Cumulative Neutrino and Gamma-Ray Backgrounds from Halo and Galaxy Mergers

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Abstract

The merger of dark matter halos and the gaseous structures embedded in them, such as protogalaxies, galaxies, and groups and clusters of galaxies, results in strong shocks that are capable of accelerating cosmic rays (CRs) to \( \gtrsim 10 \) PeV. These shocks will produce high-energy neutrinos and \( \gamma \)-rays through inelastic pp collisions. In this work, we study the contributions of these halo mergers to the diffuse neutrino flux and to the nonblazar portion of the extragalactic \( \gamma \)-ray background. We formulate the redshift dependence of the shock velocity, galactic radius, halo gas content, and galactic/intergalactic magnetic fields over the dark matter halo distribution up to a redshift \( z = 10 \). We find that high-redshift mergers contribute a significant amount of the CR luminosity density, and the resulting neutrino spectra could explain a large part of the observed diffuse neutrino flux above 0.1 PeV up to several PeV. We also show that our model can somewhat alleviate tensions with the extragalactic \( \gamma \)-ray background. First, since a larger fraction of the CR luminosity density comes from high redshifts, the accompanying \( \gamma \)-rays are more strongly suppressed through \( \gamma\gamma \) annihilations with the cosmic microwave background and the extragalactic background light. Second, mildly radiative-cooled shocks may lead to a harder CR spectrum with spectral indices of \( 1.5 \lesssim s \lesssim 2.0 \). Our study suggests that halo mergers, a fraction of which may also induce starbursts in the merged galaxies, can be promising neutrino emitters without violating the existing Fermi \( \gamma \)-ray constraints on the nonblazar component of the extragalactic \( \gamma \)-ray background.

Key words: cosmic rays – galaxies: clusters: general – galaxies: halos – gamma rays: diffuse background – neutrinos

1. Introduction

Neutrino astrophysics has made substantial progress since the IceCube Neutrino Observatory in the Antarctic (e.g., Gaisser & Halzen 2014; Halzen 2016, for reviews; hereafter IceCube) was completed. During the last half decade, scores of high-energy (HE) astrophysical neutrinos with energies between \( \sim 10 \) TeV and a few petaelectronvolts have been detected by IceCube, and the number keeps growing (Aartsen et al. 2013a, 2013b, 2014, 2015). The arrival directions of these neutrinos are compatible with an isotropic distribution even in the \( 10–100 \) TeV range, suggesting that a large part of these diffuse neutrinos come from extragalactic sources. Nonobservation of diffuse Galactic \( \gamma \)-rays from the Galactic plane and other extended regions independently suggests that the Galactic contribution (e.g., by Fermi bubbles or local supernova remnants) is unlikely to be dominant (Ahlers & Murase 2014; Apel et al. 2017; Abeysekara et al. 2017a, 2017b). However, despite extensive efforts, the physical nature of the sources of the diffuse neutrinos still remains in dispute. Possible candidates include gamma-ray bursts (GRBs); e.g., Waxman & Bahcall 1997; Meszaros & Waxman 2001; Murase 2008; Wang & Dai 2009; Baerwald et al. 2013; Bustamante et al. 2015; Tamborra & Ando 2016), low-power GRBs (Murase et al. 2006; Gupta & Zhang 2007; Murase & Ioka 2013; Xiao & Dai 2014, 2015; Senno et al. 2016; Denton & Tamborra 2017), radio-loud active galactic nuclei (AGNs); e.g., Mannheim 1995; Halzen & Zas 1997; Anchordoqui et al. 2008; Becker Tjus et al. 2014; Dermer et al. 2014; Murase et al. 2014; Padovani et al. 2015; Petropoulou et al. 2015; Blanco & Hooper 2017), radio-quiet/low-luminosity AGNs (Stecker et al. 1991; Alvarez-Muñiz & Mészáros 2004; Stecker 2013; Kimura et al. 2015), and AGNs embedded in galaxy clusters and groups.3 It is generally accepted that the bulk of astrophysical neutrinos are generated by charged pion \( (\pi^\pm) \) decays, and that these pions are the secondaries from cosmic-ray (CR) particles undergoing hadronuclear \( (pp) \) or photohadronic \( (p\gamma) \) interactions between the CRs and ambient target gas nuclei or photons. Meanwhile, these collisions also lead unavoidably to neutral pions \( (\pi^0) \) as well, which subsequently decay into a pair of \( \gamma \)-rays. Hence, the diffuse neutrino flux is expected to have an intimate connection with the diffuse CR and \( \gamma \)-ray backgrounds, and multimessenger analyses need to be applied to constrain the origin of these diffuse, high-energy cosmic particle fluxes (Murase et al. 2013, 2016; Bechtol et al. 2017).

Galaxy clusters and groups have been considered as promising candidate sources of IceCube’s neutrinos, and CR accelerators can be not only AGNs but also intragalactic sources, accretion shocks, and mergers of clusters and groups (e.g., Murase et al. 2008; Fang & Murase 2018). Star-forming and starburst galaxies (SFGs and SBGs, respectively) have also been suggested as promising candidates for HE neutrino sources (e.g., Loeb & Waxman 2006; Murase et al. 2013; Anchordoqui et al. 2014; Chang & Wang 2014; Tamborra et al. 2014; Chakraborty & Izaguirre 2015; Chang et al. 2015; Senno et al. 2015). In particular, starburst galaxies have dense gaseous environments and have been of interest as efficient CR reservoirs. Previous studies have assumed not only supernova and hypernova remnants (SNRs and HNRs, respectively) but also galaxy mergers, disk-driven outflows, and possible weak jets from AGNs as CR accelerators embedded in the star-forming galaxies (Murase et al. 2013; Kashiya &...
This paper is organized as follows. In Section 2, we introduce the halo mass function and the halo distribution that are used in the following sections. The merger rate and the CR energy input rate are given in Section 3. In Section 4 we discuss the redshift dependence of the CR maximum energies and the neutrino product efficiency, and we calculate the resulting neutrino and $\gamma$-ray fluxes. The results and implications are discussed in Section 6. Throughout, we assume a standard flat $\Lambda$CDM universe with present-day density parameter $\Omega_{m,0} = 0.3$ and Hubble parameter $H_0 = 71.9$ km s$^{-1}$ Mpc$^{-1}$ (Bonvin et al. 2017).

