) should be substituted by $\sqrt{E_c^b E_b^{\rm max}}$, and the rate of events reaches $\nu_{\rm res} = 1.3\ {\rm yr}^{-1}$. In fact, this rate is not the maximal one, because as already mentioned above, $\omega_{\rm cas}^{\rm max} = 5.8\times 10^{-7}$ eV/cm^3 was obtained for $z = 0$, while for $z_b = 20$ this limit is higher. Most probably each resonant event with tremendous energy release 6.3×10^{15} eV in the form of nuclear and electromagnetic cascades will be detected in IceCube. Ten well-identified events with equal energies during five years of observations is a good enough signature of Pop III bright phase model with $z_b \geq 20$.

V. DISCUSSION AND CONCLUSIONS

There are theoretical and observational indications that the first star formation, the galaxy formation and AGN went through a phase of enhanced activity at high redshifts, $z \sim 10$–20 (e.g. low metallicity stars and the observation of about 20 quasars at $z \sim 5$–6.4 [30]).

In this paper we concentrated upon a bright phase in the stellar evolution, in the form of Pop III stars. These astrophysical objects play an important role in at least three ways: (1) they enrich the universe with the metals that appear to be necessary for the formation of Pop II stars, thereby bridging the gap between the chemical composition of the universe as predicted by BBN and the one observed in today's universe, (2) Pop III stars may reionize the universe, as observed by WMAP [17,21], and (3) the universe may be polluted with magnetic fields from Pop III stars [34,48].

Both the metal enrichment and the reionization require that Pop III stars are short-lived and therefore very massive stars, typically with $M_* \gtrsim 100M_\odot$, which explode as supernovae. In the mass range $140M_\odot \gtrsim 260M_\odot$ the star is fully disrupted due to the e^+e^- pair-instability. In the mass range outside this interval a massive black hole and disc are produced, with ejection of gas in the form of jets. In both cases a SN explosion results in a shock where particle acceleration may take place. In the low-mass limit $M \sim (15$–$40)M_\odot$ the energy release is $W_{\rm SN} \approx 2\times 10^{51}$ erg [58], and at large masses $M > 100M_\odot$ it may exceed 10^{53} erg. In case of a very massive black hole the energy output can be much higher due to the Blandford-Znajek effect [59] (see also [60]).

Here we concentrate on the possibility that diffusive particle acceleration may take place at the shocks that are generated when Pop III stars explode. The medium in which these SN explosions occur is expected to be made mainly of hydrogen (75%), with a small contamination of helium and other light elements. Although a metal contamination must appear and increase once the bright phase starts, we expect that this represents a weak effect. Therefore, we considered protons as the main component of accelerated particles.

We discussed at length three delicate aspects of the problem of acceleration: (1) the formation of collisionless shocks around primordial stars; (2) the generation of magnetic field; (3) particle scattering and acceleration.

Pop III stars are the most natural sources of reionization of the universe due to their high temperature and luminosity [23]. The reionization of the universe is a crucial ingredient of our discussion, since in the absence of such a phenomenon the collisionless shocks associated with SN explosions would hardly form or would have very peculiar characteristics. These shocks are formed because of the mediation of electromagnetic instabilities such as the Weibel instability that was found to be effective even in the absence of a preexisting magnetic field. The generation of a magnetic field is just a different aspect of the problem of shock formation, in that the presence of a magnetic field is the very reason why particle motion is slowed down, thereby fulfilling the Rankine-Hugoniot conditions at the shock surface. Magnetic fields can also be formed downstream of the shock if the upstream plasma contains density inhomogeneities that induce shock corrugation and eventually lead to eddies that may considerably amplify a magnetic seed [61]. For relativistic shocks, amplification of the magnetic field by a macroscopic turbulent dynamo triggered by the Kelvin-Helmholtz shear instability has been investigated in [62].

It is worth stressing that in these cases the magnetic field is formed behind the shock surface and advected downstream. It is therefore not useful in terms of scattering the particles upstream of the shock, a necessary condition for particle acceleration. Magnetic fields can be produced upstream of the shock through streaming instability excited by accelerated particles. This instability can be excited in a resonant or nonresonant way, the former being in general slower for fast shocks but more effective in scattering accelerated particles. We found that these processes can hardly allow particles to reach energies in excess of 10^{19} eV for fast but nonrelativistic shocks, although higher $E_{\max}$ can be possibly obtained if numerous SNe explode in the same region.

Higher maximum energies, in excess of 10^{20} eV have been widely discussed and claimed to be achievable in the case of relativistic motion of the SN ejecta [43,45] (see however discussion in Sec. II).

UHE neutrinos in our model are produced in $p\gamma_{\rm cmb}$ collisions. The energy of CMB photons is $(1 + z_b)$ times higher than at $z = 0$, and the density of these photons is $(1 + z_b)^3$ times larger. We assume that a burst of UHE proton generation occurs at redshift z_b and consider two values $z_b = 10$ and $z_b = 20$ as benchmark cases. The generation spectrum is assumed to be $\propto E^{-2}$ with $E^{\rm acc}_{\max} = 1 \times 10^{20}$ eV or somewhat less. With these assumptions, the neutrino flux, given by Eq. (16), is fully determined by the basic parameter represented by the energy density $\omega_p(z_b)$ of UHE protons at epoch z_b, recalculated at the present epoch $\omega_p = \omega_p(z_b)/(1 + z_b)^4$ [see Eq. (16)]. In order to have the muon-neutrino flux detectable by IceCube, this parameter must be $\omega_p \gtrsim$ eV/cm^3. The dominant number of detected events is confined within a limited energy interval centered at energy $E^c_\nu = 7.5 \times 10^{15}(20/z_b)^2$ eV [Eq. (12)]. At this energy the detected flux has a weak maximum. The predicted flux with the value of ω_p given above is detectable by IceCube and respects the cascade upper limit on cosmogenic neutrinos and SN energetics for Pop III stars.

A unique signature of the Pop III burst model may be provided by $\bar\nu_e$ neutrinos interacting in a detector through the resonant reaction $\bar\nu_e + e^- \to W^- \to$ hadrons. The resonant energy of the neutrino is given by Eq. (22) as $E_0 = 6.3 \times 10^{15}$ eV, and the rate of neutrino events in a detector is determined by one unknown quantity, the neutrino flux at energy E_0, $J_{\bar\nu_e}(E_0)$, see Eq. (23).

We take this flux as to saturate the cascade upper limit corresponding to the cascade energy density $\omega^{\max}_{cas} = 5.8 \times 10^{-7}$ eV/cm3, though in the model with fixed z_b this limit is expected to be higher. The redshift of the burst z_b, or more precisely its lower limit, can be found from the condition that the energy $E_p(z_b)$ of a proton, parent of the resonant neutrino, must have energy higher than $E_c(z_b)$, at which pion energy losses at epochz_b exceed those for e^+e^- production. This condition determines the burst redshift as $z_b = 20.8$, see Eq. (24). The resonant events can be observed in IceCube through the Cerenkov light from nuclear and e-m cascades with energy E_0. In IceCube the predicted frequency of resonant events for the flux at the cascade upper limit is rather low, about 10 events in 5 years, but with a tremendous energy deposit, the same for all events. This may provide us with a reasonable signature of the model discussed here. These events are accompanied by muons with energies of order $0.5E^c_\nu \times 10^{15}$ eV.

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