### 2. Halo Mass Function

Using the formalism of Press & Schechter (1974), the halo mass function, the number of dark matter halos per unit comoving volume contained within the logarithmic mass interval $d \ln M$, is given by

$$dN = \frac{df_c}{M} d \ln \sigma_M^{-1},$$

with the background matter density $\rho = \Omega_{m,0} \rho_c$, where $\rho_c = (3H_0^2/8\pi G)$ is the critical density, and the variance $\sigma_M$ of the linear density contrast $\delta \equiv (\Delta \rho/\rho_c)$ is smoothed over the scale $R = (3M/4\pi\rho)^{1/3}$:

$$\sigma^2_M = \int \frac{k^2 dk}{2\pi^2} P(k)|\tilde{W}(kR)|^2.$$

Here, $P(k)$ is the linear matter power spectrum we calculate following Lewis et al. (2000), and $\tilde{W}(kR) = \delta_i(kR)/kR$ for a top-hat filtering function. The significance $\nu = \delta_i/\sigma_M$ is related to the linear critical density $\delta_c$ above which virialized halos can form. In the spherical collapse model, for example, a spherical region of radius $R$ collapses and virializes at redshift $z$ when the smoothed linear overdensity $\delta_K(r, z)$ exceeds $\delta_c,0 \approx 1.686$. In the flat $\Lambda$CDM universe, the linear growth factor, the time evolution of the linear density contrast, is given by

$$D(z) = \frac{\delta_i(z)}{\delta_c,0} \approx 5 \frac{\Omega_{m,0} \sqrt{\Omega_{m,0}(1+z)^3 + 1 - \Omega_{m,0}}}{\sqrt{1 + \Omega_{m,0}(1+z)^3 + 1 - \Omega_{m,0}^0}^{3/2}} dz'.$$

We normalize $D(z)$ to be unity at $z = 0$.

For the multiplicity function $f(\nu)$, we adapt the Sheth–Tormen (Sheth & Tormen 1999) fitting formula, expressed in the form

$$f_{S-T}(\nu) = A \sqrt{\frac{2\nu}{\pi}} \left[1 + (\nu^2 a)^{-p}\right] \nu \exp \left[ -\frac{a \nu^2}{2} \right],$$

with parameters $A = 0.3222$, $a = 0.707$, and $p = 0.3$, which provide the best fit to numerical N-body simulations (Jenkins et al. 2001; Reed et al. 2006; Murray et al. 2013). We show the resulting mass function $dN/d \ln M$ for redshifts between $z = 0$ and $z = 5$ in Figure 1. Our mass function slightly

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4 In some references, e.g., in Sheth & Tormen (1999), $\nu = (\delta_c/\sigma_M)^2$ is used instead of our definition here.
underestimates that from the $N$-body simulations at high masses $\gtrsim 10^{13} M_\odot$, but it does not affect the main results of this paper. We shall use the mass functions in the following sections to estimate the redshift evolution of galactic radius, gas density, and shock velocity.

3. Merger Rate and Cosmic-ray Luminosity Density

In this section, we calculate the CR input rate due to galaxy and halo mergers by using the halo mass function we have obtained in Section 2, and we estimate the energy converted into CRs from shocks in the gas component of the merging halos as follows.

There are three timescales characterizing the CR acceleration due to galactic halo mergers: the age of the universe $t_\text{age}$; the halo merger time $t_\text{merger}$, which corresponds to the average time required to undergo one merger; and the CR injection time (the shock-crossing time) $t_\text{dyn}$. These are, for a merger that happens at redshift $z$, given by

$$t_\text{age} = \int_{\zeta}^{\infty} \left| \frac{dz'}{dz} \right| dz',$$

$$t_\text{merger} = \left[ \int d\zeta \left( \frac{dN_m}{dz d\zeta} \right) \frac{dz}{dt} \right]^{-1},$$

$$t_\text{dyn} = \lambda R_g(z) \frac{v_\text{c}(z)}{v(z)}.$$  

Here, $|dt/dz| = 1/[(1 + z)H(z)]$, $dN_m/dzd\zeta$ (given by Fakhouri et al. 2010; Fialkov & Loeb 2017) is the dimensionless merger rate per redshift interval $dz$ and per unit halo mass ratio $\zeta$, $R_g(z)$ is the mean galaxy radius, and $\lambda \sim 1$ parameterizes the orientation and geometrical uncertainty of the galaxy merger. With these timescales, the probability that a halo with mass $M$ experiences a merger within the age of the universe is given by $P(M, z) = \exp(-t_\text{merger}/t_\text{age})$. Hence, assuming that the CRs are mainly protons, the comoving CR energy input rate per the logarithm of the CR energy $\varepsilon_p$ is

$$\varepsilon_p Q_{s,\varepsilon}(z) = \frac{E_\text{merger}}{t_\text{age} C} = \int_{M_{\text{min}}^{\text{max}}} ^{M} dM \times \int \frac{1}{2} \xi_g(M, z) M_\odot \left( \frac{dN}{dM \, P(M, z)} \right) dM \frac{dN}{dM \, t_\text{age}}.$$  

(6)

where $\xi_g(M, z) = M_{\text{gas}}/M_h$ is the mass fraction in gas form, $\varepsilon_p$ is the CR energy fraction (nominally taken as 0.1), and $C = \ln(\varepsilon_p^{\max}/\varepsilon_p^{\min})$ is the normalization factor for a standard flat CR spectrum $N(\varepsilon_p) \propto \varepsilon_p^{-2}$. For $z \sim 1$, the typical maximum energy, $\varepsilon_p^{\max}$, is $\sim 10^{17}$ eV, and $C \approx 18.4$ (Kashiyama & Mészáros 2014). However, $\varepsilon_p^{\max}$ varies with redshift, as we discuss in the next section.

3.1. Gas-mass Fraction $\xi_g(M, z)$

The gas-mass fraction $\xi_g$ of dark matter halos depends on the star formation rate (SFR) and on the stellar mass $M_\star = \xi_g(M, z) M_h$. Here, we obtain $\xi_g = M_\star / M_h$ from the $M_h(M_\star)$ function inferred from observations by Behroozi et al. (2013).5 We also use the gas fraction in a normal galaxy $f_\text{g} = M_\text{gas}/(M_\text{gas} + M_\text{*})$ measured in Sargent et al. (2014). Combining the two observational results, we have constructed the redshift evolution of the gas-mass fraction in dark matter halos. That is, the gas-mass fraction $\xi_g$ is related to $f_\text{g}$ through $\xi_g^{\text{evo}} = M_\text{gas}/M_h = \xi_g f_\text{g}/(1 - f_\text{g})$, and using Equation (26) in Sargent et al. (2014), we find that

$$\xi_g^{\text{evo}} = \frac{f_\text{g}}{1 - f_\text{g}} = \frac{K}{M_h^{\beta - \gamma}} \text{sSFR}^\gamma.$$  

(7)

where $K = 10^{\alpha_{\text{sSFR}}}$ is a constant and the quantity sSFR (specific star formation rate) is the SFR per unit galaxy stellar mass. For the gas fraction in normal galaxies, we use the parameters $(\alpha_{\text{sSFR}}$, $\beta) = (9.22 \pm 0.02$, $0.81 \pm 0.03)$, together with the expression for sSFR given in the appendix of Sargent et al. (2014). In Figure 2 (see red curves), we show the redshift evolution of the mean gas-mass fraction,

$$\langle \xi_g^{\text{evo}} \rangle = \frac{\int \xi_g^{\text{evo}} \frac{dN}{dM} dM}{\int \frac{dN}{dM} dM},$$  

(8)

as well as the constant gas fraction, $\xi_g = 0.05$.

In our calculation, we take the lower and upper limits of the integration in Equation (6) as $M_{\text{min}} = 10^{10} M_\odot$ and $M_{\text{max}} = 10^{15} M_\odot$, respectively. There are two main reasons to choose the lower bound $10^{10} M_\odot$. First, considering the applicability of the $M_h(M_\star)$ relation from Behroozi et al. (2013) and the gas fraction function (Equation (7)), it is safe to truncate the halo mass at $M_h \sim 10^{10} M_\odot$. Typically, dwarf galaxies reside in halos with mass less than $10^{10} M_\odot$, and we only have the constraints from observations at $z \simeq 0$. In our model, we consider the contribution from galaxy mergers up to the redshift $z = 10$ where the $M_h(M_\star)$ function is not well tested for the lower halo masses. Also, the gas fraction function

5 In Behroozi et al. (2013), the $M_h$–$M_\star$ relation from $z = 0$ to 8 is parameterized by Equation (3). Here, we extend the domain of that function to $z = 10$ considering that the uncertainty from high-redshift contributions is small.

6 In Sargent et al. (2014), the gas fraction is written as $f_{\text{mol}}$ instead.
(Equation (7)) is modeled from (normal) star-forming galaxies (Sargent et al. 2014) and may not be valid for dwarf galaxies. Second, we estimate the low-mass halo contribution to CR luminosity density by extending the lower bound to 10^8 M_\odot*, and we found that the contribution from 10^8 to 10^{10} M_\odot halos is \leq 10% of the total luminosity density at low redshift (z \leq 3), which implies that the conclusion of this paper does not depend sensitively on the mass range.

3.2. Shock Velocity v_s

In the hierarchical clustering of a large-scale structure scenario, the galactic-size halos are contained inside larger cluster-size halos. The peculiar velocities of the galactic-size halos are, therefore, on the order of the virial velocity of the cluster-sized halo. Here, we approximate the shock velocity of the galaxy merger from the pairwise velocity dispersion projected along the line of approach of two galaxies. For galaxies with a luminosity L \approx L_* (where L_* is the characteristic luminosity), Davis & Peebles (1983) showed that the two-point correlation function at r < 20h^{-1} Mpc can be approximated by a power law:

\[ \xi(r) = \left( \frac{r}{r_0} \right)^\gamma, \]

where \gamma \approx 1.7 and r_0 \approx 5h^{-1} Mpc is the correlation length, inside which galaxies are strongly correlated. Combining the hierarchical form of the three-point correlation function of galaxies (Groth & Peebles 1977) and the cosmic virial theorem derived from the Layzer–Irvine equation, the collision (or shock) velocity can be written as

\[ v_s = \sqrt{\frac{2}{\gamma}} \frac{r_0}{5h^{-1} \text{Mpc}} \left( \frac{r}{5h^{-1} \text{Mpc}} \right)^{\frac{1}{2}} \approx 0.05 \text{red solid line} \] and with \sigma_0 = 500 (blue dashed line), respectively.

\[ \sigma_0 = 500 \text{ blue solid line} \]

\[ \sigma_0 = 300 \text{ red solid line} \]

Figure 2. Redshift-dependent gas fraction \( \xi_{\text{g}}(z) \) (red solid line) compared to a constant gas fraction \( \xi_{\text{c}} = 0.05 \) (red dashed line) for a redshift-dependent shock velocity with \sigma_0 = 300 (blue solid line) and with \sigma_0 = 500 (blue dashed line), respectively.

For our calculation, it is necessary to consider the redshift dependence of the correlation function \( \xi(r) \) in the nonlinear regime. As a useful approximation, we adopt the stable clustering (SC) hypothesis (Peebles 1974; Davis & Peebles 1977), in which only the size (or separation between structures) of the clusters changes in time while the internal density structure of clusters stays intact. This leads to \( \xi(r, z) \propto (1 + z)^{-3} \) and \( r_0 \propto (1 + z)^{(3 - \gamma)/\gamma} \).

\[ \xi(r, z) \propto (1 + z)^{-3} \text{ and } r_0 \propto (1 + z)^{(3 - \gamma)/\gamma} \]

Note that we only need the evolution of the nonlinear scale \( r_0 \), which is defined by \( \xi(r_0, z) = 1 \). A more accurate treatment (Hamilton et al. 1991) describes the evolution of \( \xi(r, z) \) from the linear to the nonlinear regime, and this treatment was generalized by Peacock & Dodds (1996) and by Smith et al. (2003) using another formula for the nonlinear function. This generalized method gives \( \xi(r, z) \) in the quasi-linear regime and confirms that \( \xi(r, z) \propto (1 + z)^{-3} \) is valid in the nonlinear limit as well. Therefore, in this paper, we use \( r_0^{\text{SC}} \propto (1 + z)^{(3 - \gamma)/\gamma} \) in Equation (10) to find the shock velocities. We show the redshift dependence of the shock velocity \( v_s(z) \) in Figure 2 (see blue curves).

Note that, in our approximation of the shock velocity (Equation (10)), the galaxy separation \( r \) given by Equation (11) is overestimated, since this latter equation takes an average of the galaxies in a cosmological volume including clusters and voids. The mean separation of the galaxies in clusters, therefore, must be smaller than the value that we have adopted here, and this overestimate of \( r \) would give a slight underestimate of \( v_s \) inside clusters. As a possible way to correct for this, we note that redshift surveys give \( \sigma_0 \approx 500 \) (Jing et al. 1998; Zehavi et al. 2002; Hawkins et al. 2008) as an average value for mergers in clusters. A qualitatively appropriate correction for the cluster shock velocity may be obtained by scaling up \( \sigma_0 \) in Equation (10) from the local average value, concluding that a realistic value of \( \sigma_0 \) for clusters is in the range 500 \( \leq \sigma_0 \leq 600 \). In terms of rates, most galaxy mergers occur in the smaller mass halos containing fewer galaxies, as opposed to large clusters. Considering the observational and theoretical uncertainties, we expect that the values of \( \sigma_0 \) lie in the range of 100 \( \leq \sigma_0 \leq 1000 \), and we take 300 \( \leq \sigma_0 \leq 500 \) as fiducial values.

4. Neutrino and \( \gamma \)-ray Production

Since in our model we need to consider the neutrino/\( \gamma \)-ray production rate up to redshift \( z = 10 \), we introduce here the redshift evolution function of the gas density, \( g(z) = n(z)/n(z = 0) \). To define this function, we use the result that a sphere of gas will collapse and virialize once its density exceeds the value 1.686D(z)^{-1}p(z) (Peebles 1980). The mean density of the virialized gas is \( \Delta_\text{c}p(z) \), where \( p(z = 3H(z)^2/(8\pi G)) \), and an approximation of \( \Delta_\text{c} \) is \( \Delta_\text{c} \approx 178\Omega_{m0}^{0.45} \) (Eke et al. 1998), where \( \Omega_{m} = (1 + z)^2/\Omega_{m0}(1 + z)^3 + 1 - \Omega_{m0} \). Since clusters are the largest virialized objects in the universe, we take \( \Delta_\text{c}p(z) = g(z)n_{\text{c}}\Omega_{m0}^{0.55} \), and we assume that galaxies, halos, and clusters all share a universal \( g(z) \):

\[ g(z) = \frac{\Delta_\text{c}p(z)}{\Delta_\text{c0}p_0(0)} = (1 + z)^{-35} \times (1 + \Omega_{m0})^{0.55}. \]
In the following sections, the relations \( n_g(z) = g(z)n_{g,0}, \) \( n_{cl}(z) = g(z)n_{cl,0} \) will be used for the postshock magnetic field and \( pp \) optical depth.

### 4.1. Galaxy Mergers

The maximum energy of CRs accelerated in the merger shocks will also evolve with \( z \) due to the redshift dependence of the typical galactic radius and magnetic field characterizing the shocks. The magnetic field behind the shock in a galaxy merger is commonly parameterized as a fraction of the ram pressure (Kashiyama & Mészáros 2014), \( B_z^2 / 8\pi = \frac{1}{2} \epsilon \nabla n_p v_e^2 \sim \rho v_e^2. \) This implies a magnetic field

\[
B_z = \sqrt{4\pi \epsilon n_p g n_p g(z) v_e^2} \approx 14 \epsilon^{1/2} n_p^{1/2} g(z)^{1/2} \times \left( \frac{\nu_i}{300 \text{ km s}^{-1}} \right) \mu G. \tag{13}
\]

The magnetic field in the disk region is expected to be higher than that in the halo region. Although details depend on the geometry, for simplicity, we assume that a reasonably strong magnetic field is expected over scales between the galaxy radius \( R_g \) and a gas scale height \( h_g \), which is taken as the characteristic scale height \( h \sim (3h_g R_g^2 / 2)^{1/3} \) in this work. Then, the maximum CR energy is estimated (Drury 1983)

\[
\epsilon_p^{\text{max}} \approx \frac{1}{20} \epsilon n_p h_c \approx 1.3 \times 10^{16} \text{ eV} \left( \frac{B_z}{30 \text{ \mu G}} \right) \times \left( \frac{h}{3 \text{ kpc}} \right) \left( \frac{\nu_i}{300 \text{ km s}^{-1}} \right). \tag{14}
\]

The CRs are advected to the far downstream and produce neutrinos and \( \gamma \)-rays during the advection. In reality, one needs to calculate neutrinos and \( \gamma \)-rays from the postshock region especially when the \( pp \) optical depth in the CR acceleration region is dominant. The emissions occur during \( t_{\text{dyn}} \approx h / \nu_i \approx 9.8 \text{ Myr} \) (\( h / 3 \text{ kpc} \))\( (300 \text{ km s}^{-1} / \nu_i) \). In this work, for simplicity, we take the CR reservoir limit, in which the CRs mostly escape into the ISM and the neutrino and \( \gamma \)-ray production mainly occurs in the ISM.

After the CRs are accelerated in the shock, they will propagate in the host galaxy and cluster. In this process, neutrinos and \( \gamma \)-rays are generated from pions produced in inelastic \( pp \) collisions. The meson production efficiency is \( 1 - \exp(-f_{pp}) \), where \( f_{pp} = \epsilon n_p \sigma_{pp} g(z) \sum n_i \rho_i \) is the effective \( pp \) optical depth. In this expression, \( n_0, \rho_0 \) is the local gas density of the medium, for example, galaxies and clusters; \( \sigma_{pp} = \sigma_{pp}(\epsilon_p) \) is the \( pp \) cross section given by Kafexhiu et al. (2014); \( \rho_{p,0} = 0.5 \) is the inelasticity coefficient; and \( g(z) \) (see Equation (12)) represents the redshift evolution of the gas density.

Let us consider galaxies that are merging at \( z \). Inside the merged galaxy, \( f_{pp}' \) is determined by the time spent by the CRs undergoing \( pp \) collisions, which depends on the CR injection time and the diffusion time in the medium. The dynamical time is given by the third line of Equation (5), while the diffusion time is \( t_{\text{diff}} = h(z)^2 / (6D_g) \), where \( h(z) \) is the effective gas size at \( z \) and \( D_g \) is the diffusion coefficient in the galactic ISM gas. Here, we use a combined large and small angle diffusion expression as in Senno et al. (2015), \( D = D_L \left( \epsilon / \epsilon_{cg} \right)^{1/2} + \left( \epsilon / \epsilon_{cg} \right)^2 \), where \( D_L = \epsilon n_p / (6 \xi c_g) / 4 \) and \( \epsilon_{cg} \) is determined from \( r_L(\epsilon_{cg}) = l_c / 5. \) Here, \( r_L \) and \( l_c \) are the Larmor radius and coherence length in the galaxy environment, respectively. For local normal galaxies, the gas density in the disk is \( n_{g,0} \sim 1 \text{ cm}^{-3} \), whereas the average density in the galactic halo is smaller, \( n_{g,0} \sim 0.1 \text{ cm}^{-3} \). The magnetic field of local normal galaxies is \( \sim 4 \mu G \) and that of star-forming galaxies is \( \sim 6 \mu G \) (Keeney et al. 2006; Crocker 2012). For the density and magnetic field of merging galaxies, we take values higher than those of normal galaxies, since the galaxies may enter the starburst phase during the merger. Specifically we adopt a mean value \( n_{g,0} = 1 \text{ cm}^{-3} \). Thus, we have

\[
t_{\text{diff}} \approx 3.2 \times 10^5 \text{ years} \left( \frac{h(z)}{3 \text{ kpc}} \right) \left[ (\epsilon / \epsilon_{cg})^{1/2} + (\epsilon / \epsilon_{cg})^2 \right]^{-1},
\]

where

\[
\epsilon_{cg} \approx 1.7 \times 10^{16} \text{ GeV} \left( \frac{h(z)}{3 \text{ kpc}} \right) \left( \frac{B_z}{30 \text{ \mu G}} \right). \tag{16}
\]

Calculations of the neutrino and \( \gamma \)-ray emission depend on details of the spatial extension and time evolution of the shock region and its surrounding environment. The latter is also modified by the shock, star formation, and outflow. For simplicity, we treat a double-galaxy merger system as one CR reservoir for the injection by the merger shock, which is conservative since there should also be the emissions from the accelerator. A similar treatment for neutrino sources with active accelerators is used in the galaxy cluster model (Fang & Murase 2018; Murase et al. 2008). Then, the effective \( pp \) optical depth is estimated to be \( f_{pp}' = \epsilon n_p g n_{pp} \min[\delta_{\text{dyn}}, t_{\text{diff}}] / (30 \text{ kpc}) \), where \( B_{pp} R_{pp}^2 \sim GM_{pp}^2 / R_g \), that is, \( B_g \propto \rho g R_g \times g(z) R_g(z) \).

The typical galactic radii evolve with redshift \( z \), and considering the merger history of galaxies, it is apparent that the mean radii of galaxies at \( z \) should be smaller than \( R_{g,0} / (1 + z) \), where \( R_{g,0} \approx 10 \text{ kpc} \) is the radius of local Milky Way–like galaxies. Shibuya et al. (2015) studied the redshift evolution of the galactic effective radius \( r_e \) using Hubble Space Telescope (HST) samples of galaxies at \( z = 0–10 \), finding \( r_e \propto (1 + z)^{-1.0} - (1 + z)^{-1.3} \) with \( r_e \propto (1 + z)^{-1.0} \) as a median. Hence, in this paper, we assume that the average galactic radius evolves with respect to \( z \) as \( R_g = R_{g,0} (1 + z)^{-1.0}. \)

As for the scale height \( h_g(z) \), based on the surface photometry analysis of edge-on spiral galaxies, like NGC 4565, it has been shown that the scale height of gas in local disk galaxies is approximately \( h_g \approx 300–400 \text{ pc} \) (Bahcall & Soneira 1980; Van Der Kruit & Searle 1981). Later studies of NGC 891 (Kylafis & Bahcall 1987), NGC 5097 (Barnaby & Thronson 1992), and so on also agree with this estimate. Considering that a merger can lead to entering a star-forming phase, we assume \( h_g \approx 500 \text{ pc} \), and we assume the same redshift dependence as for \( R_g \), that is, \( h_g(z) = (1 + z)^{-1.0} h_g(0). \)

Then we take \( h = (3h_g h_z)^{2/3} \).

### 4.2. Interactions in the Host Cluster and Cluster Mergers

After escaping the galaxy, the CRs may continue to collide with the gas of the host cluster, where \( t_{\text{diff}} = R_d(z)^2 / (6D_d) \). Here, we assume a magnetic field \( B_{d,0} \approx 1 \mu G \) with a...
coherence length $l_{c,1} \approx 30 \text{kpc}$. This implies $\varepsilon_{c,1} \approx 5.6 \times 10^{9} \text{GeV}$. For a cluster of mass $10^{15} M_{\odot}$, the virial radius is $R_{c,1} = \left(3 M/(4\pi\rho_{c,1})\right)^{1/3} \approx 2.1 \text{ Mpc}$. Since $R_{c,1}$ is the approximate boundary of clustered/correlated galaxies, it should have the same redshift dependence as $n_{c,1}$. Using the SC approximation, we obtain $R_{c,1} \propto (1 + z)^{(3-\gamma)/\gamma}$. Similarly, we can calculate the diffusion time in clusters as $t_{\text{diff},1} = 1.2 [\varepsilon_{c,1}/(z)^{(3-\gamma)/\gamma}]^{1}\text{ Gyr}$. Assuming that the injection time of CRs ($t_{\text{inj}}$) at redshift $z$ is the cluster age (of order the Hubble time) $t_{\text{H}}(z)$, likewise we obtain the optical depth $\tau_{pp} = \kappa_{pp} C_{g} \rho_{c,1} t_{\text{inj}}$, where $\rho_{c,1}$, the intercluster gas density, is assumed to have the typical value $n_{c,1} \rho_{c,1} \approx 10^{-4} - 10^{-2} \text{ cm}^{-3}$ (Croston et al. 2008), which can be higher in cooling core clusters. The magnetic field may also depend on $z$ as $B \propto \rho_{c,1}^{1/2} \varepsilon_{c,1}(z)$.

Halo mergers will also lead to galaxy group and galaxy cluster mergers, after some halos have grown above a certain size that may be taken to be roughly of order $M_{\odot} \approx 10^{13} M_{\odot}$. We simplify the calculations as follows. For low-mass mergers, we expect that the $pp$ interactions occur mainly in gas with an ISM density characteristic of galaxies, while for high-mass mergers the $pp$ interactions occur mainly in gas with an IGM density characterizing the intracluster gas. In addition, there will be a component of $pp$ interactions due to low-mass merger CRs that escape from the colliding galaxy system into the IGM. Thus, we expect that the all-flavor neutrino production rate consists of a galaxy part $\varepsilon_{e} Q_{(g)}$ and a cluster/group part $\varepsilon_{e} Q_{(c)}$, plus a weaker galaxy-cluster term:

$$
\varepsilon_{e} Q_{(g)}^{(g)} = \frac{1}{2} (1 - e^{-tr_{g}}) \varepsilon_{e} Q_{p}^{(LM)}
$$

$$
\varepsilon_{e} Q_{(c)}^{(c)} = \frac{1}{2} [1 - e^{-tr_{c}}] \varepsilon_{e} Q_{p}^{(HM)}
$$

$$
+ \eta (1 - e^{-tr_{c}}) e^{-tr_{c}} \varepsilon_{e} Q_{p}^{(LM)}
$$

(17)

where the energies of the neutrinos and CR protons are related by $\varepsilon_{e} \approx 0.05 \varepsilon_{p}$. Note that the luminosity density evolution of neutrinos and $\gamma$-rays is different from that of CRs in general. In the first line of Equation (17), $\varepsilon_{e} Q_{p}^{(LM)}$ is the CR input rate (see Equation (6)) from galaxy mergers in low-mass (LM) halos, for example, $10^{10} M_{\odot}, 10^{13} M_{\odot}$. In the second line, $\varepsilon_{e} Q_{p}^{(HM)}$ is the CR input rate of the high-mass (HM) halo mergers, in the interval $[10^{13} M_{\odot}, 10^{15} M_{\odot}]$. The factors $\frac{1}{2} (1 - e^{-tr_{g}})$ are the neutrino luminosity density from CRs originating from mergers of mass $(j)$ in gas of density $i$. For our fiducial parameters, these two components constitute the largest fraction of the neutrino budget. Nevertheless, for completeness, we have included in the third line of Equation (17) the subdominant effect due to the CRs produced in galaxy mergers that may escape the host galaxies and collide with intracluster gas to produce neutrinos. (This can be important only if the $pp$ interactions in galaxy mergers are inefficient.) We introduce a parameter $\eta$ to represent the fraction of galaxy mergers that occur inside clusters, which lead to some CRs escaping into the gas halo. This can occur preferentially at lower redshifts. Since the boundary between LM and HM is ambiguous and the fraction $\eta$ can change with redshift, this parameter is very uncertain and may conservatively be estimated as being between 0.1 and at most 0.5. Fortunately, the contribution of this higher-order third component depending on $\eta$ is small compared to the first two components in Equation (17), due to the factor $e^{-tr_{c}}$. At $z = 1$, the ratio between the third line and the first line is $\lesssim 10\%$ even if $\eta$ is assumed to be unity, and it is increasingly negligible at higher redshift since $f_{p}^{2}$ increases as the gas density increases. Therefore, the exact value of $\eta$ does not significantly influence our final results.

5. Diffuse Neutrino and $\gamma$-ray Spectra

With the above, we are able to determine the CR energy input rate, $\varepsilon_{e} Q_{p}^{(g)}$ and $\varepsilon_{e} Q_{p}^{(c)}$. Figure 3 shows the CR input power over the whole mass interval $10^{10} - 10^{15} M_{\odot}$, as a function of $z$ as well as the LM and HM components of the ($\sigma_{0} = 300$, $\varepsilon_{p}^{(c)}$) scenario. As can be seen from the redshift distribution of the CR energy input, using a redshift-evolving gas fraction $\varepsilon_{g}^{(c)}$, a significant fraction of this occurs at redshifts $z \gtrsim 3$, above which a significant $\gamma$-ray attenuation of the accompanying high-energy $\gamma$-rays at $\gtrsim 20-30$ GeV energies can be expected (e.g., Chang et al. 2016; Xiao et al. 2016). In addition, from Figure 3, we find that the high-mass and low-mass components are comparable in local mergers, implying that the cluster/group merger contribution is also important. Also, the galaxy merger contribution to the CR luminosity density is more important at $z \gtrsim 2$.

Given the neutrino input rate, the all-flavor neutrino flux can be expressed as (Murase et al. 2016)

$$
\varepsilon_{e} \Phi_{\nu} = \frac{c}{4\pi} \int \varepsilon_{e} Q_{p}^{(g)} + \varepsilon_{e} Q_{p}^{(c)} \frac{dt}{(1 + z)} \frac{dz}{dz}
$$

(18)

Based on the branching ratio between charged and neutral pions, the initial diffuse $\gamma$-ray energy spectrum is expected to be given by $\varepsilon_{e} \Phi_{\gamma} = \varepsilon_{e} \Phi_{\nu} |_{\gamma} = 0.5 \varepsilon_{e} \Phi_{\nu}$. Since, however, the high-energy $\gamma$-rays can annihilate with lower energy photons, such as those from the extragalactic background light (EBL) and the cosmic microwave background (CMB), we introduce an
Figure 4. Left panel: neutrino (all flavor) and γ-ray fluxes from halo mergers with redshift-evolving gas fraction $\xi_{g}^{\mathrm{env}}$. $R_{A}=10$ kpc, $H_{A}=500$ pc. The shock velocity is obtained using $\beta_{c}(z)$ and $\sigma_{0}=300$. The magenta line is the neutrino spectrum, while the green line is the corresponding γ-ray spectrum. Galaxy and cluster contributions to the neutrino flux are illustrated as the dashed and dash-dotted lines, respectively. Right panel: same as left panel except $\sigma_{0}=500$ is utilized for $\nu_{e}$.

attenuation factor $\exp[-\tau_{\gamma}(\varepsilon, z)]$ to the integration, where $\tau_{\gamma}$ is the $\gamma\gamma$ optical depth at redshift $z$. In this paper, we use the optical depth provided by Finke et al. (2010) and Inoue et al. (2013) for low-redshift ($z \ll 5$) and high-redshift ($z > 5$) inputs, respectively. The attenuated γ-ray flux is then

$$e_{\gamma}^{2} \Phi_{\gamma} = \frac{c}{4\pi} \int \left[ \frac{2}{3} \left( e_{\gamma}^{(g)} + e_{\gamma}^{(e)} \right) \frac{d\sigma}{dz} \right] \times \exp[-\tau_{\gamma}(\varepsilon, z)] \, dz \tag{19}$$

with $\varepsilon_{\gamma} = 10 \varepsilon(1 + z)$. In addition, the electron–positron pairs produced in the $\gamma\gamma$ annihilations will subsequently scatter off the ambient diffuse photon backgrounds, leading to an electromagnetic cascade that in part compensates for the attenuation, while reprocessing the photon energy toward lower energies, which can be detected by, for example, the Fermi-LAT instrument. In this paper, for simplicity, we use the universal form for the resulting cascaded γ-ray spectrum given by Berezhinsky & Smirnov (1975; see also, e.g., Murase & Beacom 2012; Senno et al. 2015):

$$\frac{dN}{d\varepsilon_{\gamma}} \propto G(\varepsilon_{\gamma}) = \begin{cases} (\varepsilon_{\gamma}^{0})^{-1/2} & \varepsilon_{\gamma} \leq \varepsilon_{\gamma}^{0} \\ (\varepsilon_{\gamma}^{0})^{-1} & \varepsilon_{\gamma}^{0} < \varepsilon_{\gamma} < \xi_{\gamma}^{\mathrm{cut}} \\ (\varepsilon_{\gamma}^{\mathrm{cut}})^{-2} & \varepsilon_{\gamma} \geq \varepsilon_{\gamma}^{\mathrm{cut}} \end{cases}, \tag{20}$$

where $\varepsilon_{\gamma}^{\mathrm{cut}}$ is defined by $\tau_{\gamma}(\varepsilon_{\gamma}^{\mathrm{cut}}, z) = 1$ and $\varepsilon_{\gamma}^{0} = 0.0085 \text{ GeV}(1 + z)^{2} (\varepsilon_{\gamma}^{\mathrm{cut}}/100 \text{ GeV})^{2}$.

The all-flavor diffuse neutrino and γ-ray fluxes are plotted in Figure 4, together with the IceCube-observed astrophysical neutrinos. The red points and cyan points correspond to the all-flavor averaged neutrino flux (Aartsen et al. 2015, 2016) and the 6-year high-energy starting events (HESE; Aartsen et al. 2017), respectively. The Fermi-LAT observed total EGB (Ackermann et al. 2015) is shown by the blue points. The yellow area is the best fit to the upcoming muon neutrinos scaled to three-flavor. Figure 4 shows the results for an assumed redshift-dependent gas fraction $\xi_{g}^{\mathrm{env}}$, as illustrated in Figure 4(a) for $\sigma_{0}=300$ and in Figure 4(b) for $\sigma_{0}=500$, showing the effect of the corresponding different shock velocities $\nu_{s}$. In each figure, the magenta line represents the neutrino flux, while the green line illustrates the corresponding γ-ray flux after cascading down. The galaxy and cluster contributions to the neutrino flux are plotted in dashed lines and dash-dotted lines. The nonblazar (Ackermann et al. 2016) component of the unresolved extragalactic gamma-ray background is shown as the pink area.

For illustration purposes, we consider next the corresponding results using the redshift-independent gas fraction $\xi_{g} = 0.05$, which are shown in Figures 5(a) and (b). The comparison between the galaxy and cluster components indicates that the high-energy neutrinos are dominantly produced by the propagation of CRs in the clusters. This is a consequence of the rapid redshift evolution of the galaxy radius, since the size of the host galaxy limits the maximum CR energy as well as the neutrino production efficiency by restricting the diffusion time. In addition, a redshift-dependent $\xi_{g}^{\mathrm{env}}$ boosts the CR budget to a relatively higher redshift ($z \approx 3$), as can be seen from the red line in Figure 3, which as was expected leads to a reduction in the γ-ray flux. From these figures, we can also see that even with the moderate sensitivity of the results to the parameters $\tau_{b}$ and $\sigma_{0}$, the results can broadly fit a significant fraction of the IceCube data without violating the nonblazar EGB. Conversely, the γ-ray and neutrino fluxes are significantly constrained in this scenario, indicating that the halo and galaxy mergers can be regarded as promising sources of neutrinos in the context of multimessenger studies.

6. Discussion

In this work, we investigated the contribution of halo mergers to the diffuse neutrino and γ-ray backgrounds, and we tested whether the nonblazar diffuse γ-ray background Fermi constraint is violated. Our results differ from previous work by Kashiyama & Mészáros (2014) in that we studied both galaxy and cluster/group mergers out to higher redshifts, up to $z \approx 10$, by considering the redshift evolution of the average galactic radius, the shock velocity, and the gas content inside the halos, as well as the galactic/intergalactic magnetic fields.

The redshift evolution of galaxy radius implies that more protogalaxies exist, or equivalently more mergers at higher...
redshift. In fact, the merger rate calculated using our approximate approach verifies this conjecture, as well as being consistent with the Illustris simulations (Rodriguez-Gomez et al. 2015). Also, our estimates of the gas fraction $\xi_{\text{gas}}$ based on the correlation between the galactic gas content and the SFR show that the gas in high-redshift halos is relatively denser than in the current-epoch halos. The net effect is that high-redshift halo mergers can contribute a significant fraction of the CRs that are capable of producing high-energy neutrinos, as shown in Figure 3. This is crucial since the accompanying $\gamma$-ray photons in the ensuing $pp$ collisions at high redshifts can be sufficiently absorbed via $\gamma\gamma$ annihilations against CMB and EBL photons. In both cases with $\xi_{\text{gas}}$, our results indicate that high-redshift galaxy/halo mergers can explain a large fraction of the IceCube-observed diffuse neutrinos up to petaelectron-volts, with an accompanying $\gamma$-ray diffuse observed flux that is below the nonblazar Fermi constraints.

We note that, according to our calculation, the diffuse flux of CRs that survives from energy losses via $pp$ collisions is less than $10^{-8}$ GeV cm$^{-2}$ sr$^{-1}$ s$^{-1}$, which is lower than the observed CR flux around the knee or sub-ankle energy.

The CR acceleration efficiency $\epsilon_p$ is expected to be $\sim 0.1$ based on the diffusive shock acceleration theory. The redshift dependencies of gas fraction $\xi_{\text{gas}}$ and galaxy radius are relatively well modeled from current theories and observations, so our scenario can put a tighter constraint on the shock velocity of galaxy mergers. However, there are large uncertainties in the model. For example, the maximum energy depends on the magnetic field strength, which is highly uncertain. On the other hand, the fiducial value of $\sim 10$ PeV is not far from the knee energy at $\sim 3$ PeV, so our assumption is reasonable. One of the most important uncertainties is caused by the shock velocity. Our fiducial parameters ($\sigma_0 = 300$ with $\xi_{\text{gas}}$) imply a lower neutrino flux compared to the observations. This could be overcome by assuming a higher velocity with a stronger magnetic field. Or it may be possible to achieve the IceCube flux at $\sim 0.1$ PeV without exceeding the Fermi constraint by increasing the cluster contribution. However, the cluster contribution is more uncertain. Nonthermal emissions from merging or accreting clusters have been studied by various authors (e.g., Fujita & Sarazin 2001; Berrington & Dermer 2003). The Mach number of shocks on the cluster scales is so low, due to the high temperature of the intracluster medium, that the shock may not be strong enough to have a hard spectrum of $s \sim 2$.

We note that, in addition to mergers, also cluster accretion shocks and powerful jets from radio-loud AGNs can contribute to CR acceleration inside the clusters/groups, as considered in the previous literature (e.g., Murase et al. 2013; Fang & Murase 2018, and references therein). One of the generic features of the CR reservoir scenario is that different possibilities for CR acceleration are not mutually exclusive, and additional contributions from various CR accelerators may enhance the neutrino flux. Another CR source that can be associated with galaxy mergers is that the compression of the ISM gas can trigger an intense starburst. As discussed by Charbonnel et al. (2011), two processes in colliding galaxies could induce a starburst: radial gas inflows can fuel a nuclear starburst, while gas turbulence and fragmentation can drive an extended starburst in clusters. Such intense star formation can naturally lead to the injection of CRs from the ensuing massive stellar deaths, including from SNRs and HNRs. In addition, CRs may also be injected from disk-driven outflows and weak jets from AGNs (Murase et al. 2014; Tamborra et al. 2014; Liu et al. 2017). The CR contributions from these sources, which would be additional to CRs from the mergers considered here, are significantly model-dependent, and we do not attempt here a quantification of their relative importance.

One important factor that may influence the final results is the CR power-law index $s$, since the factor $\varepsilon^{2-s}$, the maximal CR energy, and a new $C = ((\varepsilon_p^\text{max})^{2-s} - (\varepsilon_p^\text{min})^{2-s})/(2 - s)$ are required to correct Equation (6) when $s$ deviates from 2.0. As presented in Equation (6), we assume that the shock is nonradiative and infinitesimally thin, and hence the Fermi first-order acceleration in the strong-shock limit implies $s = 2$. However, a finite width of the shock can steepen the spectrum to $s \gtrsim 2.0$, while a radiative shock would produce a CR spectrum with a power-law index lower than 2.0. For radiative shocks, Blandford & Eichler (1987) showed that the power-law index of the accelerated CRs is $s = (\phi + 2)/(\phi - 1)$ where $\phi$ is the compression ratio. Moreover, Yamazaki et al. (2006) assume $\phi = 7$ and $s = 1.5$ as fiducial values when studying the radiative cooling of SNR shocks. Note that the compression ratio for radiative shocks with an isothermal adiabatic index $\gamma = 1$ can be written as $\phi = M^2 \gg 1$, where $M$ is the upstream
Mach number. We assume thus $s = 1$ in this extreme case. To illustrate how the neutrino spectra are affected by the radiative cooling or the width of the shocks, we plot in Figure 6 the neutrino fluxes for four cases $s = 2.2, 2.0, 1.5,$ and $1.03$, which correspond to $\phi = 3.5, 4, 7,$ and $100$, respectively. As can be seen, a higher CR power-law index (lower $s$) will produce more high-energy neutrinos. Thus, in principle, a mildly radiative-cooling shock ($1.5 \lesssim s \lesssim 2$) can more easily achieve the high-energy neutrino flux in the range $10^{-1}$ to $10^{-5}$ TeV to electroweak neutrinos. On the other hand, $s \gtrsim 2.1$–$2.2$ is disfavored because of the damping factor $\varepsilon^{2-s}$, which is consistent with previous work (Murase et al. 2013). Note that the hard spectrum is expected for the cold gas environment that would be valid in sufficiently low-mass halo mergers. If the temperature is so high, the Mach number is expected to be low, as expected for cluster mergers. In this case, the spectrum is softer for massive clusters, and the details are beyond the scope of this work.

Additional contributions may arise from the galaxies moving through the cluster or dark matter halo, as their hypersonic peculiar motion will result in a shock as they plow through the intracluster gas, which as a result may also contribute to the diffuse $\gamma$-rays and neutrinos. Supposing as an extreme case that the loss of the galaxies’ kinetic energy due to the gravitational drag is completely converted into CR energy, we estimate a CR energy budget of

$$e_{\nu} \Phi_{\gamma}^{(GM)} = e_{\nu} \Phi_{\gamma}^{-1} \int dM_{h} \frac{4\pi G^{2}M_{h}^{2}}{v_{y}} \frac{dN}{dM_{h}} \ln \left( \frac{R_{\text{cl}}}{R_{b}} \right),$$

which is three orders of magnitudes lower than halo mergers estimated in the previous sections, because of the tenuous intergalactic gas density. Hence, these shocks contribute only a relatively small amount of diffuse neutrinos and are negligible compared to the mergers.

7. Summary

In summary, we found that the CR luminosity density by halo mergers can be comparable to that from starburst galaxies, which can be expected from galaxy mergers. In particular, the CR input from galaxy mergers and cluster/group mergers is comparable in the local universe, and the former is more important at higher redshifts, $z \gtrsim 2$ (see the dashed and dash-dotted lines in Figure 3). This emphasizes the importance of our results for CR reservoir models. We have considered the neutrino and $\gamma$-ray production in galaxy–galaxy and cluster/group merger environments and found that such mergers could explain a large portion of the IceCube diffuse neutrino flux.

Figure 6. Neutrino fluxes for different compression ratios and CR power-law indices. The black, magenta, blue, and green lines correspond to the power-law indices $s = 2.2, 2.0, 1.5,$ and $1.03$. The Astrophysical Journal, 857:50 (11pp), 2018 April 10 Yuan et al.

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Appendix

The Merger Rate

In this section, we present a comparison between our halo merger rate and the Illustris simulations (Rodriguez-Gomez et al. 2015). In our calculation, we assign a mean merger probability $P(M, z) = \exp(-\tau_{\text{merge}}/\tau_{\text{age}})$ to each dark matter halo during $\tau_{\text{age}}$. Here, $\tau_{\text{merge}}$, which can be obtained from the second equation of Equation (5), averages all possible mass ratios, that is, $\zeta \in (0, 1]$. In our calculation, we do not need to use the cumulative merger rate over mass directly; instead the factor $\frac{dN}{dM} \frac{P(M, z)}{\tau_{\text{age}}}$ in the integrand of Equation (6) is used to illustrate the number of mergers for a halo with mass $M$ and at redshift $z$. However, in order to compare our results with the simulations, it is worthwhile to estimate the average cumulative
The merger rate is shown in Figure 7, where the blue line corresponds to $10^{10} M_{\odot} < M_0 < 10^{13} M_{\odot}$, and the red line corresponds to $10^{12} M_{\odot} < M_0 < 10^{15} M_{\odot}$. In both cases, the mass ratio covers the entire interval as in the middle equation of Equation (5), which is integrated over $\zeta$ from 0 to 1. The merger rate given by Illustris simulations is shown in the lower panel. One can compare our results with the solid black lines in the right panel of Figure 2 in Rodriguez-Gomez et al. (2015), since the increase in the merger rates given by simulations (as shown as colored lines) seen at low redshifts is due to a limitation of the splitting algorithm. As can be seen, our merger rate in the same mass interval is comparable to the merger rate in the right panel with the mass ratio $\geq 1/1000$. Considering that we are using a totally different method and this approach is primarily designed to evaluate the merger probabilities of halos of various masses and at different redshifts, the moderate degree of discrepancy can be considered acceptable.

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**Figures**

Figure 7. Merger rates from Equation (22). The blue line represents the whole mass range ($10^{10} M_{\odot} < 10^{13} M_{\odot}$), and the red line corresponds to $10^{12} M_{\odot} < M_0 < 10^{15} M_{\odot}$. The merger rate using our approach:

$$ \mathcal{R}(z) = \frac{\int dM P(M_{\odot}, z) dM}{\int dM_{\odot} dM} $$

The merger rate is shown in Figure 7, where the blue line represents the whole mass range ($10^{10} M_{\odot} < 10^{13} M_{\odot}$) and the red line corresponds to $10^{12} M_{\odot} < M_0 < 10^{15} M_{\odot}$. In both cases, the mass ratio covers the entire interval as in the middle equation of Equation (5), which is integrated over $\zeta$ from 0 to 1. The merger rate given by Illustris simulations is shown in the lower panel. One can compare our results with the solid black lines in the right panel of Figure 2 in Rodriguez-Gomez et al. (2015), since the increase in the merger rates given by simulations (as shown as colored lines) seen at low redshifts is due to a limitation of the splitting algorithm. As can be seen, our merger rate in the same mass interval is comparable to the merger rate in the right panel with the mass ratio $\geq 1/1000$. Considering that we are using a totally different method and this approach is primarily designed to evaluate the merger probabilities of halos of various masses and at different redshifts, the moderate degree of discrepancy can be considered acceptable.
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